MEASURING POLARIZATION OF THE
COSMIC MICROWAVE BACKGROUND WITH
THE SOUTH POLE TELESCOPE
POLARIZATION EXPERIMENT

by

JAMES TOMLINSON SAYRE

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Dissertation Adviser: Dr. John Ruhl

Department of Physics
CASE WESTERN RESERVE UNIVERSITY

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We hereby approve the thesis/dissertation of

James Sayre

candidate for the **Doctor of Philosophy** degree *.

(signed) John Ruhl
(chair of the committee)

Corbin Covault

Glenn Starkman

Chris Mihos

(date) July 2, 2014

*We also certify that written approval has been obtained for any proprietary material contained therein.
For my family, whose unwavering support and encouragement have given me the confidence and determination to reach this point.
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Measuring Polarization of the Cosmic Microwave Background with the South Pole Telescope Polarization Experiment

by

James Tomlinson Sayre

Abstract

The South Pole Telescope Polarization experiment (SPTpol) is a camera consisting of 180 (588) pixels observing bands centered at 90 (150) GHz, installed on the South Pole Telescope in December 2012. It is a high-resolution, high-sensitivity instrument for mapping the polarized component of the Cosmic Microwave Background. In this thesis, we describe the development, testing, and deployment of transition edge sensor (TES) bolometers that make up the camera pixels, as well as the data analysis pipeline used to generate power spectra of the CMB. The tests used to measure various detector properties are described and their results displayed, and details of the analysis routines are explained. We conclude with preliminary results from SPTpol and a discussion of future directions for the experiment.
Chapter 1

INTRODUCTION

1.1 $\Lambda$CDM and the Cosmic Microwave Background

Observations of the cosmic microwave background (CMB) radiation are a pillar of modern observational cosmology. The discovery in 1964 by Penzias and Wilson of an excess antenna temperature of $3.5 \pm 1$ K at 4.08 GHz [40], interpreted by Dicke, Peebles, Roll, and Wilkinson [11] as the relic glow from a hot, dense epoch in the early Universe, gave powerful experimental support to the Big Bang theory of cosmology. A spectacular measurement of the CMB spectrum by the FIRAS instrument on the COBE satellite nearly thirty years later revealed its nearly perfect blackbody character and established its temperature as $2.726 \pm 0.01$ K [37]. COBE also made the first measurements of the minute anisotropy of the CMB brightness [46], the further study of which has driven the $\Lambda$CDM parametrization of the Big Bang model that currently stands in stunning agreement with a wide array of astrophysical observations.

Modern inflationary $\Lambda$CDM cosmology holds that the early Universe underwent a period of exponentially accelerating expansion, driven by the potential energy of a scalar field called the inflaton. For simplicity, inflation is generally described on the standard Big Bang timeline, with $t=0$ at the initial singularity (which is not
strictly necessary in the inflationary paradigm) as the “beginning of the universe”, and inflation ending at $t \approx 10^{-32}$ s. During inflation, the exponential expansion of space flattened its curvature, rarefied its contents, and enlarged quantum scales to sizes beyond the causal horizon. As a result, the post-inflation universe was nearly flat, free of exotic relics expected to be generated at the highest post big-bang energy scales, and largely homogeneous and isotropic within a given causal horizon – all properties we observe but are not generically predicted by the standard Big Bang model.

When the inflaton field reached a minimum in its potential, the Universe entered a period called reheating, when the energy density of the inflaton field was transferred to the constituents of the modern universe: $\Lambda$, or dark energy, which presently accounts for 68% of the total energy density ($\epsilon_0$) and is characterized primarily by its negative pressure that now drives an accelerating expansion of the universe; cold dark matter (CDM), non-relativistic massive particles that interact chiefly through gravitation and comprise 27% of $\epsilon_0$; and ordinary baryonic matter, photons, and neutrinos that combined make up the remaining 5%. After reheating, the $\Lambda$CDM model returns to the Big Bang paradigm of an expanding universe whose rate of expansion is determined by the constituent whose energy density dominates the overall energy density, $\epsilon_0$.

In the simplest models of inflation, quantum fluctuations in both the inflaton and graviton fields were expanded to super-horizon physical scales, where they were “frozen out” – no longer in causal contact and thus unable to evolve. When inflation ended, these frozen-out modes re-entered the horizon as the expansion of the universe again came to be driven by its radiation energy density. A causally-connected region prior to inflation expanded to become much larger than the present causal volume, so while frozen-out modes were disconnected during inflation, the differences between their thermodynamic properties were the result of quantum fluctuations on
a common background state. Super-horizon variations in the inflaton field value at reheating became fluctuations in the density of the particle fields, while the inflated graviton fluctuations re-entered the horizon as a background of gravitational waves. Details of post-inflation dynamics are beyond the scope of this work, but the ansatz of a young universe with Gaussian, largely scale-invariant, and adiabatic scalar and tensor fluctuations on a nearly homogeneous, isotropic background of standard model particles is the basis of ΛCDM modeling of the ensuing evolutionary dynamics.

When the universe was \( \sim 380,000 \) years old, it underwent a phase change, known as the epoch of recombination, that generated the CMB we observe today. Prior to recombination, photons and baryonic matter were thermally bound in a hot, dense plasma that was thermodynamically dominated by the photon component. The Jeans length, which describes the minimum size of a fluctuation in a fluid with density \( \rho \) and sound speed \( v_s \) that can collapse under its own self-gravity,

\[
\lambda_J = v_s \sqrt{\frac{\pi}{G\rho}},
\]

was set by the photon sound speed, \( v_s \sim 0.6c \), which dominated the energy density of the plasma. for the plasma was larger than the Hubble length, \( \lambda_J = 3c/H > c/H \). Over-densities in the plasma from scalar fluctuations during reheating could not gravitationally collapse and therefore evolved only as stable acoustic oscillations.

On scales up to the sound horizon, \( v_s/H \), acoustic oscillations resulted from the gravitation drawing plasma into overdense regions until the resulting increased radiation pressure caused it to explode outwards. Dark matter thermally decoupled from the plasma in the very early universe, but until \( t_{RM} \sim 70,000 \) years, the overall energy density of the universe, \( \epsilon_0 \), was dominated by its photon component. As a result, the evolution of dark matter density fluctuations was suppressed by the radiation-dominated expansion rate of the universe. At \( t_{RM} \), the CDM energy density became
the dominant component of $\epsilon_0$, at which point density fluctuations in the dark matter field could collapse evolve under the influence of their self-gravity. Though baryons were bound into acoustic oscillations by their thermal coupling to photons, the plasma density was gravitationally coupled to an evolving background of dark matter density anisotropies.

Recombination occurred when the temperature of the universe dropped low enough for stable atoms to form and as a result, the mean free path of photons increased to beyond the Hubble length in a short period of time. The CMB we observe today is overwhelmingly composed of photons that last scattered from an electron in the primordial plasma during recombination, red-shifted from their emitted wavelength by a factor $z \sim 1100$. It is essentially a frozen image of the plasma at $t \sim 380,000$ years. This so-called “surface of last scattering” captures the state of the universe before baryonic matter was able to collapse gravitationally. Fluctuations in plasma temperature from acoustic oscillations, as well as the influence of background dark matter gravitational potentials, are imprinted on the CMB image as minute spatial fluctuations in its brightness temperature. An even smaller polarized component, which is described in Section 1.2, also results from the local anisotropies as seen by scattering electrons at recombination. Tracing physics back to the end of inflation from the image of the infant universe at last scattering, we can match our observations with models describing the composition and dynamics of the post-inflation, pre-recombination universe.

In addition to providing a snapshot of the infant universe, the CMB provides a “backlight” to the subsequent evolution of cosmic structure. CMB photons observed today were emitted at last scattering 13.8 billion years ago when the universe was still largely homogeneous. They have propagated through evolving gravitationally-bound structures ever since, and small perturbations arising from interactions with intervening structure offer an important means of constraining the growth history of
1.2 Polarization in the CMB

Polarization in the CMB arises from Thompson scattering by electrons in the presence of a local temperature quadrupole during recombination. Here I will briefly describe the dominant effects in the primordial plasma that produce the temperature quadrupoles responsible for CMB polarization. For an excellent intuitive and thorough treatment of the physical origins of the polarization signal, see [23].

Acoustic oscillations that were frozen in at recombination are the principal contributor to temperature anisotropies in the CMB. Prior to recombination, photons and electrons were in local thermal equilibrium, so scattering events occurred in a largely isotropic temperature environment. The majority of CMB photons are thus unpolarized, as their final scattering was influenced by a mostly uniform local plasma temperature. As the photon mean free path approached the Hubble length during recombination, scattering electrons could “see” temperatures at distances up to the causal horizon, where the plasma had a different local equilibrium temperature. The quadrupolar projection of the resulting temperature anisotropy seen by a scattering electron induces a linear polarization in the emitted photons. This projection depends on gradients in the scalar temperature field, $\nabla \phi$, so by definition there is no rotational component in an ensemble of such quadrupoles, or $\nabla \times \nabla \phi \equiv 0$.

Bulk flow of the plasma also results in an apparent anisotropic temperature in the frame of a flowing electron. To illustrate this, we can consider an electron flowing in the direction $\hat{k}$ into an overdense region at recombination. In its rest frame, the flow toward the center of the potential imparts a net blue shift to radiation in the plane normal to $\hat{k}$ as the projection of plasma motion in that plane points towards the electron. The apparent temperature in the $+\hat{k}$ direction is lower due to a red
shifting of faster-inflowing plasma closer to the bottom of the potential well. Similarly, radiation from the $-\hat{k}$ direction is red-shifted due to the electron’s own motion. The resulting quadrupole causes the scattered electron to be polarized along the direction of motion. In the reverse case of outflow from an overdensity, the signs switch and the resulting photon polarization is perpendicular to the direction of motion. The key result of this picture is that when the scalar perturbation field is projected onto the surface of last scattering, the polarization of emitted photons is either aligned or perpendicular to the velocity vector. Since the fluid flow is along the same gradients of a scalar potential field $\nabla \phi$, responsible for the static temperature quadrupoles, the polarization field resulting from bulk fluid flow into and out of plasma over-densities is also irrotational.

In addition to the background CDM gravitational potentials, inflation predicts a stochastic background of gravity waves propagating through the primordial plasma as a result of quantum gravitational effects driven to super-horizon scales during inflation. Again consider a single electron, this time initially in thermal equilibrium with its surroundings, perturbed by a tensor gravity wave propagating along the line of sight to a modern observer. The perturbation of the passing wave induces a quadrupole in the plane transverse to its propagation direction due to “squeezing” of the plasma along one axis while simultaneously “stretching” it along the other. While the projected temperature quadrupole generating a polarized photon is equivalent to one produced from a scalar perturbation, the symmetry argument associated with the underlying source of scalar fluctuations does not hold for tensor perturbations. Thus, over the whole sky the combined polarization pattern from an ensemble of tensor-induced quadrupoles is composed of both gradient-like and curl-like patterns.

The cosmological principle that underpins ΛCDM dictates the CMB is fully described by its statistical properties, and that those properties are isotropic so there is no preferred direction or location in the Universe. Because directionality is unim-
portant, it is conventional to describe the CMB in terms of its statistics in harmonic space rather than real space. A convenient basis for such a description is the set of spherical harmonics, \( Y_{\ell m} \), onto which a scalar function, \( f \), defined on the surface of a sphere can be decomposed by
\[
f(\theta, \phi) = \sum_{\ell} \sum_{m=-\ell}^{\ell} a_{\ell m} Y_{\ell m}(\theta, \phi).
\]
The \( 2\ell + 1 \) orthogonal \( Y_{\ell m} \) basis functions for each \( \ell \) are different distributions of fluctuations having a characteristic angular extent, indexed by \( \ell \). In the case of a Gaussian random field like the CMB, the average amplitude of the \( a_{\ell m} \) coefficients for a given \( \ell \) are zero, since otherwise would imply a preferred direction in the CMB and violate the cosmological principle. The variance of the \( a_{\ell m} \) distribution at a particular \( \ell \),
\[
C_\ell = \frac{1}{2\ell + 1} \sum_{m=-\ell}^{\ell} |a_{\ell m}|^2,
\]
fully describes the statistics of a Gaussian random scalar field of fluctuations as a function of their spatial extent, \( \ell \). The \( C_\ell \) spectrum, or power spectrum, is thus the primary means of measuring and analyzing the physics of the early universe through measurements of the CMB.

Describing the CMB at a particular point on the sky in the \( \hat{n} \) direction from the origin of the observer requires three values: photon temperature, \( T \), the magnitude of net linear polarization, \( P \), and its direction \( \psi \), measured as a rotation about \( \hat{n} \) relative to a consistent zero angle. The natural basis for measuring these quantities are the Stokes parameters, \( I, Q, \) and \( U \), where \( I \) is the total intensity of incoming radiation, and \( Q \) and \( U \) are an orthogonal basis for unambiguously describing the polarization direction. Note that circular polarization (Stokes V) is not generated by Thomson scattering and is therefore generally taken to be zero in the primordial CMB. Figure 1.1 illustrates the construction of the Stokes basis for observing the CMB, using the conventions employed by SPTpol.

Temperature is a scalar function on the sphere, fully described by \( T(\hat{n}) \propto I(\hat{n}) \),
Figure 1.1: Stokes parameters on the surface of a sphere
Viewed from the origin of the coordinate system \((\theta, \phi)\), the CMB at a point in the \(\hat{n}(\theta, \phi)\) direction can be fully characterized by Stokes I, Q, and U defined on a plane, \(\varphi\), normal to \(\hat{n}\). Intensity, \(I\) is just the total absorbed power, while \(Q\) and \(U\) form an orthogonal basis for unambiguously describing the amplitude and alignment of linear polarization. Thomson scattering does not produce circular polarization, Stokes \(V\), so it is neglected here. Note that the spherical unit vector definitions are appropriate for viewing from the South Pole in standard IAU coordinates.
where \( \hat{n} \) is a direction specified by spherical coordinates \((\theta, \phi)\). The value \( T(\hat{n}) \) does not change under a rotation, \( \alpha \), about \( \hat{n} \), i.e. it has spin 0 symmetry. Thus, it is natural to write the temperature field as \( T(\hat{n}) = \sum_{\ell} \sum_{m=-\ell}^{\ell} a_{\ell m} Y_{\ell m}(\hat{n}) \), and for a given definition of the global \((\theta, \phi)\) directions, the definition is unique. Considering the isotropy argument above, even a different choice of \((\theta, \phi)\) would only result in different \( a_{\ell m} \) coefficients, but not the variance of their distributions, \( C_\ell \).

A combination of \( Q \) and \( U \) maps is necessary to unambiguously define the magnitude \( (P = \sqrt{Q^2 + U^2}) \) and orientation \( (\psi = \frac{1}{2} tan^{-1} \frac{U}{Q}) \) of the net linear polarization field. Polarization is then described by the function \( P_{\pm}(\hat{n}) \equiv Q(\hat{n}) \pm iU(\hat{n}) \) on the sphere. Because the orientations of \( Q \) and \( U \) are uniquely defined relative a fixed direction in the \( \varphi \) plane (defined in Figure 1.1), a rotation about \( \hat{n} \) does not preserve the value of \( P(\hat{n}) \), but instead mixes the \( Q \) and \( U \) components by \( P'_{\pm}(\hat{n}) = e^{\pm 2i\alpha} P_{\pm}(\hat{n}) \). Rotations by \( \alpha = n\pi \) for integer \( n \) result in \( P'_\pm(\hat{n}) = P_\pm(\hat{n}) \), so \( P_\pm(\hat{n}) \) are spin 2 functions. The spin 0 spherical harmonics do not have an accounting for \( \alpha \), and thus \( P_\pm(\hat{n}) = \sum_{\ell} \sum_{m=-\ell}^{\ell} a_{\ell m} Y_{\ell m}(\hat{n}) \) is not uniquely defined, even for a fixed definition of \((\theta, \phi)\).

Stokes parameters are a useful basis for making measurements, but the physical importance of CMB polarization is not captured in \( P(\hat{n}) \) as linear combination of \( Q \) and \( U \) maps. Recall that the full-sky ensemble of polarization fluctuations generated by scalar perturbations has no curl component, while the tensor perturbations generate both curl and gradient patterns. The definitions of curl and gradient components in the polarization field are inherently non-local, unlike the \( Q, U \) basis that comprises the \( P_\pm(\hat{n}) \) field. By applying lowering and raising operators to the spin \( \pm 2 \)-weighted functions \( P'_\pm(\hat{n}) \), \( Q \) and \( U \) maps of polarization can be decomposed into orthogonal scalar field components, \( E \) and \( B \), that can then be unambiguously described by the standard spin 0 spherical harmonics \([43]\).

The physical significance of these scalar fields is crucial, as they describe the curl-
free, or E-mode, and gradient-free, or B-mode, components of $P_{±}(\hat{n})$. Recall that B-modes are only generated in the primordial CMB by tensor perturbations from the inflationary gravity wave (IGW) background in the primordial plasma. Thus, measuring primordial B-modes offers a unique opportunity to probe a signal whose origin lies in the earliest fractions of a second in the life of the universe.

![Figure 1.2: Power Spectra of the CMB](image)

Power spectra of the CMB from the present best-fit ΛCDM model [2], along with two value of the tensor-to-scalar ratio, $r$. E-mode polarization (black curve) is generated predominantly by scalar density fluctuations in the primordial CMB. Gravity wave B-modes (blue) are generated by the background of tensor fluctuations propagating through the primordial plasma. Lensing B-modes (green) are generated by deflections of the primary CMB photons by evolving structure after recombination.

### 1.3 Cosmology with B-Mode Polarization

The measured B-mode component of the CMB is expected to have two distinct sources: tensor perturbations from the inflationary gravity-wave background in the primordial CMB (inflationary B-modes), and gravitational lensing of primordial E-modes by evolving mass fields inducing a post-recombination signal (lensing B-modes). It is clear that an unambiguous measurement of primordial B-modes would be strong
evidence for inflation, as inflation generically predicts the tensor modes capable of producing the signal. Lensing B-modes are a new means by which the CMB can be used to measure structure evolution and probe cosmology between last scattering and today.

Using the decomposition of polarized CMB maps described in the previous section, we can make an estimate of the angular power spectrum, \( C^B B \), of the B-mode field. Using cosmological analysis codes like the ones discussed in [42], the spectrum of tensor perturbations responsible for the B-mode-generating quadrupoles at last scattering,

\[
P_t = P_t(k_0) \left( \frac{k}{k_0} \right)^{n_t},
\]

(1.3)
can be inferred. Similarly, the \( C^T T \) spectrum can be referred to post-inflationary scalar fluctuation spectrum,

\[
P_s = P_s(k_0) \left( \frac{k}{k_0} \right)^{n_s(k_0)-1},
\]

(1.4)
where, following the conventions in [28], departures from scale invariance are described by \( n_t \) (0 for a scale-invariant tensor spectrum), and \( n_s \) (1 for a scale-invariant density spectrum). The pivot factor, \( k_0 \) establishes the reference scale for the spectra, and is commonly chosen to be \( k_0 = 0.002 \text{Mpc}^{-1} \), which corresponds to super-horizon scales at last scattering. The ratio of these two spectra,

\[
r = \frac{P_t(k_0)}{P_s(k_0)},
\]

(1.5)
has crucial implications for the study of inflation, as different models predict different values of \( r \) ([3], [9], [36], [30], [4]).

In standard single-field slow-roll models of inflation, measuring \( r \) yields information about two interesting parameters: the energy scale of inflation and the width of
the inflationary potential. As described in, e.g. [35], $P_s \propto \frac{V^3}{V'}$, and $r \propto \left(\frac{V'}{V}\right)^2$, which leads to $r \propto P_s V^4$, where $V$ is the inflationary potential. Normalized values of $P_s$ are well-measured [2], leading to

$$V^{\frac{1}{2}} \approx 10^{16} GeV \left(\frac{r}{0.01}\right)^{\frac{1}{4}}, \quad (1.6)$$

which relates the inflationary gravity wave signal to an energy scale far beyond anything attainable by other experimental means. In standard slow-roll inflation, it can be shown that [35]:

$$\frac{\Delta \phi}{M_{pl}} \gtrsim \sqrt{\frac{r}{0.01}}, \quad (1.7)$$

Thus, large values of $r$ imply a super-Planckian width of the inflaton field excursion during inflation, which has considerable implications for the ultraviolet completion of quantum field theory and gravity [4].

At last scattering, the CMB polarization pattern is predominantly E-modes, with a small B-mode component that depends on the amplitude of tensor fluctuations from inflation. Following recombination, photons stream freely through the universe while matter density fluctuations evolve due to their gravitational instability. As they traverse the universe, photons are affected by the evolving potentials, such that the measured CMB is a slightly warped version of the source field. The details of the warping effect, called weak lensing, can be related to important questions of both cosmology and particle physics. In particular, a B-mode signal will be produced by weak lensing even if $r$ is negligibly small, and for any reasonable value of $r$, the lensing spectrum of B-modes will dominate over the primordial gravity wave spectrum on arcminute scales. Thus, to measure the primordial B-mode spectrum on angular scales where lensing B-modes dominate, the lensing spectrum can be treated as a noise bias and subtracted off. The ability to remove the lensing signal depends on the precision with which the B-mode patterns are measured.
Starting with the two-dimensional projection of the lensing potential, \( \phi(\hat{n}) \), and measured maps of the CMB temperature and polarization fields, \( T'(\hat{n}) \), \( P'(\hat{n}) \equiv [Q' + iU'](\hat{n}) \), the relationship to the unlensed source fields is

\[
T'(\hat{n}) = T(\hat{n} + \nabla \phi(\hat{n}))
\]
\[
P'(\hat{n}) = P(\hat{n} + \nabla \phi(\hat{n})).
\]

(1.8)

Note that the sky is described in terms of \( Q \) and \( U \), rather than \( E \) and \( B \), and the reason for this underpins the utility of B-modes in studying weak lensing. Recall that \( Q \) and \( U \) are local quantities on the sky, but vary under coordinate rotation. While \( E \) and \( B \) are spin 0 functions, they are non-locally defined on the sky by integrating linear combinations of \( Q \) and \( U \) about each point \( \hat{n} \) [52]. Thus, deflections of \( Q \) and \( U \) described in Equation 1.8 cause mixing between \( E \) and \( B \). A reasonable approximation of the CMB polarization at small angular scales (where the lensing deflection effect is strongest) is a field of pure E-modes. A measured CMB spectrum is then expected to have a B-mode component from \( E \to B \) mixing resulting from weak lensing deflections.

There are two primary results promised by precise measurements of \( \phi \). The most commonly cited one, a constraint on the sum of neutrino masses (\( \Sigma m_\nu \)), is of particular interest because it probes particles physics through cosmological observations with greater precision than is possible with traditional particle-physics experiments. Neutrino mass has a subtle effect on the primary CMB at large angular scales because of the smoothing effect of relativistic neutrino density on perturbations in the primordial plasma, allowing for a constraint of \( \Sigma m_\nu < 1.3eV \) from temperature power-spectrum analysis alone [14]. However, at smaller angular scales where weak lensing becomes important, the impact of \( \Sigma m_\nu \) is much stronger. Non-relativistic neutrinos do not cluster on small scales during the growth of structure after recombination, but do contribute to the total energy density, thus increasing the expansion rate relative to
a $\Sigma m_\nu = 0$ state, effectively slowing the collapse of massive structures on small scales [31]. Including a lensing analysis in concert with other observables, but without B-modes, results in a limit from the Planck experiment of $\Sigma m_\nu < 0.23$ eV. Lensing B-modes are a clear, unambiguous result of the lensing field, suggesting potential future limits of $\Sigma m_\nu < 0.05$ eV [44].
2.1 Observing the CMB at the South Pole

2.1.1 Measuring the CMB

The CMB intensity spectrum peaks at $\sim 160$ GHz, where it is the brightest astrophysical source in the sky. As shown in Figure 2.1, fluctuations in synchrotron emission surpass the CMB anisotropies below about 20 GHz, while galactic dust emission dominates above 200 GHz. Foreground intensity varies considerably on the sky and is generally concentrated along the galactic plane, allowing for lower levels of contamination when observing smaller patches of sky at high galactic latitude. The principal fields observed by the South Pole Telescope Polarization experiment (SPTpol) lie within a low foreground region visible at moderate elevations from the South Pole, as shown in Figure 4.1. The first season deep field is centered at right ascension (RA) of 23h30m and declination (DEC) of $-55^\circ$, and is 15 degrees wide in right ascension (RA) and 15 degrees “tall” in declination (Dec). The survey field, which will be observed for the remaining three years of SPTpol, is centered at RA
0h0m, DEC -57.5 and is 15° high in Dec by 60° wide in RA.

Figure 2.1: RMS Fluctuations of Celestial Microwave Sources vs. Photon Frequency
RMS fluctuations of galactic foregrounds and the CMB as a function of photon frequency, near the SPTpol observing bands, in antenna temperature units. The foreground spectra shown are derived from WMAP 9-year data [5]. Regions near the galactic plane, where foreground emission completely dominates the CMB brightness, are masked before calculating the spectra shown above.

For a ground-based telescope observing mm-wave signals, the emission spectrum of atmospheric oxygen and water vapor, shown in Figure 2.2, has a series of strong lines that completely dominate any celestial signal, between which are windows of lower emission suitable for observing the celestial microwave sky. Near the peak CMB brightness, molecular oxygen has emission lines centered at ∼60 and 120 GHz and water emits at ∼180 GHz, defining the edges of ground-based observing bands centered at 90 and 150 GHz. Continuum emission between the lines is a strong source of mm-wave photons, which, combined with instrumental emission, sets the instantaneous instrumental sensitivity limit. Thus sites with thin, dry atmospheric conditions desirable for CMB observations. Oxygen and water are well-mixed in the atmosphere, so its effect on detector loading is a function of integrated line-of-sight airmass, which varies smoothly with elevation. In addition its elevation-dependent loading compo-
ent, water-vapor content varies considerably on both spatial and temporal scales relevant to typical observing strategies, resulting in low-frequency noise in detector data from atmospheric opacity variations due to water vapor.

![Sky Brightness Spectra](image)

**Figure 2.2: SPTpol observing bands**

The nominal 90 GHz and 150 GHz observing bands for SPTpol, plotted with the CMB intensity spectrum $B_\nu(T)$ and the effective atmospheric brightness at the South Pole for 0.26 mm precipitable water vapor [19]. Wider observing bandwidth at low background loading improves sensitivity, but atmospheric $O_2$ and $H_2O$ emission lines obscure celestial signal in addition to increasing detector loading and noise. Between these emission lines, the low continuum emission of the thin, dry, and stable atmosphere at the South Pole makes it one of the best sites in the world for CMB observations.

### 2.1.2 Observing from the South Pole

At just over 2800m elevation, the Amundsen-Scott South Pole Station sits below a thin atmosphere that is also among the driest on Earth, with a median precipitable water vapor during the winter of $\sim 0.25$ mm [7], [29]. The weather is typically very stable, with very little cloud cover and predictable, low-velocity winds generally flowing from the same direction off the Antarctic plateau. The thin atmosphere reduces the absolute loading on the instrument’s detectors, while its stability reduces the am-
plitude of fluctuations in the data from the variable opacity of the moving airmass. Further improving the conditions is the fact that the Sun is below the horizon for a full six months of every year, from mid-March through mid-September, allowing for uninterrupted observation of target fields.

The United States has maintained a South Pole base since 1957, housing astronomy and astrophysics experiments since the 1980s. The station offers infrastructure and logistical support, including a power plant and computing and internet services that are crucial to managing telescope data. In a typical season, the summer months are spent observing fields that are least contaminated by the Sun, along with shutting down the telescope for maintenance and upgrades. When the station closes in mid-February, a winterover crew of around 50 people, including two dedicated SPT personnel, maintain the station and keep the telescope operating.

2.2 The South Pole Telescope

2.2.1 The Telescope

The South Pole Telescope is a 10 m primary, off-axis Gregorian telescope located at the Dark Sector Laboratory of the Amundsen-Scott South Pole Station. The primary mirror consists of 218 individually-machined aluminum reflecting panels, mounted and aligned with a final surface RMS of 25 mm. Strips of BeCu with spring fingers to hold them in place are inserted between panels to minimize scattering from the \( \sim 2 \) mm gaps. The primary is supported by an L-shaped frame affixed to an altazimuthal mount at its vertex, with the secondary optics cryostat and the receiver cryostat containing the SPTpol camera mounted in the other arm of the frame. With this design, the entire system moves during scanning, allowing the beam to be swept across the field without varying the illumination of the primary. The drive allows for scans in excess of 2 degrees per second in telescope azimuth, so that sky signals
The South Pole Telescope. The guard ring, which prevents diffraction of low-elevation radiation around the primary, was installed prior to first light in early 2012. The comoving ground shield was installed prior to the second season of observations, beginning in early 2013. The guard ring, ground shield, and snout (shown in Figure 2.4) suppress polarized sidelobe pickup by preventing stray light from entering the optical path.

2.2.2 Optics

Radiation from the sky reflects off the primary and enters the optics cryostat through a Zotefoam window located near prime focus, then passes through two sets of IR-blocking filters that are cooled 100 K and 10 K and serve to reduce out-of-band loading on the cold optical system. Blackened optical baffles, which are also cooled to 10K, surround the beam as it approaches and reflects off the 1 meter diameter secondary mirror, terminating spillover from the secondary on a stable load and serving as the cold stop of the system. The beam then enters the receiver cryostat, which shares a vacuum with the optics cryostat, through a second set of IR blocking filters cooled to 10K and 4K, before being shaped by a weak high-density polyethylene (HDPE) lens to make the field telecentric as it couples to the focal plane. Finally, before reaching the
Figure 2.4: The SPTpol Optics and Receiver Cryostats
A cutaway view of the joined optics and receiver cryostats, which share a vacuum but are cooled by separate pulse tube coolers. The inner surface of baffles surrounding the beam inside the optics cryostat is Emerson & Cumming HR-10, a microwave-absorbing foam that helps ensure stray rays terminate on a stable, cold surface. Outside the environmental window, a warm absorbing “snout” prevents rays from wide angles scattering off the window and into the optical path.

Pixel-horn apertures, radiation passes through two “band defining” filters: a low-pass filter that determines the upper edge of the detector sensitivity band and a harmonic blocker, which is a second low-pass filter with a higher band edge but designed to block harmonics in the primary edge filter. The combination of the two filters ensures a sharp, well-controlled cut-off in the frequency response of the detectors, as well as mitigating high-frequency leakage from harmonics of the primary filter transmission band. Figure 2.4 shows a cutaway view of the optics and receiver cryostats in their mated configuration, with traces showing the paths of prime and marginal rays from detectors at the center and edges of the focal plane.

For the first season of SPTpol observations in 2012, a guard ring around the primary mirror was installed to mitigate diffraction around the primary and into the optical path. In addition, an absorbing cone called the “snout” was added near
prime focus outside the secondary cryostat to further suppress stray light coupling. After the first season of observations, co-moving ground shields were added along the telescope frame and mated to guard ring. The result, illustrated in Figure 2.3, is a three-component co-moving shield system that prevents stray light from coupling into the optical path and detector far sidelobes.

Figure 2.5: Effect of Ground Shields on Far Sidelobes
Maps made with a Gunn oscillator source mounted on a nearby building, emitting horizontally-polarized radiation. In 2012, only the guard ring around the primary and the snout were installed. In 2013, the full comoving ground shields were added, significantly reducing the effect of scattering into the optical path.

2.2.3 Cryogenics

The baffles in the optics cryostat are nested set of conical sections tracing the beam profile as it reflects off the secondary mirror. The outer set of baffles are cooled with a PT410 pulse tube cooler (PTC) to a temperature of 100K, and they serve as a
heat break between the ambient cryostat wall and the inner baffles. Nested inside the outer cones, the optical baffles and secondary mirror are cooled to 10K with the same PTC used for the outer baffles. Infrared blocking low-pass filters mounted on the skyward side of the optical baffles are affixed with aluminum plates that ensure a strong thermal link to the cooling power of the PTC. The receiver cryostat has a set of short 70K and 10K baffles that interleave with the ones in the optics cryostat when the two are joined. The two sets of cones are prevented from touching to maintain the independence of their cooling systems. Layers of super-insulation extend into the space between the receiver and optics baffles, blocking stray light without providing a significant thermal contact between them.

The receiver cryostat is cooled with a PT415 PTC that provides an intermediate temperature of 50 K and a base temperature of 4 K. A heat break jacket cooled by the 50 K pulse tube head surrounds the 4K stage and includes the outermost optical baffle. A pair of low-pass IR blocking filters are mounted on either side of the HDPE lens, and all three are cooled to around 6 K by the 4K head of the PTC. The temperatures of various elements of the SPTpol optical system are listed in Table 2.1.

The focal plane array structure is affixed to a 4 K main stage, as is the Simon Chase “Helium 10” evaporation refrigerator that provides the sub-Kelvin cooling [21]. The focal plane array itself is supported by a truss structure, comprised of the focal plane nested inside an intermediate cooling stage that then connects mechanically to the 4K base plate. Hollow, cylindrical supports made of Du Pont Vespel SP-22 polyimide resin join the various stages. Vespel has an extremely low thermal conductivity at sub-Kelvin temperatures, while providing the necessary strength to support the mass of the focal plane.
<table>
<thead>
<tr>
<th>Secondary Cryostat (PT410)</th>
<th>Temp (K)</th>
</tr>
</thead>
<tbody>
<tr>
<td>IR Shaders (3)</td>
<td>100</td>
</tr>
<tr>
<td>12 cm$^{-1}$ IR blocker</td>
<td>100</td>
</tr>
<tr>
<td>Outer baffle shell</td>
<td>100</td>
</tr>
<tr>
<td>10 cm$^{-1}$ IR blocker</td>
<td>10</td>
</tr>
<tr>
<td>Inner optical baffle</td>
<td>10</td>
</tr>
<tr>
<td>Secondary mirror</td>
<td>10</td>
</tr>
<tr>
<td>Receiver Cryostat (PT415)</td>
<td></td>
</tr>
<tr>
<td>Outer baffle shell</td>
<td>70</td>
</tr>
<tr>
<td>Outer thermal jacket</td>
<td>50</td>
</tr>
<tr>
<td>Inner optical baffle</td>
<td>10</td>
</tr>
<tr>
<td>12 cm$^{-1}$ IR blocker</td>
<td>4</td>
</tr>
<tr>
<td>HDPE lens</td>
<td>4</td>
</tr>
<tr>
<td>7 cm$^{-1}$ IR blocker(current)</td>
<td>4</td>
</tr>
<tr>
<td>Focal plane (Simon Chase He10)</td>
<td></td>
</tr>
<tr>
<td>Intermediate buffering stage</td>
<td>0.4</td>
</tr>
<tr>
<td>7 cm$^{-1}$ IR blocker(2012 season)</td>
<td>0.4</td>
</tr>
<tr>
<td>Harmonic blocking filters (2)</td>
<td>0.27</td>
</tr>
<tr>
<td>Band-defining filters (2)</td>
<td>0.27</td>
</tr>
<tr>
<td>Ultra-cooled stage</td>
<td>0.27</td>
</tr>
<tr>
<td>Detector array</td>
<td>0.27</td>
</tr>
</tbody>
</table>

Table 2.1: SPT Cryogenic Stages

Cryogenic surfaces in the SPTpol instrument. Each pulse tube cooler (PTC) has a warm and cold stage to provide two different base temperatures for elements attached to it. The IR blockers are polyamide films with an embedded metal mesh layers that act as low pass filters for incoming radiation, reflecting out-of-band power to mitigate optical loading on colder surfaces. IR shaders are a single metalized layer patterned onto a thin film to prevent thermal radiation from warming the IR blocking filters and the colder telescope stages. The focal plane cooler, mounted on a base place cooled by the 4 K stage of the receiver cryostat PTC, is a Helium-10 evaporation refrigerator [21]. The band-defining and harmonic blocking metal-mesh filters have upper transmission edges of 3.7 (5.6) and 5.6 (7.0), respectively at 90 (150) GHz. Each set of band-defining filters covers only the feedhorns for detectors in the appropriate band.

### 2.3 The SPTpol Camera

The SPTpol focal plane consists of 588 polarization-sensitive pixels at 150 GHz, arrayed in seven 84-element monolithic arrays, which are surrounded by 90 GHz pixels in 180 individual modules, grouped onto four mounting boards. Each pixel contains two polarization sensitive bolometers (PSBs), labeled “X” and “Y”, each sensing a single linear polarization state, and nominally aligned at 90° to each other. Thus, each pixel can measure a single Stokes polarization component, and two pixels
rotated 45° to each other can fully describe the polarization state at a point on the sky.

The assembled SPTpol focal plane, as seen from the sky side, without any of the IR-blocking or band-defining filters. The seven 150 GHz monolithic arrays are in the center, with the 90 GHz modules surrounding them on four mounting boards covering a quarter of the full annulus each. Each 150 GHz module has 84 pixels, containing a superconducting Nb orthomode transducer (OMT) antenna that couples orthogonal linear polarizations of light to separate bolometers. The 180 90 GHz pixel modules contain two separate detectors with PdAu resonators that dissipate photon power directly onto the detector. Each detector is sensitive to a single linear polarization, and the two detectors are mounted in a pixel at 90° to each other.

The definition of the Stokes vectors on the sky illustrated in Figure 1.1 translates to an axis on the installed focal plane that corresponds to the +Q direction. Relative to this direction, defined as the 0° angle, 90 GHz detector angles are distributed between 0° and 157.5°, in multiples of 22.5°, with the two detectors in a single pixel separated by 90°. The 150 GHz detector angles are distributed between 8° and 173°,
in multiples of $15^\circ$. Each of the four 90 GHz boards has detectors at all of the possible orientations. The 150 GHz arrays are fabricated with all pixels having one of two orientations, with $45^\circ$ relative angle between them. The spread in angles among 150 GHz detectors comes from the global rotations of the pixel arrays themselves. In the end, each detector response can be decomposed into a linear combination of its response to a fixed I, Q, U basis on the sky, as described in Section 4.3.

### 2.3.1 90 GHz Modules

<table>
<thead>
<tr>
<th>Module Structures</th>
<th>Detector Structures</th>
</tr>
</thead>
<tbody>
<tr>
<td>Module diameter</td>
<td>10.62 mm</td>
</tr>
<tr>
<td>Horn aperture diameter</td>
<td>10 mm</td>
</tr>
<tr>
<td>Horn length</td>
<td>35 mm</td>
</tr>
<tr>
<td>Waveguide diameter</td>
<td>2.35 mm</td>
</tr>
<tr>
<td>Radiation choke outer diameter</td>
<td>3.19 $\mu$m</td>
</tr>
<tr>
<td>Choke face separation</td>
<td>150 $\mu$m</td>
</tr>
<tr>
<td>SIN film thickness</td>
<td>1 $\mu$m</td>
</tr>
<tr>
<td>Detector thermal island</td>
<td>3300 $\mu$m x 200 $\mu$m</td>
</tr>
<tr>
<td>Thermal standoff legs</td>
<td>1350 $\mu$m x 10 $\mu$m</td>
</tr>
<tr>
<td>PdAu bar absorber</td>
<td>18 $\mu$m x 1300 $\mu$m</td>
</tr>
<tr>
<td>TES length x width</td>
<td>70 $\mu$m x 40 $\mu$m</td>
</tr>
</tbody>
</table>

**Table 2.2: 90 GHz Module Geometries**

*Left*: Geometries of the 90 GHz module assembly, shown in the left panel of Figure 2.7. *Right*: Geometries of individual detector structures. The thermal island spans the waveguide and extends into the choke structure, where the TES is mounted, to minimize direct coupling of radiation to the TES.

Each 90 GHz pixel is coupled to the sky through a smooth-walled, profiled horn feeding a circular waveguide that terminates in a backshort $\frac{\lambda}{4}$ (at the band center) behind the detector absorbing structures. The profiled horn produces symmetric, nearly Gaussian beams for both co-polar and cross-polar radiation in much the same way as the more traditional corrugated feeds with similar performance characteristics, but is easier and less expensive to manufacture [53]. The waveguide and feed horn are machined from a single piece of aluminum using a custom reamer that cuts out the precise horn profile. In a fully-assembled module, the horn is screwed to a mounting base that holds both the detectors and the hardware for mounting the module on the focal plane.

The module base has a $\frac{\lambda}{4}$ backshort machined into it, and also has hardware for mechanical and electrical connections to its parent circuit board. Four circuit boards
affixed to the 275 mK base plate around the 150 GHz arrays hold 45 pixel modules each and route their electrical connections to readout boards located underneath the base plate. Modules are aligned on the circuit boards with two 0.9 mm-diameter dowel pins and held in place with a single 00-80 screw. Individual detector wafers are glued to the mounting base and wirebonded to electrical pins that penetrate the module base, surrounded by an insulating sheath, and pass through the circuit board where they are soldered into place.

Individual detectors consist of a SiN membrane initially deposited on a Si wafer, from which the detector frame is etched and the TES structures deposited with standard lithography techniques. The Si base is then removed from under the sensor island to leave the detector structures suspended on the SiN membrane. The remaining Si mounting base provides mechanical support, alignment geometry, and a thermal sink. The detector itself consists of four principle components, described in more thorough detail in Section 3.1:

1. Detector membrane
   A rectangular membrane of SiN that supports the other structures and transmits power from the absorber to the sensor.

2. Absorber
   A thin bar of resistive PdAu that couples to a single polarization state of the TE01 mode in the waveguide, dissipating photon power onto the detector membrane.

3. Sensor
   A MoAu bilayer transition edge sensor (TES) that measures fluctuations in the temperature of the detector membrane resulting from changes in absorbed photon power.

4. Thermal link
Six long, thin SiN legs that restrict heat flow from the detector membrane to the thermal sink, controlling the thermal response of the detector. Superconducting electrical leads to the TES sensor run along two of the legs, contributing negligible additional thermal conductivity as they contribute only 3% of the total volume of the leg structure [51].

The various dimensions of the module and detector wafer structures are described in Table 2.2.

![Figure 2.7: 90 GHz Module Details](image)

*Left:* Cutaway-view drawing of an assembled 90 GHz module. The choke prevents radiation from leaking into the break in the waveguide where the detector wafers extend into it. Electrical connections to the detector wafers (not shown) run through the mounting base and to the other side of the focal plane mounting plate. *Right:* View of a mounted 90 GHz pixel, showing two crossed detectors from the sky side. The inset shows a closer view of the TES structure on the edge of the SiN detector island, as well the legs that provide the weak thermal link between the detector island and the Si mounting base, seen here as the black square structures. The PdAu bling and Nb dot structures are described in Section 3.1.4.
2.3.2 150 GHz Arrays

In contrast to the 90 GHz pixel design that separates individual pixels into discrete assemblies, the 150 GHz detectors are fabricated on wafers containing 84 pixels each. Radiation is coupled to the pixels through monolithic arrays of corrugated, profiled horns that feed square waveguides that transition to a short circular waveguide section before coupling to the pixels. Each horn array is comprised of 33, 500 µm-thick silicon platelets, individually etched and then aligned, stacked, bonded, and finally gold plated. As with the 90 GHz profiled feeds, the corrugated design provides the necessary systematic control over the polarized beams [6].

The design of SPTpol 150 GHz detector arrays is described in [18], but I will provide an overview here for completeness. From the sky side, each 150 GHz module consists of the following elements:

1. **Feedhorn module**
   
   A gold-plated silicon platelet array that provides a corrugated feedhorn for each pixel.

2. **Waveguide interface plate**
   
   A silicon wafer that extends the waveguide from the platelet array to within 25 µm of the orthomode transducer (OMT) membrane, minimizing leakage of photons into the waveguide break.

3. **Pixel wafer**
   
   A silicon base structure onto which a SiN membrane is deposited. Both the OMT and TES structures are then deposited onto the SiN layer and the Si behind them is removed. Microstrip transmission lines and electrical traces run along the SiN and to interface pads on the edge of the wafer.

4. **Backshort wafer**
   
   Contains the $\frac{\lambda}{4}$ waveguide backshorts, as well as surrounding “moats” filled with
ECCOSORB [15] to mitigate radiation leakage into the cavity between the pixel wafer and backshort wafer.

In addition to the 84 optical pixels on each 150 GHz wafer, there are two resistor channels used for making Johnson noise measurements and two heater resistor channels for warming the wafer. More importantly, there are five dark detectors, which are TES sensors that are not coupled to a feedhorn or an OMT. These dark detectors are extremely useful for two main purposes: One is to check the inherent coupling of the TES structure itself to stray radiation that has leaked out of the waveguide structures and into the wafer cavity. The other is to help investigate unexpected artifacts in the data, particularly whether they originate from optical signals (and would thus not be present in the dark detectors), or from some other element of the focal plane environment (such as temperature, magnetic fields, or readout electronics).

### 2.4 Readout

#### 2.4.1 Frequency Multiplexing

Readout of the TES detectors in SPTpol is performed with a digital frequency multiplexing (DfMux) system that is an upgrade to the analog fmx system employed by SPT-SZ and described in detail in [12]. The DfMux system is comprised of the same basic components as the analog system, and their operation is conceptually and practically very similar. In this section, I will provide a brief overview of frequency multiplexing, followed by a discussion of the details of the SPTpol DfMux implementation. A thorough discussion of TES detectors follows in 3.1, but for the purposes of explaining the readout, it is sufficient to think of a TES as variable resistor whose fluctuations we wish to measure.

In the frequency-domain multiplexing system, a series of bias voltage signals, \( \sum_i V_{B,i} e^{j\omega_i t} \), are sent down a single pair of wires to a chain of TES detector channels.
Figure 2.8: The Digital Frequency-Multiplexing Circuit
The basic frequency multiplexing (fMux) circuit for a single SQUID amplifier. The DfMux motherboard generates a “comb” of bias signals and associated nuller signals, the former of which is routed to a series of combination of LRC detector circuits. Bias resistors $R_{B,C}$ and $R_{B,N}$ convert the combs from voltage to current sources. The bias current then flows across the parallel combination of the LRC circuits and a small shunt resistor, $R_{SH}$, which draws most of the current, setting up a stiff voltage bias across the parallel LRC array. Fluctuations in the TES resistances ($R_{TES}$) modulate the current flowing through each LRC leg, and the nuller current cancels an unmodulated version of the original bias, leaving the residual current modulations, which have a much lower amplitude than the bias component, to be read out by the SQUID transimpedance amplifier. The SQUID is placed in a negative-feedback flux-locked loop (FLL) that further nulls the input current to $L_{SQ}$, while converting the current fluctuations into an amplified voltage signal that is then routed back to the DfMux where it is digitized, demodulated, and downsampled. Section 2.4.2 describes the physical components, labeled here in grey, in greater detail.

wired in parallel, as illustrated in Figure 2.8. The combination of bias voltages is referred to as a “comb”, in reference to the peaked structure of its PSD. Each detector channel consists of an LRC series circuit, where the TES acts as a variable resistor with $\langle R_{TES} \rangle \approx 0.85\Omega$, $L$ is constant $24 \mu H$ for all channels, and each channel has a different capacitance, $C_i$, that puts the LC resonance, $\omega_i = \sqrt{1/ LC_i}$, between 300-1200 kHz.
The total LRC impedance, for channel $i$ is:

$$Z_i(\omega) = R_{TES,i} + i\omega L - \frac{i}{\omega C_i},$$  \hspace{1cm} (2.1)

which reduces to $Z_i = R_{TES,i}$ when $\omega = \omega_i$. In order to provide a fixed voltage bias across each channel, the LRC channel array is placed in parallel with a 30 mΩ shunt resistor, $R_{shunt}$. Bias signals are initially generated on the DfMux board as voltages, $V(\omega_i)$, then turned into currents by a bias resistor $R_B \gg R_{shunt} \parallel \Sigma LRC$. The incoming stiff bias current flows mostly through $R_{shunt}$, so that the voltage across channel $i$ is fixed at $V_{b,i} \simeq I_{bias,i}R_{shunt}$, and $V_{b,i}$ across the TES is related to $V_{B,i}$ from the DfMux output by a scalar factor.

With a fixed bias voltage set across each channel, modulations in $R_{TES,i}$ induce fluctuations in the resonant current, $I_i(\omega_i)$. Modulated current signals are recombined after passing through their associated detector channels. The magnitude of current fluctuations induced by sky signals is much smaller than the $V_{b,i}/\langle R_{TES} \rangle$ bias amplitudes, so after passing through the detector channels, a scaled version of the input bias waveform is subtracted from the combined current output. Ideally, this nuller signal completely removes the input bias current component and leaves only sideband amplitude modulations. In practice, changes in $\langle R_{TES} \rangle$ result in some fraction of the bias signal remaining in the nulled current. Still, nulling vastly reduces the dynamic range of the readout current signal, allowing for greater amplification.

We measure the minute fluctuations in TES current with series-array SQUID transimpedance amplifiers manufactured at NIST in Boulder, CO. Nulled current signals pass through a 100 nH inductor, $L_{SQ}$ in Figure 2.8, such that the resulting magnetic flux $\phi$ is coupled to the SQUID. A DC current bias applied to the SQUID sets up a DC voltage level that changes as a function of $\phi$, as shown in Figure 2.9. We choose a DC bias current that maximizes the modulation depth of the $V - \phi$ curve,
and inject a DC current into the detector readout path to servo the SQUID voltage to a point where $\frac{\partial V}{\partial \phi}$ is approximately linear. Without further treatment, detector current is simply transformed into an output voltage according to the $V-\phi$ curve, so only the smallest changes in $\phi$ result in a linearly-scaled output voltage.

Figure 2.9: SQUID V-Phi Curve

V-Phi curve of a SQUID. Voltage across the SQUID array, $V_{SQ}$ or ”SQUID Bias” is induced by an applied DC bias current. Increasing the bias raises the DC level of the SQUID response curve and changes its modulation depth. Increasing the current through an inductor, $L_{SQ}$, coupled to the SQUID increases magnetic flux, $\phi$, through the SQUID itself, inducing the sinusoidal response shown. The period of the SQUID output voltage, labeled here as ”Voltage Offset” vs $\phi$, or simply $V-\phi$, curve corresponds to a change in flux of a single flux quantum, $\phi_0 = \frac{h}{2e}$, which is induced by a change in the current through $L_{SQ}$ of $\sim 25 \text{ } \mu$A for the SPTpol SQUIDS.

Linearity in the SQUID transimpedance is improved by placing it in a negative-feedback op-amp flux locked loop (FLL), where the output voltage drives a flux bias current through $R_{FB}$ in Figure 2.8 that servos the input to $L_{SQ}$ towards zero. This FLL operation limits the bandwidth of the squid response due to to a phase shift from the impedance of wires running from the output of the op-amp (at 300 K) back to the flux bias injection point near the SQUID (at 4K). As a result of
the bandwidth limit, the highest stable bias frequency for multiplexed channels using SPTpol wire lengths is approximately 1.2 MHz. The lowest possible frequency is set by AC-coupled components in the readout electronics at approximately 100 kHz.

As an alternative to linearizing SQUID output with the analog FLL, the DfMux electronics allow for a precise real-time digital active nulling (DAN) of the summed signals before they pass through the coupling inductor [10]. In DAN operation, the nuller amplitude and phase for each detector channel are updated dynamically, with an effective bandwidth of \( \sim 10 \) kHz around each bias frequency. This essentially cancels both the bias signal and the modulations within the signal bandwidth of \( \sim 100 \) Hz, such that the nuller signal can be read out as the TES current output. Instead of serving as the readout for TES current fluctuations, the SQUID effectively becomes a monitor for the DAN integrator, obviating the need for the FLL to dynamically inject a flux bias to cancel SQUID input current. DAN suppresses the amplitudes of signals near the bias frequencies, while the LRC circuits suppress signals far from any bias frequency, so in nominal operation the SQUID input inductor is only ever coupling small-amplitude currents.

An important consideration in frequency multiplexing is the response of a detector channel to off-resonant signals from other channels on the comb. The impedance of detector \( i \) seeing a signal from at the bias frequency of a different detector, \( j \), is:

\[
Z_i(\omega_j) = R_{TES,i} + i\omega_j L - \frac{i}{\omega_j C_i}. \tag{2.2}
\]

The total current in channel \( i \) from its own bias signal and the one from channel \( j \) is then:

\[
I_i(\omega_i) = \frac{V_i e^{i\omega_i t}}{Z_i(\omega_i)} + \frac{V_j e^{i\omega_j t}}{Z_i(\omega_j)}, \tag{2.3}
\]

and the magnitude ratio of the leaked current to the on-resonance current, assuming
$V_i = V_j$ is given by:

$$
\frac{|I_i^{\omega_j}|}{I_i^{\omega_i}} = \frac{R_{TES,i}}{\sqrt{R_{TES,i}^2 + \left(\omega_j L - \frac{1}{\omega_j C_i}\right)^2}}.
$$

(2.4)

In the SPTpol DfMux system, $R_{TES} \sim 0.85\Omega$, $L = 24\mu H$, and the bias frequencies range between 300-1200 kHz, requiring $0.7nF < C_i < 12nF$. We can make the approximation

$$
\sqrt{R_{TES,i}^2 + \left(\omega_j L - \frac{1}{\omega_j C_i}\right)^2} \approx 2\Delta \omega L,
$$

where $\Delta \omega = |\omega_j - \omega_i|$ is the spacing between adjacent channels in the bias comb [12]. Current leakage from $j$ acts as a current bias on channel $i$, which can lead to instability if the leakage component is too large, as the assumption of a constant, strong voltage bias is violated. For the nominal SPTpol channel spacing of $\sim 60$ kHz as shown in Table 2.3, the percentage of electrical bias power on channel $i$ from a next-closest-neighbor, $j$, with a similar resistance and bias voltage amplitude, is less than 1%. Combining the contributions of all other similar detectors on a bias comb, the total off-resonance leakage power is less than 3%.

Along with inducing bias carrier leakage, the presence of multiple bias signals on the same comb induces low levels of crosstalk between the output signals of each detector. Two inductors physically next to each other on the board containing the LC chips have a mutual inductance that can lead to signal crosstalk. However, in SPTpol, we arrange the inductors such that if they are in close physical proximity on the board, they are associated with resonant circuits at least $2\Delta \omega$ apart, so the crosstalk is heavily suppressed by the different LRC impedances.

Leakage current from channel $i \rightarrow j$ is modulated by changes in $R_{TES,j}$, so that the final, demodulated signal at bias frequency $\omega_i$ has a crosstalk component from the fluctuations in the resistance of detector $j$. The ratio of the magnitudes of the current fluctuations at $\omega_i$ from modulations in $R_{TES,i}$ to those from modulations in
\( R_{TES,j} \) can be approximated as:

\[
\frac{|\delta I_j|}{|\delta I_i|} \approx \frac{R_{TES,i}^2}{(2\Delta \omega L)^2},
\]  

(2.5)

again assuming the resistances and bias voltages are roughly equal, as they are
de- signed to be in the DfMux system [12]. The fraction of leakage signal from modula-
tions in nearby detectors for the SPTpol design \( R_{TES} \approx 0.8 \Omega, \Delta \omega \sim 2\pi \times 60kHz \). Stray impedances in the fMux circuitry can also induce crosstalk signals, but com-
ponents are chosen to minimize such strays, rendering the resulting crosstalk component
negligible.

2.4.2 SPTpol Readout Architecture

Here I will give a brief overview of the connections between the different stages in
the SPTpol DfMux readout chain, illustrated schematically in Figure 2.8 for a single
SQUID amplifier.

![Figure 2.10: LC Board Circuit Schematic](image)

A circuit schematic of one side of an LC board, prior to the addition of the circuit
components. On the left, a stripline that terminates in a 25-pin Micro D connector
routes signals to and from the SQUID card. Two pins for each SQUID are broken out
to the 12 LRC combinations fed by a single comb. The traces going to the inductors,
fabricated as 8-element arrays, are staggered so that nearest neighbors in frequency
are on separate inductor arrays. Channels using adjacent inductors on an array are
spaced by at least two spacings, \( \Delta f_{bias} \sim 60 \text{ kHz} \).
• **Focal Plane** From the 1599 individual readout elements on the SPTpol focal plane, superconducting electrical leads run along the 90 GHz module mounting PCBs and 150 GHz Si pixel wafers to wirebonding junctions on the periphery of each. Wirebonds then transfer the circuits to superconducting striplines that terminate in zero insertion force (ZIF) connectors that are plugged into one of 18 LC boards mounted behind the focal plane.

• **LC Boards** Mounted behind the focal plane and connected to the detector arrays by tinned-copper ZIF stripines, each two-sided LC board contains the series L, C combinations for 8 fMux combs with 12 channels per comb. The two-sided LC boards are designed such that for each comb, the LC components for channels nearby in frequency are spaced far apart physically. The two combs generated by each DfMux mezzanine have different frequency schedules, listed in Table 2.3, to minimize crosstalk in warm wiring between the DfMux Mezzanine and SQUID controller boards. A single LC Board has two offset circuit blocks on each of its sides, as shown in Figure 2.10, with each block containing the L,C components of the two frequency schedules for a single mezzanine’s two combs. Each LC board corresponds to 8 frequency combs and reads out 96 detectors.

• **SQUID Cards** Eight SQUID amplifiers are mounted on each 4 K SQUID card. A single pair of wires from each LC board contains a comb of modulated detector signal current, which is nulled at 4K on the SQUID card before coupling into \( L_{SQ} \) for readout. Without the DAN system running, the FLL feedback current comes from the warm SQUID controller board, necessitating short wire lengths to minimize phase shifts in the feedback current at high frequencies.

• **SQUID Controller Boards** A single SQUID controller board is used to bias and read out each SQUID card. Signals from a total of 4 mezzanine boards are routed through each SQUID controller and into the receiver cryostat. Amplified
SQUID voltages are fed back from the SQUID controller to the DfMux.

- **Mekzanine Boards** The analog waveforms for both the bias carrier and bias nulling combs are generated on the mezzanine board and sent to the SQUID controller, where they are converted from a voltages to currents. The bias current is then sent through to the parallel combination of $R_{SH}$ and the LRC array.

- **DfMux Motherboard** A single DfMux motherboard holds two mezzanine boards, and performs the bias generation and signal demodulation for 4 SQUIDs’ worth of detectors. Each motherboard has a Xilinx Virtex-4 field programmable gate array (FPGA) clocked at 25 MHz, running a Linux operating system. Embedded code is a combination of C and Python, and both the data and user interface are over ethernet.

The entire readout architecture is encoded in a “hardware map”, which consists of a series of XML files describing the parent/child relationships of each component in the readout chain. For a given comb, the hardware map lists the resonant frequencies for each bolometer, the individual SQUID used to read out the comb, and the DfMux board that performs the oscillation/demodulation.
<table>
<thead>
<tr>
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<th>$f_{bias}$ (kHz)</th>
<th>$C_i$ (nF)</th>
<th>$\Delta f$ (kHz)</th>
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<td>1</td>
<td>297.1</td>
<td>14.35</td>
<td>-</td>
</tr>
<tr>
<td>2</td>
<td>363.8</td>
<td>9.57</td>
<td>66.8</td>
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<td>6.88</td>
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<td>3.47</td>
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<td>2.92</td>
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<td>714.1</td>
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<td>1.81</td>
<td>57.2</td>
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<td>11</td>
<td>905.2</td>
<td>1.55</td>
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</tr>
<tr>
<td>12</td>
<td>958.0</td>
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<table>
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<th>$C_i$ (nF)</th>
<th>$\Delta f$ (kHz)</th>
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<tr>
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<td>386.8</td>
<td>8.46</td>
<td>57.2</td>
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<td>6.26</td>
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<tr>
<td>12</td>
<td>1006.1</td>
<td>1.25</td>
<td>69.7</td>
</tr>
</tbody>
</table>

Table 2.3: DfMux Frequency Schedule

Schedule of bias frequencies for comb associated with a single squid. Two different schedules are used to minimize crosstalk between the two combs generated on each DfMux mezzanine board. The capacitance values are achieved in most cases by soldering two or three separate capacitors in parallel. The measured cryogenic inductance is $L = 20\mu\text{H}$, rather than the warm specification of $24\mu\text{H}$. 
Chapter 3

DEVELOPMENT, DEPLOYMENT, AND CHARACTERIZATION OF SPTPOL

3.1 Transition Edge Sensor Bolometers

In this chapter, I will present an overview of transition edge sensor (TES) bolometer detectors, following the thorough treatment presented in [24]. The purpose is to describe the various components of a TES detector, illustrate their relationships to one another in the context of their use in the SPTpol pixels, and establish terminology that will be used in Section 3.2. Therefore, certain simplifications will be introduced to the general derivations in order to arrive at the notation used during the SPTpol detector development program.

Bolometers measure changes in incident power through a corresponding change in the sensor temperature. As total power detectors, bolometers do not have any inher-
ent spectral sensitivity, making them well-suited to observations over wide bandwidths where total photon count is the key parameter. As we will see in Section 3.1.5, observing over a wide bandwidth improves sensitivity, which is crucial when making observations of the $\approx 10\%$ linear polarization of $\mu K$-scale fluctuations in the CMB.

Transition edge sensors are extremely sensitive thermometers by virtue of a steep drop in their resistance over a narrow range in temperature, as shown in Figure 3.1. A TES is fabricated from a resistive material that becomes superconducting at low temperatures, and the addition of normal metal to the sensor allows for control over the temperature at which the superconducting transition occurs, called the critical temperature, or $T_C$. In the case of SPTpol pixels, two different architectures are employed: the 90 GHz detectors use a bilayer of Mo and Au, with a $T_C$ set by the relative layer thicknesses, while the 150 GHz detectors use Al doped with Mn, where the doping fraction sets $T_C$. Minute changes in temperature translate into large changes in resistance because of the steep $R(T)$ profile, so monitoring resistance fluctuations is a sensitive probe of TES temperature fluctuations. As we will see in the following sections, the tightly-coupled interactions between the thermal properties of the bolometer and the electrical properties of the TES are non trivial and require a significant amount of design and testing effort in order to produce sensors appropriate for use in a CMB camera.

### 3.1.1 Steady-State TES Bolometer Properties

The principle of any bolometer is to convert fluctuations in incident optical power to a measurable signal response in the detector. In the absence of any feedback, changes in optical power change the temperature and therefore resistance of the TES, with the drawback of having a very non-linear response as a function of sensor’s location in its $R(T)$ curve. As we will see in Section 3.1.2, the so-called negative electro-thermal feedback (ETF) keeps a TES at a particular bias point in its $R(T)$ curve, greatly
 Resistance versus temperature of a MoAu bilayer transition edge sensor (TES). Above the transition, the TES acts as a normal resistor, with a very small slope, $\Delta R/\Delta T$. As the sensor cools, it undergoes a steep transition in its resistance from the normal value to its superconducting state. Within the transition, small variations in sensor temperature correspond to large variations in resistance. Measuring the sensor resistance allows it to be used as a precise thermometer. Coupling a TES to a mm-wave absorber then provides a mechanism for converting small fluctuations in optical power to fluctuations in resistance that can be measured precisely with modern electronics. Improving the stability and linearity of the bolometer. The improved performance possible with negative ETF make TES bolometers excellent detectors for SPTpol observations.

Photons are dissipated onto the sensor, changing its temperature according to $\Delta T = \Delta Q / C$, where

$$C \equiv \frac{dQ}{dT} \ [J/K] \ (3.1)$$

is the heat capacity. The power then flows to a low-temperature bath through a weak thermal link, which in SPTpol detectors is provided by long, thin SiN legs.
that suspend the sensor away from the thermal bath. The weak link supports a temperature difference between the sensor and bath of $\Delta T = \frac{P}{\mathcal{G}}$, where

$$\mathcal{G} \equiv \frac{P}{\Delta T} \quad [\text{W/K}]$$

(3.2)

is the mean thermal conductance through the link between the two temperatures separated by $\Delta T$. Note that this is distinct from $G = \frac{\partial P}{\partial T}$, the dynamic conductance for small fluctuations around the steady-state power and temperature, defined at the sensor temperature, $T_C$. The combination of the heat capacity and the dynamic conductivity sets a natural time constant for the sensor response to a fluctuation in power:

$$\tau = \frac{C}{\mathcal{G}},$$

(3.3)

which parametrizes the exponential response to a delta function input of power. In the limit that the TES $R(T)$ is flat, i.e. a simple resistor, its temperature response is fully described by Equation 3.3.

In the voltage-biased state used by the SPTpol readout, fluctuations in in the optical power incident on the bolometer are tracked by monitoring the current through the TES. Section 3.1.3 describes the so-called TES responsivity, $\frac{\partial I}{\partial P}$ that maps power fluctuations to TES current fluctuations. The basic TES readout circuit is shown in the left panel of Figure 3.2. A voltage bias $V_{bias}$ is set across $R_{TES}$ by placing it in parallel with a shunt resistor $R_S \ll R_{TES}$, which draws the majority of the current from a source feeding the bias circuit. The Thevenin-equivalent circuit is shown in the right panel of Figure 3.2, and it combines the shunt resistance with any stray resistance in series with the TES into the equivalent load resistance, $R_L = R_{sh} + R_{stray}$. This circuit is useful for deriving the equations describing TES detector performance, and will be referred to for the remainder of this section.

The SPTpol readout circuit shown in Figure 2.8 differs from the basic circuit...
Figure 3.2: Simplified TES Sensing Circuits

*Left:* A basic voltage-biased TES circuit, with a shunt resistor to draw the majority of the bias current and therefore set up a stable a voltage bias across the TES. An inductor in series with the TES is used to convert current through the TES to a magnetic field coupled to a SQUID pre-amplifier, as described in Section 2.4. *Right:* A Thevenin-equivalent of the TES sensing circuit, where the parasitic stray resistance in series with TES is combined with the shunt resistance into a single, series load resistance, $R_L$. For the remainder of this chapter, we will be referring to this circuit, except where noted. For the purposes of discussing basic TES bolometer performance, the DfMux LRC circuits can be approximated by what is shown here.
shown in Figure 3.2 in that the bare $R_{TES}$ is replaced in the former with the series combination of $L_i$, $C_i$, and $R_{TES}$. When channel $i$ is biased with a sinuisoid at frequency $\omega_i = \sqrt{1/LC_i}$, the combined impedance of the $L$ and $C_i$ components goes to zero and the fMux circuit reduces to the basic circuit used here.

Sensing the TES resistance by sending current through it with a constant voltage bias leads to Joule heating power:

$$P_e = \frac{V_{bias}^2}{R_{TES}} = I_{TES}^2 R_{TES}. \quad (3.4)$$

The total power sensed by the bolometer $P_0 = P_{opt} + P_e$, where $P_{opt}$ is the optical power coupled to the TES and $P_e$ is the Joule heating power supplied through the bias circuit. In the steady-state limit, the total power absorbed by the sensor raises its temperature, $T_{TES}$, relative to the bath, $T_B$, according to:

$$P_0 = \kappa (T_{TES}^n - T_B^n), \quad (3.5)$$

where the exponent $n$ reflects the mechanism of heat transport through the thermal link. For normal metals, $n = 1$, while for insulators like the SiN used in SPTpol detectors, $n \approx 3$. In the steady-state limit, where both $T_{TES}$ and $T_B$ are stable, the heat transport equation can be simplified as

$$P_0 = P_e + P_{opt} = \overline{G} (T_{TES} - T_B), \quad (3.6)$$

where $\overline{G}$ is the average thermal conductance between the TES and thermal bath.

For small fluctuations about the steady-state TES temperature such as those induced by fluctuations in optical power, the general equation $G \equiv \partial P/\partial T$, gives us $G(T) = n\kappa T^{n-1}$. Solving where $T_{TES} \simeq T_C$, we find the dynamic $G$ that is important for modeling the small-signal, frequency-dependent thermal behavior of the TES.
At values of $P_0$ sufficient to drive $T_{TES} > T_C$, the TES behaves as a normal resistor with a “normal resistance” $R_n$. In this state, $R(T)$ is flat, and the dynamic properties of the TES are suppressed. In the narrow range of $P_0$ where $T_{TES} \simeq T_C$, the detector is said to be in its transition. The highest value of $P_0$ for which the TES is in the superconducting transition phase of its $R(T)$ profile is commonly referred to as the “turnaround power” $P_{\text{turn}}$, above which the detector is said to be “saturated” or simply “normal”.

While Equation 3.6 does not describe the response of the TES to small fluctuations such as CMB anisotropies, it is crucial for determining its ability to operate in a particular optical environment. Atmospheric and instrumental loading, detector bandwidth, and cryogenic base temperature can all be measured/estimated to provide an expected $P_{\text{opt}}$. From there, appropriate values of $G$ and $T_C$, which are set by the detector geometry, are chosen to provide a suitable ratio of $P_e : P_{\text{opt}}$. If the ratio is too small, optical power fluctuations will drive the TES normal. If it is too large, the thermal fluctuation noise (TFN), or “G noise” described in Section 3.1.5 will be unnecessarily large. We design our detectors to have a large enough $G$ for stable, linear operation under expected loading conditions, but as small as possible given that constraint and a reasonable safety margin. For SPTpol detectors, a typical target $P_{\text{sat}}$ is between 25-40 pW.

### 3.1.2 Electro-Thermal Feedback

The combination of electrical and optical power in the total power measured by the bolometer results in a coupling of thermal and electrical properties of the TES. With a voltage bias, this coupling takes the form negative electro thermal feedback (ETF). Negative ETF refers to a difference in sign between fluctuations in optical and electrical power, as an increase (decrease) in the former drives a decrease (increase) in the latter. For example, an increase in incident power $\delta P_{\text{opt}}$ changes the right
hand side of Equation 3.6 and drives up the temperature of the TES, which in turn increases $R_{TES}$. Since $P_e = V_{bias}^2 / R_{TES}$, the electrical power decreases, counteracting the total power increase from $\delta P_{opt}$. At a fixed $V_{bias}$, the change in $R_{TES}$ induces a change in $I_{TES}$, which is the quantity measured directly by the sensor readout circuit. Understanding and controlling the details of how $I_{TES}$ responds to fluctuations in $P_{opt}$ is central to the development and testing process that will be discussed in the following sections.

### 3.1.3 TES Responsivity

In order to make use of TES bolometers as the pixels in the SPTpol camera, we must characterize the transduction of optical power fluctuations to readout current fluctuations, known simply as the power-to-current responsivity:

$$s_I(\omega) \equiv \frac{\partial I}{\partial P}.$$  \hspace{1cm} (3.7)

In order to derive the responsivity, I will follow closely the treatment in [24], with simplifying assumptions made to reflect the design parameters of SPTpol bolometers and the SPTpol DfMux readout, discussed in Section 2.4.

Responsivity is a function of the sensor’s electrical and thermal properties and the properties of the electrical readout circuit, and is effectively a measure of the spectral characteristics of the ETF induced by optical power fluctuations. The underlying assumption of the following analysis is that the power fluctuations are small and the responsivity is adequately described by first-order expansions of the detector parameters about their steady state values. Given the $\approx \mu K$ CMB power fluctuations with a background loading of $\approx 10K$, the assumption of small signal perturbations is reasonable.

We start with the two differential equations that describe the behavior of the
TES in response to changes in its thermal and electrical properties. The thermal
differential equation describing the evolution of fluctuations from the steady state
described is:

\[ C \frac{dT}{dt} = -P_0 + P_e + P_{opt}, \quad (3.8) \]

where \( P_0 \) is the total power defined in Equation 3.5. The electrical differential equation
describing the evolution of changes in current through the squid coupling inductor
shown in Figure 3.2 is:

\[ L \frac{dI}{dt} = V_b - (IR_L + IR(T, I)). \quad (3.9) \]

We then expand \( R_{TES} \) in terms of small perturbations in the TES temperature and
the bias current flowing through it:

\[ R_{TES}(T, I) = R_0 + \frac{\partial R}{\partial T} \delta T + \frac{\partial R}{\partial I} \delta I \quad (3.10) \]

Taking the derivative of the logarithm of \( R(T) \) results in the conventional parametriza-
tion of the TES transition steepness:

\[ \alpha \equiv \frac{T_0}{R_0} \frac{\partial R}{\partial T}, \quad (3.11) \]

where the derivative is evaluated at \( I_0 \), the steady-state bias current value \( (I_0 = V_{bias}/R_0) \). Hereafter, all zero subscripts refer to the steady-state values of a “biased”
TES, i.e. a sensor that has been tuned with an appropriate bias voltage amplitude,
\( V_{bias} \) such that by Equation 3.6, \( T_0 \approx T_C \).

The same procedure for the current-dependent term in Equation 3.10 evaluated
at \( T_0 \) yields

\[ \beta \equiv \frac{I_0}{R_0} \frac{\partial R}{\partial I}, \quad (3.12) \]

48
but for SPTpol detectors, the current-dependence parametrized by $\beta$ is very small and can reasonably be ignored.

In analogy to an op-amp run in negative feedback mode, we introduce a so-called loopgain,

$$\mathcal{L} \equiv \frac{P_e \alpha}{GT_0},$$

(3.13)

which reflects the strength of the ETF feedback at a given bias point for a detector with particular $\alpha$ and $G$. Recall that $G = \partial P/\partial T$ with $P$ from Equation 3.5, evaluated at $T_0$. As we will see, loopgain is one of the most important characteristics of a TES bolometer.

We now proceed to describe the thermal and electrical responses described in Equations 3.9 and 3.8 to small fluctuations in the TES. The thermal differential equation for small fluctuations is:

$$\delta P_0 = C \frac{d\delta T}{dt} = \delta P_e + \delta P_{opt} - G \delta T,$$

(3.14)

where we have made the substitution $\delta P_0 \simeq G \delta T$, where $G$ is the dynamic conductivity, $\partial P/\partial T$, evaluated at the bias point. From Equations 3.4 and 3.10 $\delta P_e$ introduces a coupling to the electrical bias circuit:

$$\delta P_e = 2I_0 R_0 \delta I + I_0^2 \left( \frac{\partial R}{\partial T} \delta T \right),$$

(3.15)

Referring back to the basic TES bias circuit shown in Figure 3.2 and Equation 3.9, the time evolution of a small fluctuations in the bias current through the TES is described by:

$$L \frac{d\delta I}{dt} = -(R_L + R_0) \delta I - I_0 \frac{\partial R}{\partial T} \delta T,$$

(3.16)

where I have neglected fluctuations in the voltage, $\delta V_B$. For a strong voltage bias, we assume that fluctuations $\delta V_{bias} \simeq 0$, and I will drop them here for clarity. However,
in principle, $\delta V_{bias}$ can be retained as a source term like the $\delta P_{opt}$ we are interested in.

We can rearrange Equations 3.14 and 3.16 and substitute our simplifications (3.11, 3.13) to arrive at two coupled, linear differential equations:

\[
\frac{C}{d} \frac{d\delta T}{dt} = -G(1 - \mathcal{L}) \delta T + 2I_0R_0\delta I + \delta P_{opt} \quad (3.17)
\]

\[
\frac{L}{d} \frac{d\delta I}{dt} = -\frac{\mathcal{L}G}{I_0} \delta T - (R_L + R_0)\delta I \quad (3.18)
\]

With the substitution of $\tau_{el} = L/R_0 \approx L/(R_0 + R_L)$ as the electrical time constant from the LR resonance in Figure 3.2, and recalling that $\tau = C/G$ the equations can be written in matrix form as:

\[
\begin{pmatrix}
\frac{d}{dt} C\delta T \\
\frac{d}{dt} L\delta I
\end{pmatrix} =
\begin{pmatrix}
(1 - \mathcal{L}) & -2I_0 R_0 \\
\mathcal{L}G & (R_L + R_0)
\end{pmatrix}
\begin{pmatrix}
\delta P_{opt} \\
0
\end{pmatrix} +
\begin{pmatrix}
\frac{1}{\tau} \\
\frac{1}{\tau_{el}}
\end{pmatrix}
\begin{pmatrix}
C \delta T \\
L \delta I
\end{pmatrix} \quad (3.19)
\]

which can be simplified to:

\[
\frac{d}{dt} v = -Av + s. \quad (3.20)
\]

Our ultimate goal is to parametrize the responsivity, $\frac{\partial \delta I}{\partial \delta P_{opt}}$, as a function of fluctuation frequency, so we replace an arbitrary optical power fluctuation with a sinusoidal variations, $\delta P_{opt} \rightarrow \delta P_{opt} e^{i\omega t}$, and rearrange Equation 3.20 to get:

\[
\begin{pmatrix}
\delta P_{opt} \\
0
\end{pmatrix} e^{i\omega t} = \left( \frac{d}{dt} + A \right) v e^{i\omega t}. \quad (3.21)
\]

We can introduce a change of variables $v \rightarrow \tilde{v}$, such that $v$ is transformed into a linear combination of the eigenvectors $u_{\pm}$ of $A$ with associated pre-factors ($F_{\pm}$).

\[
\begin{pmatrix}
C \delta T(\omega) \\
L \delta I(\omega)
\end{pmatrix} = (F_+ u_+ + F_- u_-) e^{i\omega t}. \quad (3.22)
\]
The eigenvectors of $A$ and their associated eigenvalues are:

$$u_{\pm} = \begin{pmatrix} \frac{I_0}{\tau} (\lambda_{\pm} - \frac{1}{\tau_e}) \\ 1 \end{pmatrix}$$

$$\lambda_{\pm} = \frac{\text{tr}(A) - \sqrt{\text{tr}(A)^2 - 4|A|}}{2},$$

(3.23) (3.24)

Solutions to the homogeneous form of Equation 3.20 take the basic form of $a_i e^{-\lambda_i t}$, from which we can draw the interpretation of the eigenvalues in Equation 3.24 as inverse time constants describing the rise and fall times of the detector response to an impulse fluctuation in optical power, $\lambda_{\pm} = 1/\tau_{\pm}$. With this interpretation, and accounting for the assumptions invoked above that $R_{TES} \gg R_L$, we can approximate the two time constants as the high-frequency response roll off of the TES due to the bias circuit bandwidth, $\tau_+ \to \tau_e \approx \frac{L}{R_{TES}}$, and the settling time of the detector in response to a power fluctuation, $\tau_- \to \tau_{eff}$. The effective thermal time constant, $\tau_{eff}$, describes the effect of the ETF loopgain on the thermal response of the TES, according to:

$$\tau_{eff} = \frac{\tau}{\tau_e + 1}.$$  

(3.25)

Qualitatively, the electrical time constant describes the suppression of electrical signals at high frequency due to the LR impedance in the bias circuit. In SPTpol, the LRC bandwidth ($\delta \omega = 2L_{max}/R$) is much narrower than the $L_{SQ}/R$ from the TES and the coupling inductor, but the physical outcome is the same. The effective thermal time constant describes the decreased thermal response time of the sensor under the influence of negative ETF; compensation effect of the ETF reduces the actual energy increase of the TES in response to a change in optical power.

Using the above substitutions, Equation 3.21 is then transformed to:

$$\begin{pmatrix} \delta P_{opt} \\ 0 \end{pmatrix} = \begin{pmatrix} F_+ (1 + i\omega \tau_e) \\ \frac{I_0}{\tau_e} (\tau_e - \frac{1}{\tau_e}) \end{pmatrix} + \begin{pmatrix} F_- (1 + i\omega \tau_{eff}) \\ \frac{I_0}{\tau_{eff}} (\tau_{eff} - \frac{1}{\tau_e}) \end{pmatrix}.$$  

(3.26)
Solving first for the pre-factors $F_\pm$ in terms of $\delta P$ and then solving 3.22 for $\delta I(\omega)$ and $\delta T(\omega)$ in terms of the power fluctuation $\delta P_{opt}$ for a given frequency mode, we arrive at our equations for the power-to-current and power-to-temperature responsivities, $s_I(\omega)$ and $s_T(\omega)$, respectively:

$$s_I(\omega) = -\frac{1}{V_0} \frac{\mathcal{L}}{\mathcal{L} + 1} \left( \frac{1}{1 + i\omega \tau_{eff}} \right) \left( \frac{1}{1 + i\omega \tau_e} \right)$$

(3.27)

$$s_T(\omega) = \frac{1}{G} \left( \frac{1}{1 + \mathcal{L}} \right) \left( \frac{1}{1 + i\omega \tau_{eff}} \right)$$

(3.28)

At high loopgains, $\mathcal{L}/\mathcal{L} + 1 \rightarrow 1$ and frequencies below $1/\tau_{eff}$ and $1/\tau_e$, $s_I(\omega) \rightarrow -1/V_0$, and the current response to changes in optical power is perfectly linear. This state is ideal for camera pixels, since the details of the sensor state do not change its gain. This is reflected a different way in $s_T(\omega)$, which goes to zero as loopgain increases. In other words, at high loopgain, the TES temperature becomes insensitive to optical power fluctuations, as the ETF compensation effectively corrects the total power load faster than any imbalance can register as a change in the TES temperature.

High loopgain has an important negative consequence. Recalling the eigenvalue Equation 3.24, in order for the eigenvalues $\lambda_\pm$ to represent a decay in the detector responsivity in homogeneous solutions to Equation 3.20, they must be real and positive. Otherwise, power fluctuations either induce oscillations or exponential growth in the detector response, driving the TES out of its transition. This imposes the constraints $\text{tr}(A)^2 > 4|A|$ and $\text{tr}(A) > \text{tr}(A)^2 - 4|A|$ for stable, overdamped responsivity. Solving for the particular case of $\lambda_+ = \lambda_-$ and assuming moderate loopgain ($\mathcal{L} \gg 1$), we find the limit for stability for a given electrical bandwidth and natural thermal bandwidth:

$$\mathcal{L} < \frac{\tau}{5.8\tau_e}.$$  

(3.29)

Equation 3.29 is a key design boundary for TES detectors. Readout considerations
(multiplexing factor, signal bandwidth, readout bandwidth) drive the value of $\tau_e$, which for SPTpol is $\sim 0.05$ ms. This reduces the stability criterion to $0.3L < \frac{C}{G}$, and while operating at high loopgains is desirable, tuning $C$ and $G$ to provide a stable natural time constant can be challenging.

### 3.1.4 Extensions to the Basic TES Architecture

Developing detectors that adhere to constraints derived in Section 3.1.3 is non-trivial. In this section I will discuss the SPTpol extensions to the basic TES architecture, developed to adhere to the stability requirement of Equation 3.29 within the practical limits of tuning $\tau$ for our detectors. Given an optical loading budget, based on a nominal observing bandwidth and detector efficiency, shown for SPTpol in Table 3.1, both the DC optical power from Equation 3.6 and the noise contribution from background photon loading (discussed in Section 3.1.5) can be calculated. The former sets a minimum feasible $G$, while the latter sets a maximum value that ensures photon loading is the largest contribution to the noise budget.

For SPTpol-type detectors, $G$ is determined by the length and width of the SiN legs from which the sensor island is suspended, so it can be tuned to a wide range of values within the limits set by our loading and noise constraints. There is some design freedom in tuning the TES heat capacity, but even then, achieving a target $\tau$ is not as simple as adding more TES material, as the sensor geometry and chemistry determine not only $\tau$ but also $R_{TES}$ and $T_C$.

An inherent assumption of the basic TES model is that the heat capacity is contained entirely in the TES structure itself and that a single heat capacity coupled to a single thermal link describes the entire system. In this basic structure, the only way to increase $C$ for a particular material is to make the TES bigger. Simply adding more material is impractical, given that doing so affects other important parameters. An alternate solution employed first by SPT-SZ and again in SPTpol is the addi-
tion of a dedicated heat capacity, known as “bling”, to the sensor. Bling is made of a higher heat-capacity material than the TES (PdAu for SPTpol detectors) to increases $C$ rapidly with bling volume, while being deposited on the detector in such a way that it does not affect other sensor parameters, namely $R_{TES}$ and $T_C$. Adding bling introduces extra connections between thermal elements relative to the basic TES model, which can have a significant impact on detector performance and complicate the model derived in Section 3.1. A thorough treatment of generic multi-body TES performance can be found in Appendix A of [33]. However, subject to that analysis framework, the eventual goal of adding bling is to increase $C$ without altering the other TES parameters.

Adding bling to the detector island introduces an additional heat capacity, $C'$, and thermal link, $G'$, to the system, as shown in Figure 3.3. Both 90 GHz and 150 GHz detectors had bling added, and a thorough description of the testing program for the 150 GHz detectors can be found in [16]. The details of the thermal link between a PdAu bling and TES are sufficiently complicated that an iterative testing program involving different test geometries is required to effectively tune detector performance. Qualitatively, we want the bling and TES to act a single lumped element in response to fluctuations at all frequencies within the LCR bandwidth. An insufficiently large $G'$ results in the bling and TES thermal responses decoupling at high frequencies. In essence, the slowing effect of the bling is lost and the TES is able to respond to high-frequency fluctuations according to its base time constant, causing instability at target loopgains.

A second extension to the basic design is to add superconducting structures on the TES itself in order to tune the left side of Equation 3.29 independently. Such structures modify the interaction between the superconducting and normal materials, widening the temperature range over which the superconducting transition occurs, without changing its central temperature. The net effect of such a broadening is to
reduce $\alpha$, providing a means of controlling $L$ independent of other TES parameters. As with bling designs, analytical treatments are less effective than experimental iteration for determining the ideal shape and size of the structures.

Figure 3.3: Two-Body TES Detector Schematic
A schematic representation of the two-body TES architecture that includes bling to increase the detector heat capacity. Optical power is coupled through the PdAu absorber (for SPTpol 90 GHz detectors) in the middle of the SiN detector island, shown as the red box. It then flows to the edge of the island where the TES and bling are located. The thermal connection to the island for both the TES and bling dominate all other conductivities considered. Electrical power is deposited directly onto the TES, while optical power couples from the SiN island to both the TES and bling. Power flows through the weak thermal link, $G_0$, which has both an average $\overline{G}$ for steady-state power flow and dynamic $G$ for frequency-dependent power fluctuations.

3.1.5 TES Noise
The fundamental TES bolometer sensitivity limit is set by the noise inherent in the statistics of arriving photons. If the total noise of a deployed detector is dominated by the contribution of this statistical photon noise, the detector is said to be “background
limited”. Operating in the background limit is a common design criterion for CMB experiments, as it provides a convenient target for detector design parameters once basic experimental performance details like observing bandwidth and loading environment have been chosen. SPTpol detectors are designed to be background limited under the typical observing conditions at the South Pole, summarized in Table 3.1, with nominal bandwidths shown in Figure 2.2 and high per-detector optical coupling efficiency. In this section, I will describe the sources of noise that we consider while designing and developing our detectors.

### Photon Noise

Optical power in the mm-wave spectrum comes from the CMB and other celestial sources, along with thermal emission from the atmosphere and surfaces in the optical path of the telescope. A useful way of paramatrizing detector noise is with noise-equivalent power (NEP), which is defined as the size of a fluctuation in incident power on the detector that can be resolved with a S/N of 1 in 1 s of integration time [41]. Given a blackbody source at a temperature $T$, illuminating a detector over a frequency band $\nu$ with an optical throughput of $A\Omega$, the total optical power incident on the detector with a spectral response $\zeta(\nu)$ from all sources in the optical path is

$$P_{opt} = \eta \sum_i \epsilon_i \tau_i \int_\nu A\Omega \frac{B(\nu, T_i)}{2} \zeta(\nu) d\nu,$$

where $\frac{1}{2}B(\nu, T) = \frac{h\nu}{\lambda^2} \frac{1}{e^{\frac{h\nu}{k_B T}} - 1}$ is the Planck blackbody intensity for a single polarization state, $\epsilon_i$ is the emissivity of each source and $\eta$ is the optical coupling efficiency between a single bolometer and the total incident power on it. The factor $\tau_i$ is the product of the transmissivities of all subsequent elements in the optical path from source $i$ to the detector. In a single-moded, diffraction-limited optical system such as SPT, $A\Omega = \lambda^2$. Thus, from $B(\nu, T)$, the number of photons with energy $h\nu$ absorbed
by the detector in a second of integration time from a single source is:

\[
    n_\gamma = \eta \tau \epsilon \frac{1}{e^{\frac{h\nu}{kT}} - 1}.
\]

(3.31)

From [26], the variance in the number of photons arriving per second is given by:

\[
    \langle \Delta n_\gamma^2 \rangle = n_\gamma (1 + \varepsilon n_\gamma),
\]

(3.32)

where \( \varepsilon \) has been introduced to account for the degree of “bunching” in the arriving photons. In the Wein limit when \( \frac{h\nu}{kT} \gg 1 \), \( \varepsilon \rightarrow 0 \) and the photon arrivals obey Poisson statistics, while for Rayleigh-Jeans sources when \( \frac{h\nu}{kT} \ll 1 \), \( \varepsilon \rightarrow 1 \) and the variance term approaches \( n_\gamma^2 \). For ground-based CMB observations where the photon sources are the CMB and the low-emissivity atmosphere and telescope optics, which span the temperature range of 0.25 - 250 K, neither limit is universally applicable. For a basic calculation of photon noise, \( \varepsilon = 1 \) provides the worst-case noise limit, while the true value depends on the relative contributions of loading sources, their temperatures, and their emissivities, and the optical efficiency at the detector.

The NEP from a single source for an SPT-like optical system (diffraction-limited, single-moded) is given by:

\[
    NEP_\gamma^2 = \int h^2 \nu^2 \langle \Delta n_\gamma^2 \rangle d\nu \left[ \frac{W^2}{Hz} \right],
\]

(3.33)

which expresses the variance in the number of arriving photons times the energy per photon. The photon loading budget for the 2012 season of SPTpol observations is given in Table 3.1. The bandwidths and detector optical efficiencies used are nominal values based on design targets and simulations of detector architecture. The resulting \( NEP_\gamma \) values for a typical SPTpol 90 (150) GHz detector are 65 (55) \( \text{aW}/\sqrt{\text{Hz}} \).
Readout Noise

Various components in the current readout system, from the SQUID through the warm electronic components, produce statistical noise of their own. These noise components can be referred back to the power fluctuations in the detector that would have produced them, in order to make a consistent comparison with other forms of TES-localized noise. Because the actual origin of the noise is not on the detector, it is assumed that the detector properties are constant when referring the readout noise, such that for a given current noise PSD, $S_I(\omega)$, in the readout electronics, the resulting readout NEP is:

$$NEP_{\text{readout}}^2 = V_0^2 S_I(\omega) \left[ \frac{W^2}{Hz} \right].$$

(3.34)

Given an operational voltage bias, $V_0$, the readout noise can be controlled by the choice of components used in the signal readout path, to ensure that its contribution is sub-dominant to other sources of noise.

Johnson Noise

Thermal fluctuations in the resistive components of the TES readout circuit (including the TES itself) result in Johnson noise terms that take the form of a temperature-dependent voltage fluctuation in series with an ideal resistor. The power spectral density of the voltage fluctuations from Johnson noise in the resistors is

$$S_V^2(\omega) = 4k_BTR.$$

(3.35)

The impedance of the basic TES readout circuit from Figure 3.2 in series with the Johnson noise voltage source is $R_s + Z_{TES}(\omega) + i\omega L_s$. The impedance of the TES, in the context of a noise calculation, the ratio of a voltage fluctuation to the resulting
fluctuation in TES current, depends on where the voltage fluctuation occurs. If it is external to the TES, the result is a perturbation to the bias on the TES, whereas noise internal to the TES dissipates power that then directly affects its thermoelectric properties. From [24], the two impedances are

\[ Z_{\text{ext}}(\omega) = \frac{\mathcal{L}}{\mathcal{L} - 1 \cdot s_1(\omega) I_0 (1 + i\omega \frac{\tau}{1 - \tau^2})}, \] (3.36)

and

\[ Z_{\text{int}}(\omega) = \frac{1}{-s_1(\omega) I_0 (1 + i\omega \tau)}. \] (3.37)

The Johnson noise spectrum in voltage, \( S_V^2 \), can then be transformed into a current noise spectrum,

\[ S_I^2(\omega) = \frac{S_V^2(\omega)}{|Z_{\text{TES}}(\omega)|^2}, \] (3.38)

which can then be referred back to a noise at the TES by dividing by the square of the responsivity. In the moderate loopgain limit where TES bolometers operate (\( \mathcal{L} \approx \mathcal{L} + 1 \)), the resulting NEP contribution from the external resistance in series with the TES, \( R_S \), is then:

\[ \text{NEP}_{j,R_S}^2 = 4k_B T_R S I_0^2, \quad [W^2/Hz] \] (3.39)

and for the noise in the TES itself:

\[ \text{NEP}_{j,R_{\text{TES}}}^2 = 4k_B T_0 P_e \left(1 + \omega^2 \tau_e^2\right) \frac{1}{\mathcal{L}^2} \left[\frac{1}{W^2/Hz}\right]. \] (3.40)

The NEP for the Johnson noise originating from TES resistance fluctuations is suppressed by ETF, subject to the bandwidth of the detector responsivity. At lower frequencies and higher loopgains, the suppression is strongest. Note that the treatment here supposes a DC bias, which differs from the circumstances present in a
readout system the SPTpol DfMux. However, in Appendix A of [33], it is shown for the component of the TES current that is demodulated in phase with the carrier input, which is what we retain as our detector signal, Johnson noise is suppressed as in the DC system.

**Thermal Fluctuation Noise**

For a heat capacity $C$ in thermal contact with a bath through a thermal link $G$, it can be shown as in [26, Chapter 3], that the system is subject to white noise energy fluctuations with a variance of $\langle \Delta E^2 \rangle = k_B T_0^2 C$. These energy fluctuations are related to temperature fluctuations in the TES by $C \equiv \frac{\Delta E}{\Delta T}$, such that the variance in the temperature fluctuations is just $\langle \Delta T^2 \rangle = \frac{k_B T_0^2}{C}$, and the variance is related to an integrated spectrum of fluctuations, $S_T(\omega)$, by:

$$\langle \Delta T^2 \rangle = \frac{1}{2\pi} \int S_T(\omega) d\omega.$$  \hspace{1cm} (3.41)

Since we are referring noise terms to the power on the TES, we are interested in the spectrum of power fluctuations associated with the temperature fluctuations. The underlying process is a change in the energy of the bolometer system, so the spectrum of fluctuations in power is flat like that of the energy fluctuations. However, the bolometer temperature is related to the power incident on it by the temperature-to-power responsivity (3.28), which can be reduced in the zero loopgain (i.e. no electrothermal feeback) limit with $\tau_+ = \tau_- = \tau$ to

$$s_T(\omega) = \frac{\Delta T}{\Delta P} = \frac{1}{G(1 + i\omega\tau)}.$$  \hspace{1cm} (3.42)

Because the source fluctuations are actually in the temperature of the TES, they are naturally referred to changes in the power flowing through $G$, so we do not include an ETF term to refer the noise to an NEP. Using (3.42), we get the thermal fluctuation
NEP from the relation \( S_T(\omega) = S_P(\omega)(|s_T(\omega)|^2) \), where again \( S_P \) is a white noise spectrum. Combining (3.41) and (3.42) and distributing the fluctuations over the bandwidth set by the detector natural time constant \( \tau = \frac{C}{G} \), we recover the NEP due to thermal fluctuation noise,

\[
NEP_{TFN}^2 = 4k_B^2 T_0^2 G [W^2/Hz].
\]  

(3.43)

### 3.1.6 Noise Equivalent Temperature

It has been convenient to this point to refer noise to the power fluctuations in the detector, and indeed units of power on the TES are the common reference during detector testing and development. However, our interest lies in the sensitivity of a detector to temperature fluctuations of the CMB, so we refer the amplitude of absorbed power fluctuations than can be measured with signal-to-noise of 1 to corresponding fluctuations in source temperature required to induce them. Thus for a particular detector NEP, the corresponding NET comes from:

\[
NET = \frac{NEP_{TES}}{\frac{dP}{dT}} [K/\sqrt{Hz}].
\]  

(3.44)

The precision of a measurement of the CMB at a point on the sky is function of how long that spot is observed, as noise fluctuations integrate down while the CMB signal remains constant. Thus, it is common to refer to NET in terms of integration time rather than bandwidth, which brings a factor of \( \sqrt{2} \) to Equation 3.44, since the same noise level is integrated for two seconds to get its equivalent level in one Hz of bandwidth. From Equations 3.30 and 3.31, the derivative (where \( A\Omega = \lambda^2 \) and we consider a single polarization per detector) can be rewritten as:

\[
NET = \sqrt{2} \frac{NEP_{TES}}{\eta \int \frac{d}{dT} \left( h\nu n_\gamma(\nu, T) \right) d\nu} [K\sqrt{s}].
\]  

(3.45)
Here, the efficiency, $\eta$, in the denominator includes the accounting of all the component $\tau_i$ factors from Equations 3.30 and 3.31. For Rayleigh-Jeans sources like the atmosphere where $\frac{h\nu}{kT} \ll 1$, $h\nu n_\gamma(\nu, T) \rightarrow k_B T$, so the NET is independent of source temperature and is inversely proportional to the product of detector bandwidth and optical efficiency. Even though the CMB is not a Rayleigh-Jeans source in the mm waves, necessitating a full calculation of $\frac{dn_\gamma(\nu, T)}{dT}$ at $T_{CMB}$ to find the NET in units of equivalent CMB fluctuations, the inverse relationship between NET and the efficiency-bandwidth product still holds, and is a primary driver of experiment design.

3.1.7 SPTpol Noise Budget

Given the choice of observing in the atmospheric windows shown in Figure 2.2, the nominal target observing bandwidths for SPTpol were 30 (45) GHz for the 90 (150) GHz bands. Simulations of our detector architectures suggested optical coupling efficiencies of $\sim 90\%$, while an optical loading profile expected for typical observing conditions on the telescope is shown in Table 3.1. Estimates of these parameters, along with the known cryogenic performance of the telescope and other similarities between the SPTpol and SPT-SZ experiments, were sufficient to derive target detector parameters for SPTpol detectors, shown in Table 3.2.

3.2 Development and Testing of SPTpol Detectors

Development of SPTpol pixels happened largely in parallel as the 90 GHz and 150 GHz detector programs were centered at two different institutions. Argonne National Laboratory (ANL) fabricated the 90 GHz detectors while the 150 GHz arrays were built at National Institute of Standards and Technology (NIST). My work was largely in coordination with ANL and the University of Chicago on the development and
<table>
<thead>
<tr>
<th>Loading Source</th>
<th>$T$ (K)</th>
<th>$\epsilon$</th>
<th>$\eta_{\text{cum}}$</th>
<th>$P_{\text{opt}}$ (pW)</th>
</tr>
</thead>
<tbody>
<tr>
<td>CMB</td>
<td>2.73</td>
<td>1.00 (1.00)</td>
<td>0.457 (0.344)</td>
<td>0.21 (0.12)</td>
</tr>
<tr>
<td>Atmosphere</td>
<td>230</td>
<td>0.12 (0.08)</td>
<td>0.520 (0.372)</td>
<td>5.90 (3.79)</td>
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<tr>
<td>Primary window</td>
<td>220</td>
<td>0.02 (0.02)</td>
<td>0.530 (0.380)</td>
<td>0.96 (0.99)</td>
</tr>
<tr>
<td>Environmental window</td>
<td>300</td>
<td>0.01 (0.01)</td>
<td>0.536 (0.384)</td>
<td>0.66 (0.68)</td>
</tr>
<tr>
<td>100 K filters</td>
<td>100</td>
<td>0.09 (0.09)</td>
<td>0.589 (0.422)</td>
<td>2.15 (2.19)</td>
</tr>
<tr>
<td>10 K filters</td>
<td>10</td>
<td>0.05 (0.05)</td>
<td>0.620 (0.444)</td>
<td>0.10 (0.09)</td>
</tr>
<tr>
<td>Secondary mirror</td>
<td>10</td>
<td>0.12 (0.37)</td>
<td>0.704 (0.704)</td>
<td>0.28 (1.08)</td>
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<tr>
<td>Receiver filters</td>
<td>6</td>
<td>0.09 (0.09)</td>
<td>0.774 (0.774)</td>
<td>0.12 (0.13)</td>
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<tr>
<td>Sub-K filters</td>
<td>0.300</td>
<td>0.14 (0.14)</td>
<td>0.900 (0.900)</td>
<td>0.00 (0.00)</td>
</tr>
<tr>
<td>Total</td>
<td></td>
<td></td>
<td></td>
<td>10.4 (9.07)</td>
</tr>
</tbody>
</table>

Table 3.1: Principle Sources of Optical Loading at 90 (150) GHz

Optical loading budget for elements in the optical system, based on the first year SPTpol observations in 2012. Additional loading terms are introduced by scattering, but the telescope shielding and tilting of secondary optical filters reduces the coupling to scattered radiation from warm optical elements by an order of magnitude. Scattering off of the secondary mirror and receiver optics contributes negligibly to the total loading. The numbers shown here assume flat detector response bands with nominal edges, as shown in Figure 2.2, and the estimated optical efficiency from simulations of the detector architecture. The cumulative efficiency $\eta_{\text{cum}}$ accounts for the fact that each surface, in addition to supplying an additional loading term, only transmits a fraction of the power incident on it.

testing of 90 GHz detectors, so they will be the focus of the following sections, in which I will discuss the chief aspects of the detector development program. However, most of the material is directly applicable to the 150 GHz detector development, and in fact I performed many of the same tests on 150 GHz prototypes. However, particularly as the development transitioned from the early prototyping phase to pre-deployment validation of finalized designs, the allotment of resources had Case Western testing capabilities dedicated solely to preparing 90 GHz detectors for the final camera.

### 3.2.1 Optical Design

#### Optical Efficiency

The optical coupling efficiency, $\eta$, from Equation 3.30, defined as the fraction of the incident power in a single polarization state that is absorbed and measured by the
<table>
<thead>
<tr>
<th>Parameter</th>
<th>90 GHz Target</th>
<th>150 GHz Target</th>
</tr>
</thead>
<tbody>
<tr>
<td>Band center</td>
<td>90 GHz</td>
<td>145 GHz</td>
</tr>
<tr>
<td>Bandwidth</td>
<td>30 GHz</td>
<td>45 GHz</td>
</tr>
<tr>
<td>$\eta_{det}$</td>
<td>0.90</td>
<td>0.90</td>
</tr>
<tr>
<td>$\eta_{optics}$</td>
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<td>0.90</td>
</tr>
<tr>
<td>$T_C$</td>
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<td>470 mK</td>
</tr>
<tr>
<td>$R_0$</td>
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<td>0.85 Ω</td>
</tr>
<tr>
<td>$\overline{G}$</td>
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<td>130 $\mu W/K$</td>
</tr>
<tr>
<td>$n$</td>
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<td>3</td>
</tr>
<tr>
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<td>9.1 pW</td>
</tr>
<tr>
<td>$P_0$</td>
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<td>20 pW</td>
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<tr>
<td>$NEP_\gamma$</td>
<td>65 $aW/\sqrt{Hz}$</td>
<td>55 $aW/\sqrt{Hz}$</td>
</tr>
<tr>
<td>$NEP_{TFN}$</td>
<td>48 $aW/\sqrt{Hz}$</td>
<td>44 $aW/\sqrt{Hz}$</td>
</tr>
<tr>
<td>$NEP_J$</td>
<td>30 $aW/\sqrt{Hz}$</td>
<td>30 $aW/\sqrt{Hz}$</td>
</tr>
<tr>
<td>$NEP_{readout}$</td>
<td>44 $aW/\sqrt{Hz}$</td>
<td>33 $aW/\sqrt{Hz}$</td>
</tr>
<tr>
<td>$NET_{CMB}$</td>
<td>$500 \mu K_{CMB}\sqrt{s}$</td>
<td>$400 \mu K_{CMB}\sqrt{s}$</td>
</tr>
</tbody>
</table>

Table 3.2: SPTpol Target Noise Budget

Approximate target sensitivity values and associated detector property targets for the SPTpol detector development program. Details of the photon loading calculation are shown in Table 3.1. These values assume a simple top-hat detector frequency response, and a dynamic $G$ equal to the $\overline{G}$ values shown (a good approximation based on Equation 3.51). The $NEP_\gamma$ photon noise term sets the noise limit for other terms, which in turn constrains TES properties and readout design to ensure that photon noise is the largest component of the overall noise budget. Other NEP terms are based on estimates of target detector properties.
Figure 3.4: Detector Testing Setup
Cross-sections of two configurations of the testing set-up used for detector development. The outer blue layer is the 300 K vacuum enclosure, inside of which is a 77 K heat break stage, followed by the 4 K enclosure. Left: Optical testing set up, with the detector coupled to sources outside the cryostat. IR-blocking low-pass filters at the 77 K and 4 K breaks, along with an IR-shading reflective film at 4K, reduce the out-of-band loading inside the cryostat. For a 90 GHz detector, the band defining filter stack, mounted at 250 mK, consists of 7, 5.1, and 3.7 icm (inverse-centimeter) filters. The 3.7 icm filter is removed when testing a 150 GHz detector. Each of the low pass filters is approximately 95% transmissive and the IR shader is 99% transmissive in the range of our detector bands. Right: An alternate set up for performing low-loading tests. The IR blocking filters are mounted at 4 K in front of a source whose temperature can be tuned to between 4-30 K. The source is machined from a block of Al and its interior surface is coated in emissive epoxy.
Figure 3.5: TES Load Curves

TES “load curves”, showing the relationship between the voltage applied across the TES and the current flowing through it. Such curves are used extensively in measuring the performance parameters of prototype detectors. *Left:* Current through the TES as a function of bias voltage, called an “IV curve”. It begins with the voltage bias high enough to saturate the TES, corresponding to the right side of the plot. $V_{bias}$ is then decreased until the total power is low enough that $T_{TES} \simeq T_C$, at which point the TES drops through its superconducting transition at a nearly constant temperature and total power load. The power remains roughly constant as the reduced R allows more current to flow, shown in the up-turn in $I_{TES}$ as $V_{bias}$ continues to drop. *Right:* Converting the data shown in the left plot to TES resistance and electrical power dissipated on it. The up-turn in the current through the TES in the IV curve happens as the resistance drops. The shape of this so-called “RP” curve mirrors the TES $R(T)$ curve from Figure 3.1.
detector, is a critical parameter in the overall detector performance. A closer look at Equation 3.45 reveals that the numerator contains a factor of $\sqrt{\eta}$ from taking the square root of Equation 3.33, while the denominator contains a factor of $\eta$ from the derivative of Equation 3.30 with respect to CMB temperature. Thus, greater optical efficiency, while increasing the NEP due to a larger number of loading photons, results in a $\sqrt{\eta}$ reduction in the final NET, which is the fundamental figure-of-merit for detector performance. Furthermore, in order to accurately estimate detector optical loading and thus determine the optimal value for $G$, a reasonable estimate for $\eta$ is necessary.

In order to isolate $\eta$, we simplify Equation 3.30 to $P_{opt} = \eta P_{inc}$, where $P_{inc}$ is the total power incident on the detector from a known ensemble of sources. This allows us to write Equation 3.6 as

$$P_e + \eta P_{inc} = G \Delta T,$$  \hspace{1cm} (3.46)

where $G \Delta T$ does not vary for a given TES as long as $T_B$ is constant, which it is for the purposes here. If we can measure the change in $P_e$ as we change $P_{inc}$, Equation 3.46 can be rewritten as:

$$\eta = \frac{-\Delta P_e}{\Delta P_{inc}}.$$  \hspace{1cm} (3.47)

To extract $\eta$, we change $P_{inc}$ and measure the corresponding changes in $P_e$, using a consistent reference $R_{TES}$ in the detector transition to ensure we are comparing things at a constant $T_{TES}$ and therefore $P_0$. We vary the optical loading using a temperature-controlled thermal blackbody, coupled to the test device as shown in Figure 3.4. The source is machined from a block of aluminum and its detector-facing surface is coated in highly emissive ECCOSORB epoxy[15]. Its shape is designed to ensure a minimum of three scatterings off of a blackened surface for each ray originating in the detector beam. To control the detector spectral response and avoid out of band loading, a series of infrared blocking and band-defining filters shown in Figure 3.4. There are 5
low pass filters that are estimated to be 95% transmissive in our bands of interest and
one 99% transmissive IR shader, with a combined transmission \( \eta_{filt} = 0.77 \). Thus,
with an estimated detector spectral response, \( \varsigma(\nu) \), we can reasonably estimate the
source power as

\[
P_{\text{inc}} = \frac{1}{2} \eta_{filt} \int \varsigma(\nu) B(\nu, T) d\nu
\]

for a single linear polarization.

\[ (3.48) \]

Figure 3.6: Load Curves vs. Optical Source Temperature
A series load curves taken with varying temperature of a beam-filling optical black-
body load. Changing the load temperature varies \( P_{\text{opt}} \), while \( P_0 \) remains constant,
so comparing the changes in the measured electrical power at a given point in the
transition to expected changes in the optical power as \( T_{\text{load}} \) varies allows us to solve
for optical efficiency using Equation 3.47.

With the blackbody load set to a stable temperature between 5-30 K, we take a
load curve to measure \( P_e \). For each such curve, we calculate \( P_{\text{inc}}(T_{\text{load}}) \) from Equation
3.48 and extract \( P_e \) from the load curve at the chosen \( R_{\text{TES}} \) value. A series of R-P
curves vs. load temperature are shown in Figure 3.6. Figure 3.7 shows a “by eye”
fit of \( \eta \) to load curve vs. \( T_{\text{load}} \) data for two detectors measured before deployment.
The recovered estimate for optical coupling efficiency to a single polarization is 87%.
The result is in line with predictions from simulations of the detector design, which indicated an efficiency of $\eta \sim 0.90$ [38].

![Optical Efficiency for Crossed-Pair ANL Pixel](image)

**Figure 3.7: Detector Optical Efficiency Measurement**

Taking load curves at a series of optical source temperatures illustrates the linear relationship between changing optical power and the resultant change in electrical power. The slope of the relationship, illustrated here with a by-eye fit to show the approximate slope for two detectors in a crossed pair, is the optical coupling efficiency, or the fraction of incident optical power that is absorbed in the sensor.

**Polarization Efficiency**

The ability to precisely measure the faint CMB polarization signal depends on individual detectors having good discrimination of linear polarization. Temperature fluctuations in the CMB are one or two orders of magnitude greater than polarization fluctuations, so even a few percent confusion in rejecting unpolarized light completely dominates the polarization signal. In order to validate the polarization
fidelity of the detector design and fabrication, we illuminate test detectors with a chopped thermal source shined through a rotating wire grid polarizer with very low crosspolar transmission. Given a detector sensitivity direction, $\psi$, and a source polarization direction $\theta$, the detector response, $S$, as a function of the difference between the two directions, $\xi = \theta - \psi$ is

\[
S(\xi) = A \cos^2(\xi) + B, \tag{3.49}
\]

where $A$ is the modulation depth of the response and $B$ is the response when the polarized source direction is orthogonal to the detector sensitivity direction. A non-zero minimum response results either from coupling of the detector to cross-polar radiation or an imperfectly polarized source. Assuming a perfectly polarized source, we define polarization efficiency as

\[
\rho = \frac{A}{A + B}, \tag{3.50}
\]

which expresses the extent to which a detector can reject cross-polar radiation. For a proper beam-filling optical signal from a highly polarized source, a plot like the one shown in Figure 3.8 indicates better than 99% polarization efficiency.

**Detector Frequency Response**

In order to check that the detector response is limited to the expected frequency band, we measure its response spectrum with a Martin-Puplett interferometer employed as a Fourier Transform Spectrometer (FTS). The FTS has two input ports: one a blackened 300K surface and the other typically cooled to 77K. A beam splitter separates a light beam and sends half to a stationary roof mirror and the other half to a moving mirror. The light is then recombined and sent to an output port. As the translating mirror moves, the interference pattern of the recombined light is modulated at a
Measurements of polarization efficiency for two SPTpol 90 GHz detectors in their mounted configuration, with a nominal 90° relative rotation. The detector response amplitude is measured at each angle of the rotating polarized source, and a sinusoid fit to each, according to Equation 3.49. With the fit normalized to the maximum response as it is here, the polarization efficiency, \( \rho \), defined in Equation 3.50, is > 99%.

different audio frequency for each photon wavelength. With an appropriate mapping between the photon frequencies and the audio frequency of their modulations, the frequency transform of a detector’s response to the FTS output signal encodes its response to amplitude modulations of radiation at different frequencies.

Figure 3.9 shows the frequency response of a 90 GHz prototype pixel, indicating both that the in-band portion of the response is free from any drastic features and that out of band response is negligible. Final response bandwidths for the deployed instrument are determined by the band-defining low-pass filters at the upper edge and the circular waveguide diameter at the lower edge. However, for the purposes of
Figure 3.9: Detector Spectral Response
A measurement of the in-band coupling of a test detector using a Martin-Puplett interferometer as a Fourier transform spectrometer (FTS). Different mirror speeds alias photon-frequency modulations to different audio frequency bands, so the detector frequency response features will overlap for the different spectra, while audio-band noise will manifest as features at different photon frequencies for different mirror speeds.

testing other detector properties, particularly optical efficiency and noise levels, it is important to both verify response within the expected bandwidth and ensure there is no out-of-band response.

3.2.2 Thermal Design

The thermal conductivity of the bolometer is an important factor in both its stability and noise performance. The steady-state average conductance, $\overline{C}$, is determined by the geometry of the legs that separate the sensor island from the bath, which can be tightly controlled in the lithography process. The target value for $\overline{C}$ is bounded on the high side by detector noise performance requirements and on the low end by the requirement of robust performance under expected loading conditions. High values
of $G$ come with a noise penalty, as $NEP_{TFN}^2$ scales linearly with $G$ (Equation 3.43) and $NEP_j^2$ (Equations 3.39 and 3.40) scales with the bias current required to supply the $P_e$ component of $P_{sat}$. On the other hand, loading conditions at the South Pole impose a minimum $G$ necessary to avoid saturating the detector with background power alone, and at least a 1:1 of $P_e$ to $P_{opt}$ ensures robust detector performance through normal variations in weather. The change in DC atmospheric loading due to variations in airmass from elevation changes in a typical SPTpol observing field is near 20%.

Measuring detector $G$ values as a function of $\Delta T$ between the TES and the base fridge temperature follows a similar procedure to the $\eta$ measurements described in the previous section, but instead of changing the balance described by Equation 3.6 by varying $P_{opt}$, we vary $T_B$. More generally, we can start from Equation 3.5, recalling the approximation of $T_{TES} = T_C$ through the superconducting transition, and record $P_0(T_B)$. For these tests, detectors are not coupled to waveguides and are instead enclosed in a sealed, 250 mK enclosure, so $P_{opt} \approx 0$. With $P_0 = P_e$, load curves as a function of $T_B$ provide the necessary data to fit for $\kappa$ and $n$ in Equation 3.5, where $n$ is nominally 3 for SiN. Fitting the curve shown in Figure 3.10 allows us to directly calculate the dynamic $G$ at any base temperature, as well as providing a secondary estimate of $T_C$, which is usually measured simply by changing $T_B$ while supplying a small current to the TES and recording when its transition occurs.

### 3.2.3 Detector Stability

One of the chief areas of effort in the development of SPTpol detectors was the tuning of various parameters to ensure stable ETF operation within the constraints imposed by loading and readout. As discussed in Section 3.1.4, a bling structure was added to the detector island to increase the natural time constant $\tau$, reducing the inherent thermal response bandwidth and allowing for a higher loopgain while still adhering
Figure 3.10: Measurement of Thermal Conductivity

For a series of load curves taken at different values of $T_B$, the measured $P_e(T_B)$ values at a fixed $R_{TES}$ in the transition are plotted as the blue points. We enclose the test detectors in a 250 mK box so $P_{opt} \approx 0$ and $P_e \approx P_0$, allowing us to fit the measured $P_0(T_B)$ data to the model in Equation 3.5. Evaluating Equation 3.51 at $T = T_C$ with the fit values of $\kappa$ and $n$ from Equation 3.5 yields the operational dynamic $G$ at the detector operating point.

to the stability requirement set by Equation 3.29.

The most difficult part of adding bling is ensuring that it is sufficiently well thermally coupled to the TES. In other words, $G'$ between the bling and the TES must be much larger than the link from the TES to the bath, $G$. The schematic of the two-body TES model is shown in Figure 3.3, illustrating where power is deposited onto the detector island before flowing through the thermal link to the bath. As long as $G' \gg G$, the TES and bling act as a single heat capacity and the sensor behaves as
a single-element TES. We parametrize the conductivities by $G' = \gamma G$, such that the total frequency-dependent conductivity of the two-body system $G(\omega)$ can be written, following [16], as:

$$G(\omega) = G \frac{\gamma}{1 + \gamma} \left( \frac{1 + i\omega C'G}{1 + i\omega (1 + \gamma)G} \right),$$

(3.51)

where $C'$ is the combined TES-bling heat capacity, which is assumed to be dominated by the bling contribution, and $G$ is again the dynamic thermal conductivity, $\partial R_b/\partial T$ between the sensor and the bath. For large values of $\gamma$, the combined thermal conductivity reduces to the simple one-body TES approximation, $G(\omega) = G(1 + i\omega\tau)$. In the simple case of zero loopgain, $G(\omega) = s_T(\omega)^{-1}$ from Equation 3.28. As the frequency of power fluctuations increases, the TES thermally decouples from the bling, following the fluctuations according to the bare TES time constant, which is much smaller than the combined TES-bling time constant. Low values of $\gamma$ lower the frequency at which this bling decoupling occurs, while higher values ensure that the TES-bling combination remain thermally coupled to higher frequencies.

In order to probe the effective thermal structure of the multi-body TES, we employ an algorithm developed for SPT-SZ, which is described in [34]. This so-called “tickle” measurement allows us to inject a power modulation onto the sensor through the TES itself, mimicking an optical power modulation with a controlled amplitude and frequency. Using two carrier channels, we bias the detector with a carrier voltage $V_{bias} = V_b e^{i\omega t}$, with $\omega = \frac{1}{\sqrt{LC}}$, then inject a second signal at a nearby frequency, $\omega + \delta\omega$, with a much lower amplitude, $V_{tickle} = V_t e^{i(\omega + \delta\omega)t}$. The combination of the two signals results in a total electrical power on the detector that is, to first order:

$$P_e = \frac{V_b^2}{R_0} + \frac{2V_b V_t}{R_0} e^{i\delta\omega\tau_e} \cos(\delta\omega t),$$

(3.52)

which gives us a power fluctuation at a frequency $\delta\omega$, $\delta P = \frac{2V_b V_t}{R_0} \frac{e^{i\delta\omega\tau_e}}{1 + i\delta\omega\tau_e}$, that stands in for $\delta P_{opt}$ in the equations from Section 3.1.3. The $(1 + i\delta\omega\tau_e)$ term comes from the
fact that the electrical readout bandwidth attenuates the power fluctuation amplitude before it reaches the detector. Recalling the definition of responsivity as $s_I(\omega) \equiv \frac{\delta I}{\delta P_{opt}}$ and substituting the tickle power signal, we measure the amplitude of the output from the measurement as:

$$|I(\delta\omega)| = \frac{2V_b V_t}{R_0|1 + i\delta\omega \tau_e||s_I(\delta\omega)|},$$

where we are probing the responsivity to verify that $s_I(\omega)$ behaves as expected for a single-body TES bolometer. In particular, we wish to verify that the bling-to-TES thermal link is much stronger than the link to the bath, i.e. that $\gamma \gg \mathcal{L}$ from Equation 3.51. In that case, the form of $s_I(\omega)$ matches the simple case of 3.27, and the data fits a model with a loopgain-dependent thermal pole and a resistance-dependent electrical pole.

Early in the testing program, the CWRU test cryostat was configured to operate a single detector per SQUID channel without LC elements in series with the TES. Without the comb inductors, $\tau_e$ is based only on the squid coupling inductance and $R_{TES}$. Since $L_{squid} \ll L_{comb}$, the electrical bandwidth is significantly wider in the configuration we used than in a fMuxed system. As a result, we can probe the detector responsivity well beyond the Mux bandwidth, as shown in Figure 3.11, and verify that the bling and TES remain coupled until the fMux readout bandwidth suppression takes over the responsivity.

A second element affecting the stability requirement for the 90 GHz detectors is the value of $\alpha$, which is a function of the TES bilayer architecture. The basic MoAu bilayer architecture with a geometry that satisfies the $T_C$ and $R_N$ requirements results in an $\alpha$ that causes excessive loopgain throughout the transition. To fix this, we add superconducting structures on top of the TES bilayer, which alter the shape of the transition without shifting $T_C$ or $R_N$. The superconducting structures interact with the TES bilayer to widen the transition and thus reduce $\alpha$ such that we can reliably achieve a low enough loopgain to ensure stability.
Measurements of detector electro-thermal feedback (ETF), as a function of depth in the TES transition. The data shown here was taken in a system without resonant LC elements in series with TES. As a result $\tau_e$ is much less than it is in an fMux comb, allowing us to probe the detector response to a much higher bandwidth. The flattening of the response curve above $\sim 5$ kHz comes from the decoupling of the TES and the bling as the internal $G'$ between them is insufficient for high-frequency fluctuations to drive both the TES and the bling together. As a result, the effective $\tau_e$ of the TES reduces to its own internal $C'/G$, which is much smaller than the time constant of the coupled bling-TES system.

We analyzed a series of patterns for the $\alpha$-reducing Nb structures, including a “finger”, or single 11 $\mu$m-wide bar that completely bisected the TES layer halfway between the electrical leads, a series of 11 $\mu$m-wide bars laid in parallel and each only partially spanning the TES, and a series of 3$\mu$m-diameter “dots” spaced at a pitch of 11 $\mu$m. All three patterns reduce $\alpha$ with respect to a bare TES reference detector, as shown in the $R(T)$ curves in Figure 3.12, but the greatest effect was seen with the dots. As a result, the deployed SPTpol 90 GHz detector TES architecture includes the dots, which can be seen in the microscope image of a TES in Figure 2.7.
Figure 3.12: $\alpha$-Reducing Niobium Structures

$R(T)$ curves for several different patterns of superconducting Nb structures deposited directly onto the TES. The reference TES had no such structures, and exhibits a sharp drop from $R_n$ into the transition. Adding the Nb structures both widens the superconducting transition and extends the regime of a significant slope in $R(T)$ to a higher temperature. The resulting high-temperature responsivity is useful for making measurements when background loading would otherwise drive the TES completely normal.

3.2.4 Detector Noise

Ensuring background-limited detector noise performance is crucial for maximizing instrumental sensitivity. We establish a baseline for the background loading based on the expected optical loading of the deployed instrument as shown in Table 3.1, and characterize the detector noise to ensure non-photon noise sources are in fact sub-dominant. Mounting test detectors without optical feeds and enclosing them inside a sub-Kelvin enclosure reduces the integrated optical loading, even assuming perfect optical efficiency, to $\ll 1\ \text{pW}$, while the target $P_0$ for SPTpol detectors is between 20-40 pW. We then measure the detector noise spectrum at a variety of
operating points to help characterize the different components, as they differ in their
dependence on $R_{TES}$ and $\mathcal{L}$. Readout noise has a constant current spectral density
as a function of frequency, which is then modified by the responsivity when the
noise is referred to NEP. The spectrum of Johnson noise NEP for the TES resistance
changes as increasing loopgain suppresses the noise fluctuations in the TES. Thermal
fluctuation NEP has flat spectrum in power on the detector as it arises from random
fluctuations in the energy transport from the TES.

An example of noise data gathered for a prototype 90 GHz detector is shown
in Figure 3.13. We estimate the expected noise levels from the sources described
in Section 3.1.5, excepting the photon loading, and compare them to the expected
photon noise at the South Pole. In this test configuration, the detectors show noise
spectra consistent with background-limited performance.

3.3 Deployment and Characterization of the Receiver

In November 2011, I made my first trip to the South Pole as part of the early crew
tasked with decommissioning the SPT-SZ receiver and preparing the telescope for
the new SPTpol receiver, which was assembled and installed between December and
January. Our primary job was to replace the cold optical baffles ("cones") in the
optics cryostat with a set that had been modified at Case Western Reserve to hold the
prime-focus IR blocking filters tilted at an angle of 25° relative to normal incidence.
Experience with SPT-SZ had shown considerable reflections between the filters and
structures between the horns on the focal plane, leaking into the detector beams as
sidelobe pick up. Tilting the filters ensures that reflected rays terminate on the cold
baffles instead of coupling back into the optical path, reducing the fraction of the
beam that sees reflected power by a factor of 10.
Noise power spectra for a test 90 GHz pixel, taken at several operating points in the TES transition, with a 4K optical load. Here the noise is referred to the current at the squid and then multiplied by the voltage bias to get units of power, eliminating the spectral dependence of the responsivity that factors into the full NEP calculation. In this form, the thermal fluctuation noise varies as $L / L + 1$. Since we are primarily interested in the white noise level here, we can neglect the effect of responsivity.

A secondary task for the early season crew was the installation of a mount to hold a new external shield near prime focus that I designed, known as the “snout”. The snout is a blackened, conical shield that opens from the Zotefoam environmental window towards the primary, surrounding the beam as it enters the secondary cryostat. It was installed to mitigate coupling to the galaxy and ground as a result of scattering.
of rays off of the environmental window and prime focus filters and into the beam.

Following successful assembly of the SPTpol receiver on site, the collaboration embarked on a period of characterization and calibration of the deployed instrument. A subset of the tests described in Section 3.2 were carried out to validate detector performance in the receiver cryostat prior to installing it onto the telescope. As part of the first season of observations, we also performed several extensive calibration measurements to complete our model of the focal plane for use in subsequent data analysis. In the following sections, I will describe some of the key efforts in characterizing the on-sky performance of the instrument.

3.3.1 Polarization Calibration

As we will see in 4.3, the value of a single CMB map pixel is the averaged time-ordered-data from many different detectors. Each sample from detector $i$, trained on the $\hat{n}$ direction in the sky, describes its instantaneous response to incoming Stokes $I$, $Q$, and $U$ by:

$$d_i(\hat{n}) = G \left( (1 - \rho_i^2)I + \frac{\rho_i^2}{2} (Q \cos(2\psi_i) + U \sin(2\psi_i)) \right), \quad (3.54)$$

where $\rho$ is the polarization efficiency parameter and $\psi$ is the angle of the detector’s primary sensitivity axis relative to a to the coordinate system that defines $Q$ and $U$ [39]. Polarization efficiency, $\rho$, follows the convention described in Equation 3.50, where it describes the ratio of cross-polar to co-polar coupling, with no such coupling corresponding to perfect polarization efficiency.

In SPTpol, we establish the purely elevation direction as $0^\circ$, corresponding to the $+Q$ axis in the Stokes basis illustrated in Figure 1.1, and measure $\psi$ as positive in the counter-clockwise direction looking outward at the sky. Each map pixel value for map $p \in [T, Q, U]$ is then $P(\hat{n})_p = \sum_i d_i(\hat{n})_p$, where the detector response is decomposed
into its components in the T, Q, and U basis. In order to recreate maps accurately, we need a measurement of \( \rho \) and \( \psi \) for each detector to decompose time-ordered-data according to Equation 3.54. An error in the measurement of \( \rho \) reduces sensitivity to polarization but maintains its sense on the sky. However, a rotation of the actual detector angle relative to the value used in 3.54, \( \delta \psi = \psi_{\text{actual}} - \psi_{\text{nom}} \), mixes sky Q and U into measured \( \tilde{Q} \) and \( \tilde{U} \) by
\[
\tilde{Q} \pm i\tilde{U} = e^{\pm 2i\delta \psi} (Q \pm iU). \tag{3.55}
\]
The mixing between Q and U leads to a mixing between E and B, such that the eventual reconstructed power spectra \( \tilde{C}_{\ell}^{BX} \) that include a B-mode component are contaminated by
\[
\begin{align*}
\tilde{C}_{\ell}^{EB} &= \frac{1}{2} \sin(4\delta \psi) (C_{\ell}^{BB} - C_{\ell}^{EE}) \\
\tilde{C}_{\ell}^{TB} &= -\sin(2\delta \psi)C_{\ell}^{TE} \\
\tilde{C}_{\ell}^{BB} &= \cos^2(2\delta \psi)C_{\ell}^{BB} + \sin^2(2\delta \psi)C_{\ell}^{EE},
\end{align*}
\tag{3.56}
\]
where the spectra on the right are the CMB source spectra and those on the left are what is measured. Cosmological \( EB \) and \( TB \) are expected to be zero and can be used to “self-calibrate” for a global rotation of the focal plane [25]. Recall that power in \( EE \) is expected to be roughly two orders of magnitude larger than \( BB \) over the \( \ell \) range of interest, so the \( EE \to BB \) leakage spectrum suggests that in order to make an un-contaminated measurement of \( C_{\ell}^{BB} \) without any self-calibration, \( \sin^2(\delta \psi) < 0.01 \). Assuming the total error \( \delta \psi \) is simply the error on the mean of individual detector angle errors, \( \delta \psi_i \), we set a target for the polarization calibration to measure each detector angle to statistical precision of 2°.

An ideal calibration source for measuring the on-sky polarization response of each detector would be in the telescope far field and at an elevation comparable to our CMB fields. This would provide a direct measurement of the transfer function between a celestial source polarization and detector response with the array in its nominal...
observing state. Unfortunately, the far field of SPT with the full primary illuminated is 100 km away at 150 GHz. Construction and deployment of a far-field source, particularly one at observation-field elevations, is logistically prohibitive. A more technically feasible source was constructed at the South Pole and used to make the necessary measurements described in the remainder of this section.

The polarization calibration source is embedded in a 10m-high flat reflecting panel set at angle $30^\circ$ relative to vertical, that redirects the portion of the telescope beam that does not couple into the source to an elevation of $60^\circ$, reducing loading from the horizon and near-horizon atmosphere. Set 3 km from the telescope, it is close enough for easy access but far enough that the telescope can focus on it. The calibration signal comes from an unpolarized thermal load, chopped at 7.5 Hz and collimated by an HDPE lens, shining first through a stationary wire grid polarizer to establish a clear baseline polarization, then through second wire grid polarizer on a rotating stage that modulates the output polarization. The full aperture of the source is 4 inches in diameter, but a reflecting plate with a 1.5 inch hole is placed over the aperture to reduce loading on the detectors and keep them linear.

The plane of the rotating polarizer is tilted $2^\circ$ from vertical, while the fixed grid is the vertical plane, such that the axis about which the rotating grid turns points $2^\circ$ in elevation, rather than directly back at SPT. This ensures that any signals reflected from SPT are directed back to the atmosphere, so the polarization state orthogonal to the calibration source emission is sourced by the atmosphere above the telescope, rather than the ground. To further suppress reflections, a wooden fence is placed halfway between the source and the telescope, so that rays from the source cannot reflect off the snow and couple back into the telescope beam.

An Arduino Mega 2650 micro-controller board mounted inside the calibration source is used to control the rotating grid angle and record the chopper phase. Timestamps from an Adafruit Ultimate GPS receiver are sampled at the same 40 Hz rate
used by the Arduino to record source data. The recorded source data is transmitted to the telescope control room over a XBEE wireless adapter, where it is logged to a local computer. In addition to recording data, the Arduino runs an embedded control program to step the polarizing grid from $0^\circ$ to $165^\circ$ (in $15^\circ$ increments) and then back down to $0^\circ$ (corresponding to Stokes $+Q$ in the basis defined on the sky from the telescope). The grid is held fixed at each angle for 3 s, and it takes $\sim 1$ s to step between them.

A preliminary polarization calibration run in February of 2012 yielded no useful measurements of detector angles, but revealed several systematic issues inherent in the process. The primary issue was that properly aligning the beam of the pixel under test to the calibration source is quite difficult. The basic pointing model for each pixel assumes static angular offsets relative to the telescope boresight, but we found that the pointing drifts by large fractions of the beam width on hour time scales. The exact reasons for the pointing variation are not clear, as the circumstances of the measurement – observing a source in fixed ground coordinates, near the horizon, with the sun up – are quite different from typical CMB field observations for which the basic pointing model is defined.

In addition to the pointing errors, we discovered that observing the full source aperture overloaded most detectors. Finally, because there are no wires running from the source to the telescope, it was originally operated autonomously, with someone manually starting/stopping the polarizer rotation program before/after a suite of measurements. Arduino data was logged to a memory card that retrieved after measurements were completed. Because the rotation program ran continuously, detectors had to be tuned while viewing a modulated optical load, causing problems with finding an appropriate detector bias point.

In November 2012, I deployed as part of the team charged with using the lessons of the preliminary run to make an improvement polarization calibration measurement.
We derived a scheme to address the pointing difficulties where the beam of the pixel under test is pointed near the reflector based on the naive pointing model, then scanned across it, first from left to right in azimuth and then from the top down in elevation. The change in detector loading as the beam crosses the edge of the reflector is sharp enough to derive precise locations of the reflector edges. Knowing the instantaneous pointing coordinates to three sides of the reflector allows us to derive the coordinates of the source and accurately align the pixel under test to it. In addition to adding the 1.5 inch-diameter aperture plate to reduce the source flux, we used the communication capabilities afforded by the wifi link to change the polarizer rotation cadence. During detector tuning, the rotating polarizer is set to its $90^\circ$ alignment, blocking the modulated source signal and providing a stable background loading. After tuning, a single cycle of the rotation program is commanded, stepping from $0^\circ$-165$^\circ$ and back.

Analyzing polarization calibration measurements begins by up-sampling the rotating polarizer position and the chopper phase data to the same 190.73 Hz sample rate as the detector data. We use the GPS timestamps recorded both by the Arduino and the telescope DAQ system to align the two data sets. We then generate a sinusoidal mixer signal at a frequency determined by the chopper phase data and use it to lock in to the detector response at each discrete angle. We calculate the response as the mean of the per-chop-cycle amplitudes and the error as their standard deviation. With these values and errors, we perform a Markov-Chain Monte Carlo (MCMC) fit of the data to the model described in the next paragraph, from which we recover both the best-fit values and associated errors for $\rho$ and $\psi$ for each detector.

The calibration source output is linearly polarized by each of the two wire grids each with less than 0.1% cross-polar transmission in power. The fixed polarizer ensures a well-known reference polarization state that is then modulated by the second, rotating grid. Reflections between the two grids produce an angle-dependent
MCMC-derived fits of Equation 3.60 to processed polarization calibration data for a single detector. The rotating grid program steps from 0°-165° and back, so there are two samples for each grid angle. The errors come from the scatter on the per-chop response amplitudes. The best-fit polarization angle is marked as the vertical line.

output that is polarized in the transmission direction of the rotating grid, θ_g, with an intensity of:

\[ I(\theta_g) = \frac{2\cos^2(\theta_g)}{(1 + \cos^2(\theta_g))}. \]

Decomposing the calibrator source signal at an arbitrary θ_g into the Stokes basis, 3.57 becomes:

\[ S(\theta_g) = I(\theta_g) \begin{pmatrix} 1 \\ \cos(2\theta_g) \\ \sin(2\theta_g) \end{pmatrix} \]

Substituting this output from the calibration source for the Stokes components in Equation 3.54, we can rewrite the detector response as:

\[ d_i = AS(t, DC) \left( (1 - \rho_i^2)I(\theta_g) + \rho_i \left( \cos(2\theta_g) \cos(2\psi_i) + \sin(2\theta_g)\sin(2\psi_i) \right) \right), \]

or equivalently,

\[ d_i = AS(t, DC)I(\theta_g) \left( \cos^2(\theta_g - \psi_i) + (1 - \rho_i)\sin^2(\theta_g - \psi_i) \right), \]
where $A$ is the amplitude of the modulation in detector response as a function of rotating grid angle, $\theta_g$, $\rho$ and $\psi$ are the polarization efficiency and detector angle we are interested in. Here we have converted from a general Stokes basis to one defined by a particular detector response, so the $(\cos^2(\theta_g - \psi_i))$ describes the coupling to co-polar radiation while the $(1 - \rho_i)\sin^2(\theta_g - \psi_i)$ term describes the cross-polar coupling. The $S(t, DC)$ term is a gain scaling factor that can be allowed to vary linearly as a function of time or the DC level of the chopped detector response signal. We find that allowing the gain factor to vary either in time or with DC level has a negligible impact on the fits, and generally results in a very small slope. This reaffirms the fact that most detectors are biased at a stable point in their transition and remain there throughout the observation with very little change in their responsivity. Tracking the gain variation as a function of DC level is more likely to reveal nonlinearity, so we opt to include $S(DC)$ in the final analysis.

$$
\begin{array}{|c|c|}
\hline
\delta \psi & < 2^\circ \\
\delta \rho & < 5\% \\
\chi^2 \text{ PTE} & > 10^{-9} \\
\Delta S/S & < 20\% \\
A & < 250 \\
\hline
\end{array}
$$

Table 3.3: Polcal Cut Parameters
Cut parameters for polcal analysis. In addition to keeping only fits with low statistical errors on the recovered $\psi$ and $\rho$ parameters, we eliminate fits that have abnormal secondary statistics. Extremely low PTE values from a $\chi^2$ fit to the data indicate improper error estimation on the data or a poor adherence to the model. A large fractional change in the gain parameter, $S$, indicates either a nonlinear detector or data that is dominated by a drift instead of coupling to the calibrator source signal. Finally, extremely large modulation amplitudes ($A$) require an un-physically large detector gain and indicate a problem with the data. The cut levels on the final three parameters were determined empirically by examining the distributions and limiting outliers.

We derive the final polarization calibration values, $\psi$ and $\rho$ for each detector by gathering all of the fit parameters from all observations, and cutting the results according to the limits in Table 3.3. The primary limits are the statistical errors
\(\delta\psi\) and \(\delta\rho\), which indicate low signal-to-noise data, but further cuts are added to ensure that various forms of un-physical or corrupted data that result in spuriously low error bars are eliminated. Recovered values for each detector are then combined by taking the inverse-variance-weighted mean of the recovered parameters from all observations of that detector. Some detectors, particularly those with a consistently-strong response to the calibration source, were observed multiple times as we used them to determine if our data acquisition and analysis were working properly. Other detectors did not have any good observations, so we assigned them the the median \(\psi\) and \(\rho\) from all good measurements of detectors in the same module with the same nominal angle.

Over the course of two polarization calibration runs, one in November 2012 with the receiver in its first-season state and one in January 2013 after a series of upgrades, we gleaned measurements of \(\rho\) and \(\psi\) for 971/1523 active detectors. The remaining detectors were assigned values based on distributions of similar detectors as described above.

To check for possible systematic contamination of the calibration measurement, we performed a series of calibration observations using a subset of well-behaved detectors with various changes to the measurement configuration. We first checked whether our assumption that the measurement was not strongly sensitive to telescope focus position by intentionally de-focusing the telescope. We then removed the anti-reflection fence to check whether reflections from the source contributed a considerable bias to the measurements. The normal bias routine for a pixel under test involved tuning both detectors simultaneously, so we performed the measurement with only one detector tuned at a time to ensure that no readout crosstalk affected the calibration results. Finally, we removed the fixed polarizing grid to leave only the rotating polarizer modulating the source output polarization angle and altered the fit function accordingly. The results for the mean difference between the angles derived
Figure 3.15: Polarization Calibration Fits

Top: Histogram of the statistical error on the MCMC fit to the rotation angle, $\psi$. Shown here is a portion of the full distribution, which has a tail up to $\sigma_\psi \sim 70^\circ$, but those fits are performed on very low signal-to-noise data, where the detector was not responding to the calibration source. Bottom: Histogram of errors on $\epsilon$, the fractional coupling to cross-polar radiation, rather than the polarization efficiency, $\rho \equiv 1 - \epsilon$. 

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<table>
<thead>
<tr>
<th>Test</th>
<th>$\langle \Delta \psi \rangle$, degrees</th>
<th>$N_{\text{obs}}$</th>
<th>Significance</th>
<th>$\sigma_m$</th>
</tr>
</thead>
<tbody>
<tr>
<td>De-focus</td>
<td>0.78 (0.55)</td>
<td>4 (13)</td>
<td>1.04 (1.32)</td>
<td>0.75 (0.42)</td>
</tr>
<tr>
<td>Fence removed</td>
<td>0.68 (0.55)</td>
<td>5 (15)</td>
<td>1.01 (0.14)</td>
<td>0.67 (0.39)</td>
</tr>
<tr>
<td>Single detector tuned</td>
<td>0.32 (-0.21)</td>
<td>5 (8)</td>
<td>0.48 (-0.40)</td>
<td>0.67 (0.53)</td>
</tr>
<tr>
<td>Fixed grid removed</td>
<td>0.07 (-0.07)</td>
<td>98 (664)</td>
<td>0.46 (-1.20)</td>
<td>0.15 (0.06)</td>
</tr>
<tr>
<td><strong>Total Uncertainty</strong></td>
<td></td>
<td></td>
<td><strong>1.22 (0.78)</strong></td>
<td></td>
</tr>
</tbody>
</table>

Table 3.4: Polarization Calibration Systematic Errors

Systematic error accounting for 90 (150) GHz detectors during the polarization calibration process. For each detector observation in a systematic test, the best-fit angle is compared to the aggregate best-fit angle derived from nominal observations, and the mean of all such angle differences is shown as $\langle \Delta \psi \rangle$. The number of observations used to calculate $\langle \Delta \psi \rangle$ for each test are shown as $N_{\text{obs}}$. The statistical error on each fit is taken to be $1.5^\circ$ and the error-on-the-mean, $\sigma_m$ of each measurement is $1.5^\circ/\sqrt{N_{\text{obs}}}$. Significance is calculated as the factor $\sigma_m$ each average angle difference is from zero. The final systematic uncertainty in the mean angle is then taken as the quadrature sum of the $\sigma_m$ values for each test.

in the nominal configuration and during each systematic test are shown in Table 3.4.

During the polarization calibration analysis, we realized that due to an ambiguity in the mounting scheme, many 90 GHz detectors were installed with polarization angles rotated $90^\circ$ to the nominal value; essentially the X and Y detectors within some pixels had been swapped. Using the calibration data and resulting fits, we were able to determine which detectors were reversed and correct the appropriate configuration files. As long as one detector from each pixel had an unambiguous fit – meaning even if it failed the cuts listed in Table 3.3, the fit was good enough to establish which of the two possible nominal angles best fit the data – the appropriate median value could be assigned.
Chapter 4

COLLECTING AND REDUCING DATA

There are three main sources of data recorded during normal telescope operations: bolometer time-ordered data (TOD or “timestreams”), output from telescope motion encoders, and thermometry. All data is collected by the Generic Control Program (GCP), which is essentially the operating system of the telescope, performing the interface for control, monitoring, and data acquisition [49]. GCP combines the various data streams with a consistent time base derived from a 100 Hz IRIG-B signal generated on a dedicated timing card that is synchronized to GPS time. Data is stored locally at the South Pole in “arcfiles”, which are a series of registers describing individual data frames, stored in fixed-length binary files. Each day, a copy of the approximately 60 GB of new arcfile data is transmitted via satellite to computers at the University of Chicago, where the processing and analysis happens.

The SPTpol analysis pipeline is written largely in Python 2.7, with elements of C and IDL. The primary interface with the data is a so-called “dataReader” object, which takes a start and end time as inputs, reads the binary data from the associated range of arcfiles, and arranges it into a record structure that allows for straightforward
access and manipulation.

4.1 CMB Field Observations

4.1.1 Field Depth

A common parametrization of the noise in a CMB field map with an area of $A_{\text{sky}}$ is the so-called “depth”, defined as the standard error on the mean measurement of temperature in a given map pixel. For a single detector observing a single spot on the sky, the variance on a measurement of sky temperature per second of integration time is just the $\text{NET}^2$ from Equation 3.44. Since the NET is essentially a white noise level, more observing time reduces the error on the recovered sky temperature by $\sqrt{t_{\text{obs}}}^{-1}$.

The definition of NET includes an implicit integration over the angular response of the detector, commonly characterized by the FWHM of the radially-averaged beam $\sigma_b$, so that each data sample is measuring a spot $\approx \sigma_b^2$ in size.

The natural map pixelization is the number of independent measurements of the field that can be made by an instrument with a given angular resolution, or $n_{\text{pix}} \sim A_{\text{sky}}/A_{\text{pix}}$. For $n_{\text{det}}$ detectors in the array, the time spent observing each map pixel is $t = n_{\text{det}}t_{\text{obs}}/n_{\text{pix}}$, where $t_{\text{obs}}$ is the total observing time. The simple analytical estimate for pixel variance, using assuming uniform detector NETs and beam widths is then:

$$
\sigma_T^2 = \frac{A_{\text{sky}} \text{NET}_{\text{det}}^2}{\sigma_b^2 n_{\text{det}}t_{\text{obs}}} \quad [\mu K^2/\text{pixel}].
$$

(4.1)

Clearly for a fixed amount of observing time, deeper maps are produced by observing smaller fields with more detectors. Given that the pixel number is fixed for a given instrument, the remaining free parameter in selecting CMB fields is their size. Several factors serve to counteract the advantages of smaller, deeper fields. Assumed in Equation 4.1 is uniform coverage of the map by all detectors. In reality, the finite
field-of-view of the telescope means that the region of uniformly-sampled depth is an area inset from the border of the field by field-of-view width.

A more important consideration in the determination of field size is the effect of sample variance on estimates of the CMB power spectrum. Mapped over the whole sky and decomposed into the spherical harmonic basis $Y_{\ell m}(\theta, \phi)$ with amplitudes $a_{\ell m}$, the CMB angular power spectrum $C_\ell = \frac{1}{2\ell+1} \sum_{m=-\ell}^{\ell} |a_{\ell m}|^2$ is the measure of power in fluctuations with angular scale indexed by $\ell$. The component $a_{\ell m}s$ are Gaussian random numbers distributed about zero with a variance $C_\ell$, and there are $2\ell+1$ $a_{\ell m}s$ in each $C_\ell$, so the statistical variance on each estimate of $C_\ell$ is $\sigma^2_{C_\ell} = \frac{2}{2\ell+1} C_\ell$. In the case of observing only some fraction $f_{\text{sky}}$ of the full sky, the number of samples contributing to each $C_\ell$ measurement is reduced, so the effective sample variance is $\sim \frac{2}{f_{\text{sky}}(2\ell+1)} C_\ell$.

4.1.2 Observation Strategy

An altazimuthal bearing at the South Pole is also an equatorial mount, since telescope motion about its elevation (el) axis traces out arcs along longitudinal great circles and motion about the azimuth (az) axis trace arcs of constant latitude. Telescope coordinates are in the az, el of the bearing, with zero azimuth along the prime meridian and zero elevation at the horizon. Celestial coordinates account for the motion of the earth, so there is a time-varying relationship between the local telescope $\phi_{\text{az}}$ and the celestial $\alpha_{\text{ra}}$. Thus, the mapping between telescope $(\phi_{\text{az}}, \theta_{\text{el}})$ and celestial $(\alpha_{\text{ra}}, \delta_{\text{dec}})$ coordinates is:

$$\Delta \phi_{\text{az}} = \Delta \alpha_{\text{ra}}, \quad (4.2)$$

$$\theta_{\text{el}} = -\delta_{\text{dec}}. \quad (4.3)$$

The actual angle subtended on the sky for a change in pure elevation is just
\[ \Delta \psi_{\text{dec, sky}} = \Delta \theta_{\text{el}}, \] while sweeps in azimuth subtend an elevation-dependent angle \[ \Delta \psi_{\text{ra, sky}} = \Delta \phi_{\text{az}} \cos \theta_{\text{el}}. \] This is important because it means that scans at a constant velocity in azimuth map to different sampling rates of sky pixels depending on the elevation of the scan. In other words, the sky is sampled more finely at a lower declination (higher el) for a fixed field width \( \Delta \delta_{\text{ra}} \) in celestial coordinates.

The scan pattern for observations with SPT is a series of back-and-forth sweeps in the azimuth (RA) direction across the width of the field with a small increase in elevation after each such scan. Thus, each distinct elevation position has both a “left-going” and “right-going” scan associated with it. This is useful for determining noise by differencing maps made from only left-going scans and those with only right-going scans. Each pair of scans at a given elevation have the same detector pointings, so they sample the same fixed celestial signal, while both the white noise and (nominally) the atmospheric fluctuations between them are uncorrelated, so differenced maps should contain only residual, uncorrelated noise.

Pickup from stationary ground sources will also behave largely as a stationary correlated sky signal that will remain as a systematic residual in summed maps. As a precaution against this we observed during the first two seasons in “lead-trail” mode, where the two side-by-side halves of the field are each mapped independently. The first half of the field is observed as a particular section of horizon rotates beneath it. When the first half of the observation is complete, the horizon has rotated relative to the sky such that the second half of the field is observed over the same section of the horizon as the first. To the extent that it is fixed in azimuth, any ground pickup will be present identically in each half of the map, while all other signals (including the sky) are uncorrelated. We can then measure and remove a correlated ground-synchronous template from both halves of the map. Otherwise, the two half-fields can simply be combined and the resulting map treated as a single observation.

Individual pixel beams are separated on the sky, such that the field of view image
is under-sampled at each instant. The scan strategy described above naturally fully samples the field in \( \text{ra} \), but in order to cover the whole field in a minimal amount of time (or, in the case of lead-trail observing, the amount of time set by the rotation of the ground under half the field), the most efficient el step for fully covering the field in elevation in a minimal amount of time is a factor of ten larger than the detector beam FWHM, at roughly 10 arcminutes. The focal plane is mounted with an \( 8^\circ \) rotation of its vertical axis of symmetry to ensure that the elevations of each pixel beam are staggered relative to each other to improve per-scan coverage. We also subdivide the full el step between scans into 12 or 18 (depending on the field) “dither steps”, which are slight offsets in elevation applied to the field scan profile. We cycle through the dither for each observation, such that sequential field observations are offset from each other in elevation by one dither step, wrapping back to the original scan schedule after cycling through each set of steps. This ensures complete sampling of the field in elevation by each detector over the course of each dither cycle.

Power spectrum analyses with SPTpol are performed on two fields, the so-called “deep field” observed over the 2012 observing season and for three months early in 2013, and the “survey field”, which will be observed for the remaining three years of SPTpol operations. Figure 4.1 shows the location of SPTpol fields, imposed on a map of galactic dust emission temperature measured by the Planck satellite. Table 4.1 shows the depths of the main SPTpol fields both currently and projected for the remainder of the experiment.

We observed the deep field for the 2012 season in order to get as deep a field as possible in a single year of observations, particularly in the \( \ell \) range relevant to lensing B-modes. For longer-term analyses, the survey field has a lower cosmic variance term owing to the greater number of modes measured and the relaxed filtering requirements for removing 1/f noise from the data, improving the sensitivity to gravitational wave B-mode spectrum. In addition to the primary fields used for polarized power spectrum
The deep field was observed throughout 2012 and for the first two months of 2013. A typical winter observing season lasts 8 months, during which time approximately 65% of the time is spent observing CMB fields. Analysis, we spend the summer months observing fields that are not contaminated by the proximity of the sun. These summer fields are not deep enough to contribute to polarization analyses, but are deep enough in temperature to be used for CMB temperature science.

Figure 4.1: SPTpol Fields
The location of SPTpol observation fields, in equatorial celestial coordinates, plotted over the Planck temperature map of composite dust emission [2]. The fields are chosen both to minimize the effect of galactic emission, which is brightest along the galactic plane (visible as the bright streak along the top of the projected image), and to overlap with other astrophysical observations.
A typical observing cycle for the survey field in 2013. Single horizontal lines indicate el nod, cal stare combinations, which are performed before any other observation. Multiple horizontal lines indicate fast scans of mat5a and RCW-38 along with attendant cal stares and el nods. The bracketed numbers noted for the survey field observations indicate which dither step (out of 18) the observations used. Details of each auxiliary observation type are described in Section 4.1.3.

Each “observing cycle” is defined as the amount of time between refrigerator cycles. Each fridge cycle lasts approximately six hours, after which the focal plane remains at \( \sim 270 \) mK for 28 hours on average. Most of the 28 hours of “hold time” is devoted to field observations, but a series of auxiliary observation, described in Section 4.1.3, are also performed. A typical schedule for a single cycle is shown in Table 4.2.

### 4.1.3 Auxiliary Observations

Each observing day includes a series of short observations that are used for monitoring the observing conditions and telescope performance. Immediately prior to each field observation, the telescope is held still at the mid-field elevation, and a shutter in the secondary is opened for one minute to measure each detector’s responsivity to a 6 Hz
chopped signal from a stable thermal calibration source. These “calibrator stares” or simply “cal stares” are used both to find detectors with low responsivity and to correct for day-to-day detector gain variations, as described below in Section 4.2.4. At least once per observing cycle we record calibrator stares at a series of six elevations within the approximate field elevation range (“cal, el noe, noise vs. el”), in order to measure any change in the response to a 6 Hz optical signal as the DC optical loading changes due to variations in atmospheric opacity with viewing elevation. Finally, once per observation cycle we run a “calibrator sweep”, where a series of calibrator stares are done sequentially with increasing chop frequencies between 5-75 Hz to monitor detector time constants.

Along with each calibrator stare we perform an “el nod”, where the telescope sweeps down two degrees in elevation at a constant velocity, then back up to its initial position. The signal induced in the detectors by the change in atmospheric opacity with elevation is then fit to a line, the slope of which reflects the gain of the detector to a change in atmospheric column air mass, which can be modeled for high elevations as:

\[ M_a = M_{a,\text{zen}} \frac{1}{\sin\theta_{el}}, \]

(4.4)

where \( M_{a,\text{zen}} \) is the zenith airmass. For a fixed range of \( \theta_{el} \), changes in detector response \( S(M_a) \) can be the result of individual detector tunings, but averaged over the array, we can measure variations in \( S(M_{a,\text{zen}}) \) between observations.

Detector noise is measured during the “cal, el nod, noise vs. el” schedule, during which a 2 minute calibrator stare, el nod, and 2 minute noise stare are performed at five elevations spaced evenly between the top and bottom elevations of the CMB field. These noise stares are useful for measuring baseline detector noise performance for comparison with field observations. Variations in elevation induce different levels of static background loading, which allows us to probe noise performance as both a function of optical load and, when the detectors are not re-tuned with elevation, as a
function of depth in the transition. At start of each fridge cycle, we take a ten minute
noise stare pointed at the field center elevation to build an archive of high-resolution
detector noise spectra.

The other class of regular auxiliary observations are maps of bright astrophysical
sources near our CMB fields. We make maps of the galaxy Centaurus A, NGC3576
(MAT5a) nebula, and the RCW38 star cluster for use in reconstructing the exact
telescope pointing over time, as described in the next section. We employ two different
schedules for observing these sources, one with a fine elevation spacing between az
scans to build deeper maps, the other at a coarser elevation spacing to avoid diverting
time from CMB field observations. The HII regions MAT5a and RCW38 were used
as calibration sources in SPT-SZ along with several other CMB experiments, as their
properties have been well-studied and cataloged [8]. In the next chapter I will describe
the various processing routines that make use of these auxiliary observations.

4.2 Data Pre-Processing

Detector data is stored in raw DfMux ADC units, along with the meta data necessary
to convert it into physical TES units. Along with detector data, a wide variety of
auxiliary information about the telescope motion and the observing environment is
recorded. As our understanding of the data has evolved, a suite of pre-processing
routines that condition and transform the data into a more useful state has been
established. In the following sections, I will describe some of the basic operations
performed on raw data before CMB mapmaking. A large part of my work since
deployment has been in developing and improving these routines.
4.2.1 Glitch finding

Over the course of an observation, detector timestream data will occasionally include spurious events unrelated to the desired sky signal that must be identified and removed from inclusion in sky maps. We flag these glitches as the first step in the data processing chain to prevent them from being altered by subsequent filter steps. Data is read in at full resolution with no pre-processing, the glitch finding described below is then run, and flags indicating which detectors had glitches for each left/right scan are stored.

We also analyze the RMS of each scan and store the flags indicating high-noise scans. Because detector timestreams can have considerable low-frequency drifts from atmospheric fluctuations, it is only sensible to measure RMS on filtered data. After glitch finding, per-scan timestreams have a fourth-order polynomial removed before the RMS is calculated. We do not wish to retain this filtering in the raw data, so only the RMS values for each filtered scan are saved. With the results of the glitch finding and RMS calculation stored, the data is re-read with the normal pre-processing routine, and the flags applied to the new dataReader object.

Glitch events fall into two broad categories: standard glitches and DC jumps. Standard glitches are characterized by large spikes in the timestream that happen over usually less than 10 samples, after which the detector returns to its previous state and behaves normally. Based on the sign of the detector response, standard glitches are assumed to be cosmic rays depositing power directly onto the TES, resulting in a reported jump in temperature over a shorter time scale than the beam extent and scanning speed would permit for even the brightest astrophysical sources. Furthermore, sky signals are correlated across nearby detectors in a predictable way, whereas standard glitches typically only show up in a single detector; concurrent glitches found in more than one detector are rare, and generally attributable to one detector responding to a glitch on the other through electrical crosstalk.
In the case of DC jump glitches, the same single-sample excursion from the mean sample-to-sample variance occurs, but instead of quickly returning to its initial state, the detector DC level remains at the peak level of the excursion. Also unlike standard glitches, these DC jumps are often seen in other detectors on the same SQUID, with a lower amplitude.

Standard glitches can have amplitudes several orders of magnitude greater than the detector timestream RMS, resulting in a spurious high-temperature measurement of the sky pixel being observed during the glitch. In addition to biasing random map pixels towards high temperatures, glitches present problems for other data processing steps in the pipeline. Timestream filtering is performed on scan-length sections of the data, which are \( \approx 1000 \) samples long after downsampling. Some large glitches have amplitudes high enough to significantly bias the fit to a polynomial that is then subtracted as a high-pass filter, distorting data from the entire scan and introducing a larger effect on the final map than a biased measurement of a single pixel temperature.

Identifying and flagging glitches without including false positives from reasonable statistical noise fluctuations in a way that is robust to variations in the overall timestream profile is rather nuanced. Expected low-frequency noise from atmosphere and drifts in the refrigerator base temperature cause the DC level of a typical detector timestream to drift considerably with respect to the high-frequency noise level over the course of a single scan. Furthermore, glitch finding requires examining hundreds of detector timestream segments for hundreds of scans across the field, so minimizing the run time of the glitch finding algorithm is crucial. Fortunately, the sample-to-sample difference in a normal timestream is effectively a high-pass filter on the timestream data, producing a nearly Gaussian distribution of values from which glitches can be clearly identified as outliers.

In order to decrease sensitivity to expected hash on the timestreams, we take the difference between five leading and five trailing samples at each point in the
Figure 4.2: Timestream Glitches

*Top:* Detector scans with standard glitches. The vertical spacing between timestreams is inserted for clarity. Glitches happen over the span of a few samples, with amplitude excursions well in excess of typical astrophysical sources. *Bottom:* A DC-jump glitch shown for multiple detectors on the same SQUID comb. The amplitude of the glitch is reduced the farther in frequency a detector is from the one with the largest jump, consistent with the reduced crosstalk impedance between them.
timestream, to build up a “difference timestream” for each detector scan. The mean and standard deviation of the difference timestream values is then computed, and samples that lie more than six standard deviations from the mean difference are flagged. Finally, for each flagged glitch, twenty samples on either side are averaged to determine whether the glitch was “standard” or a DC jump. In the event of a DC jump, the data for all other detectors on the same SQUID comb is also flagged for that scan as a precautionary measure.

Glitch finding is performed on a scan-by-scan basis, and the detectors seeing a glitch during a given scan are flagged so that when observation maps are made, data from flagged detector scans are not used. If a glitch is found in more than 10% of scans for a given detector during a given observation, that detector is flagged as bad for the entire observation and its data is not used. In a given observation, approximately 1% of all total scans are flagged, and it is rare for a detector to meet the 10% of scans threshold and be cut for the entire observation. The brightest point sources in the observing field can induce a response that is flagged as a glitch, which is problematic since the location of the source is constant in map space, so if a large number of detectors are flagged for encountering a glitch when they in fact scanned the source, a strip of poor coverage will be present in the final map. To avoid this, we use a catalog of known point source locations in concert with basic detector pointing information to exclude glitches near sources.

Scans corrupted by glitches or high RMS data are recorded as meta-data fields in the dataReader object. The data itself is not altered, but subsequent analysis steps like timestream filtering and mapmaking exclude flagged scans. In the normal processing of an observation, the glitch and RMS flags are copied from the dataReader and a new dataReader is created for the same observation, using the processing steps described in the following sections. The glitch and RMS scan information is then copied back into the new dataReader, preserving the information about which scans
were corrupted independently of any processing that might have changed the performance of the glitch finding and RMS calculations.

4.2.2 Downsampling

Once the flags for each detector indicating which scans had glitches or high RMS are generated, the dataReader object for an observation is re-loaded, but down-sampled during read out. The default sample rate of the timestream data is 190.74 Hz, which comes from the dFmux board clock rate (25 MHz) divided by a digital downsampling factor of $2^{17}$. For a 1 deg/s scan rate in telescope az at elevation of 55°, the Nyquist frequency of the timestream data corresponds to spatial fluctuations corresponding to an $\ell_{\text{max}} \approx 60000$, which is far in excess of the multipole range of interest for CMB science and considerably smaller than the detector beam size. Downsampling the data by a factor of 4 keeps $\ell_{\text{max}}$ above 10000, which is still above interesting CMB scales but fine enough to avoid affecting interesting data, while considerably reducing data storage requirements and readout times. Prior to actually downsampling the stored data, a FIR low pass filter slightly below the downsampled Nyquist frequency is applied to ensure that aliased high frequency noise is not retained in the downsampled data.

4.2.3 Notch Filtering

Microphonic pickup from the pulse tube coolers attached to the receiver and secondary cryostats results in spectral lines that are coherent in time across the array, with a large enough amplitude that they average down more slowly than other, less persistent signals. In an array-average of detector timestream PSDs from a typical CMB observation, the FWHM of a Gaussian fit to the pulse tube line is typically between 1-2 mHz, which is approximately the frequency resolution for the full length of a CMB observation. In order to minimize the effects of the PTC signals on sub-
sequent filtering steps, the notch filtering is performed immediately after reading in the down-sampled detector data, prior to any further processing.

The notch filtering algorithm takes a full-resolution power spectrum of each detector timestream for a given observation and averages them to produce a single, array-averaged spectrum at the full frequency resolution. Individual spectra are calculated after padding the data to a length at least a factor 1.2 longer than the original record length in order to avoid ringing. Beyond that threshold, the exact padded length is chosen to maximize the speed of the FFT algorithm. The array-averaged PSD is then searched in a 0.5 Hz bandwidth around the programmed pulse tube frequency to identify the precise center frequency of the line. Finally, we fit a Gaussian centered on the pulse tube frequency to the surrounding spectrum bins and flag a width of $3\sigma$ on either side of the PTC frequency to be notched. In the event that the fit is unusually wide, we restrict the maximum bandwidth notched for each line to 7 mHz. Each detector spectrum then has the resulting bandwidth zeroed in its complex FFT before inverting the data back into real space.

4.2.4 Timestream calibration

Detector data is stored in archive files as ADC counts at the DfMux output. Using a basic DfMux circuit model in concert with a known $V_{bias}$ for each detector, we can convert ADC counts to a rough estimate of electrical power in Watts. The default detector timestream units in a dataReader are Watts, based on this rough conversion. We relate the data in Watts to celestial brightness temperature $T_{CMB}$ using an astrophysical source at a known temperature.

RCW38 is a well-measured star cluster that is bright in mm-waves and, at coordinates of (RA 08h 59m 19.2s, Dec -47 30’ 22”), sits near the lower-elevation portion of SPTpol CMB fields. As shown in Table 4.2, dedicated observations of RCW38 are interspersed with CMB field observations providing regular measurements of each
Figure 4.3: Effect of Notch Filtering

Averaged detector PSD for the array, before and after notch filtering. The line centers are found in the averaged PSD, and we then fit a Gaussian to the line, removing 3 \( \sigma_{\text{fit}} \) of bandwidth on either side of the center frequency, with a maximum \( \delta f \) of 7 mHz from individual detector FFTs.

detector’s response \( \rho_{\text{RCW38}, i} \) in W. In the most basic form, our calibration to convert detector timestreams, \( TOD_i \) from units of Watts to \( K \) is:

\[
TOD_i[K] = TOD_i[W] \frac{T_{\text{RCW38}}[K]}{\rho_{\text{RCW38}, i}[W]},
\]

(4.5)

where \( T_{\text{RCW38}} \) is a published map of RCW38 in brightness temperature units in observing bands corresponding to those used by SPTpol [48]. Individual detector observations of RCW38, \( \rho_{\text{RCW38}, i} \), are performed at high resolution once per observing cycle, yielding a template of each detector’s response in raw detector power units to the published map.

Internal calibrator stares are performed prior to every sky observation, while
RCW38 observations are scattered throughout each observing cycle. Therefore, we relate every detector calibrator response, $\rho_{\text{cal},i}$, to CMB temperature in order to make finer adjustments to the gain factors. We convert from $T_{\text{RCW38}}$ to $\rho_{\text{cal}}$ using the ratio of season-averaged medians of RCW38 ($\langle \rho_{\text{RCW38},i} \rangle$) and calibrator responses ($\langle \rho_{\text{cal},i} \rangle$) for each detector to relate $T_{\text{RCW38}}$ to the median calibrator response of each detector.

Further corrections account for variations in atmospheric opacity and the calibrator brightness. Atmospheric opacity variations are measured by comparing the module-averaged RCW38 template $\rho_{\text{RCW38, mod}}$ of each RCW38 observation to its season-long median, $\langle \rho_{\text{RCW38, mod}} \rangle$. While the calibrator is expected to be stable over the course of an observing season, variations in its brightness are tracked by comparing a module-averaged response of each individual stare, $\rho_{\text{cal, mod}}$ to the season-long median value $\langle \rho_{\text{cal, mod}} \rangle$. Thus, for each observation, the known RCW38 template is compared to the nearest RCW38 observation, calibrator stare, and their season-long medians by:

$$TOD_i[K] = TOD_i[W] \frac{T_{\text{RCW38}}}{\langle \rho_{\text{RCW38},i} \rangle} \left( \frac{\langle \rho_{\text{cal},i} \rangle}{\rho_{\text{cal},i}} \right) \left( \frac{\langle \rho_{\text{RCW38, mod}} \rangle}{\rho_{\text{RCW38, mod}}} \right) \left( \frac{\rho_{\text{cal, mod}}}{\langle \rho_{\text{cal, mod}} \rangle} \right).$$

(4.6)

4.2.5 Pointing Corrections

Position encoders attached to the azimuth and elevation bearings are sufficient for a rough estimate of the telescope boresight pointing. Observations of RCW38 shortly after deployment were used to measure nominal detector beam pointing offsets relative to boresight. This so-called “online pointing” is sufficient for directing observations of known sources, but is too coarse for use in reconstructing maps with the precision allowed by the instrumental resolution. The “offline pointing” fits for a six-parameter model that accounts for flexing of the telescope boom as a function of elevation, known tilts relative to nominal of both the azimuth and elevation bearings, and changes due to weather conditions and telescope temperature. Periodic observations
of bright HII regions described in Section 4.1.3 are then fed into the offline pointing model to identify the solution that minimizes the spread of the source images due to pointing errors. The offline pointing parameters are measured throughout the course of observations correcting pointing uncertainty to $\approx 10''$ RMS.

4.2.6 Intermediate Data Files

The processing operations described above are computationally intense and require considerable resources. Loading a single observation from arcfile data consumes upwards of 20 GB of RAM and takes upwards of 10 minutes. Pre-processing routines, particularly glitch finding and notch filtering, can each take nearly as long as it does to load the data. Furthermore, arcfiles are contiguous in time, with no inherent indication of where observations start and end. To reduce the overhead of reading and interacting with data, we periodically generate intermediate data files (IDFs) of individual observations.

To make an IDF, we read the arcfile data, perform glitch-finding and per-scan detector RMS calculations, then re-read the data. The second readout downsamples the data and performs notch filtering on it. We typically leave timestreams in the basic electrical power units, converting to temperature units when making maps. Offline pointing is applied during the second readout as well, since it does not effect bolometer data. The resulting dataReader object for each observation is then saved in HDF5 table format.

Loading the IDF of a single field observation takes less than 10 seconds. The only irreversible changes to the data saved in an IDF are the notch filtering and downsampling; calibration and pointing corrections can be reversed since the factors applied to data during the corrections are saved with the IDF. Glitch finding results and RMS calculations are saved as flags that can be ignored or modified.

In addition to its full IDF, each observation also has an associate “stub” IDF,
which is a copy of everything except the actual bolometer TOD. These stubs include all of the metadata associated with each detector, including calibration factors, responses to auxiliary observations, and noise properties. These can be used, for instance, to convert simulated CMB skies into simulated timestreams with characteristics similar to the real data.

4.3 Making CMB Maps

Once raw data has been pre-processed and stored in IDFs, maps with a variety of post-processing options suited to a particular type of analysis can be created. The scan strategy described in Section 4.1.2 includes a series of scans across the field at a constant telescope az velocity, one left-going (decreasing ra) and right-going (increasing ra) for each elevation. We do not use data taken during the turnarounds and el steps of the telescope to maps because of resonant oscillations in the telescope that show up in detector data during the accelerations. The fundamental unit of a field map is then the individual field half scans, and after the appropriate processing is applied to each detector’s per-scan data, they are averaged into pixels based on their pointing information and the desired pixel size. In the following sections, I will describe the process of converting pre-processed detector time-ordered data into CMB field maps.

SPTpol is optimized for the study of small angular scales, and CMB polarization analysis requires deep maps, so typical fields cover less than 1% of the sky area, allowing the use of a flat sky approximation to simplify the analysis [22]. It is of course possible to expand the analysis to a spherical sky with an $f_{sky}$ coverage fraction, but since we have not found the need to do so, this discussion will proceed using the flat sky approximation. In this limit the formalism of spherical harmonic decomposition is replaced with simple 2-dimensional Fourier transforms, and the angular power
spectrum $C_\ell$ is simply averaged annuli corresponding to the magnitude of a particular wave number, $k$, in harmonic space. The approximate mapping between map space and $\ell$ space is $\ell \simeq 2 \pi \kappa |\vec{k}|$, where $\vec{k} = (k_x, k_y)$ is a vector in 2-d Fourier space.

4.3.1 Detector Cuts

The first level of cuts removes detectors that fail tests of overall data quality for the entire observation. During IDF processing, detectors that fail the bias tuning algorithm, are intentionally left un-tuned (in which case they are kept normal with excess bias current to avoid degrading the performance of other detectors on the same comb), or are read out by a squid displaying poor performance get flagged for the entire observation. Next, outliers in the distributions of responses to the nearest calibrator stare and elnod are also removed. Failing one of these two cuts accounts for the majority of detectors removed, as the cuts are applied sequentially and most other forms of undesirable behavior also result in poor response to auxiliary observations. Finally, detectors with a pixel partner that has been cut are excluded to ensure proper reconstruction of polarized maps.

A second level of cutting removes only bad single scans across the field. Glitch finding described in Section 4.2.1 only flags those scans where a detector sees a glitch. Similarly, the RMS of each detector, calculated after removing a fourth-order polynomial, is recorded and compared to its full-observation distribution, and scans with aberrant values of RMS are cut. During the calculation of per-scan cuts, if a detector has more than 10 % of all scans in an observation flagged, the entire observation is cut for that detector.

4.3.2 Filtering

Signals in the native audio frequency space of detector sampling, $f_A$ (in Hz) can be mapped to the spatial scale on the sky, $\ell$, to which they correspond by the approximate
relationship:

\[ \ell \approx \frac{360}{\lambda_{\text{sky}}} \approx \frac{360}{v_{\text{sky}} f_A}, \tag{4.7} \]

where \( v_{\text{sky}} \) is the telescope scan speed in the number of degrees per second on the sky. For a typical constant-elevation scan, this is just given by

\[ v_{\text{sky}} = v_{\text{az}} \cos(\theta_{el}), \tag{4.8} \]

where \( v_{\text{az}} \) is the telescope scan speed in number of degrees per second in azimuth.

Information on smaller scales than \( \sim 1' \) is suppressed through convolution with the detector beam response, so SPTpol is not sensitive to modes above \( \ell \sim 10000 \) or \( f_A \sim 15 \text{ Hz} \) for the 500 \( \text{deg}^2 \) lead-trail scan speed. Similarly, the spatial extent of the field puts a lower limit on the multipole range that can be measured. On a rectangular field, the largest mode that can be measured has a wavelength of \( \theta_{\text{max}} \), where \( \theta_{\text{max}} \) is the width of the largest dimension of the field on the sky. In practice, 1/f fluctuations in the data mean the largest modes are completely noise dominated, while at high spatial frequencies the power in CMB fluctuations is reduced and instrumental noise dominates.

We suppress high-frequency noise by applying a harmonic low-pass filter based on a chosen maximum \( \ell \). The cut-off is converted from \( \ell \) to audio frequency in the timestream at the bottom of the field, where \( v_{\text{sky}} \) is at its minimum. A low-pass filter is then applied to FFTed data after an appropriate anti-aliasing filter is applied, removing modes above the audio frequency corresponding to \( v_{\text{sky}} \). For E-mode analyses, we set this filter cutoff at \( \ell = 10000 \), while for B-mode power spectra, \( \ell_{\text{max}} = 6600 \).

The primary low-frequency contaminant is expected to be atmospheric fluctuations, both from scanning across volumes of air with varying opacity and from the motion of the atmosphere itself. We employ two different types of high-pass filters: an
FFT high-pass filter, in which the data is FFTed and low-frequency bins are zeroed, and a timestream polynomial subtraction using the Legendre polynomial basis. Polynomial subtraction allows a more direct removal of atmospheric fluctuations, since simple expected modes, such as a linear drift across a scan can be complicated in harmonic space. However, polynomial subtraction introduces slight leakage between frequency-domain modes in the maps, owing to the non-locality of the polynomial in frequency space. The harmonic-space high-pass filter inherently avoids such leakage, so the final filtering choices are a balance between the effective removal of 1/f noise and the retention of low $\ell$ signal. For the deep field EE and BB power spectra, we remove a fourth- and third-order polynomial, respectively, from each scan during map making.

4.3.3 Weights

In order to properly combine the contributions of each detector to the measured temperature of each map pixel, they are weighted according to the individual detector noise. The noise weighting for each detector is the inverse of the average PSD in a chosen frequency band, typically 0.3-2 Hz, that corresponds to the audio frequency of desired CMB sensitivity. Dedicated pre-observation noise stares can be used for extracting the weights, or else they are constructed by averaging differenced (leftgoing - rightgoing) scans to subtract out the sky signal. For the deep field power spectrum analyses, we use the latter option to most accurately capture the instantaneous noise on the CMB field.

4.3.4 Polarized Maps

Each detector has two fundamental quantities associated with it: a polarization angle, $\psi$, defined relative to the purely elevation direction of the telescope, and polarization efficiency $\rho$, the fractional coupling to the polarized component of the incident signal.
For a perfect linearly polarized detector, $\rho = 1$, while a detector that has no ability to discriminate polarization has $\rho = 0$.

In describing the construction of polarization maps, I will adopt the notation convention that most closely matches the processing used in the software pipeline. Before polarized maps are created, we calibrate the data as described in Section 4.2.4, which couples detector response to a signal in brightness temperature. The resulting gain factor, $C$ includes a factor of $(1 - \rho)$ since any cross-polar coupling increases the detector’s response to unpolarized light. Including a detector noise term $n_t$, each detector sample is then:

$$d_t = C(I_s + \frac{\rho}{2-\rho}Q_s \cos(2\psi) + \frac{\rho}{2-\rho}U_s \sin(2\psi)) + n_t.$$  \hspace{1cm} (4.9)

Each detector has a “Stokes coupling” that describes how its samples project onto the I, Q, and U basis defined on the sky:

$$SV = \begin{pmatrix} 1 \\ \frac{\rho}{2-\rho} \cos(2\psi) \\ \frac{\rho}{2-\rho} \sin(2\psi) \end{pmatrix}.$$  \hspace{1cm} (4.10)

During the map making procedure each detector data sample, $d_t$ is assigned to a pixel based on its corrected pointing information, and multiplied by its timestream weight (described in Section 4.3.3). For a given timestream weight factor, $W_i$, the contribution of a detector to the output maps is:

$$\begin{pmatrix} I \\ Q \\ U \end{pmatrix} = d_{t,i}W_iSV.$$  \hspace{1cm} (4.11)

We record the pixel weights as 3x3 matrices for each pixel, where the distribution of
the detector weight contributions to each element come from:

\[
\begin{pmatrix}
II & IQ & IU \\
QI & QQ & QU \\
UI & UQ & UU
\end{pmatrix}
= d_{t,i} W_i S V S V^\dagger,
\]

where the diagonal elements are the map pixel weights and the off-diagonals are their covariances.

### 4.3.5 E and B Estimators

Constructing E and B mode maps out of the Q and U maps as described in Section 1.2 does not work in the flat sky approximation. The non-locality of E and B, coupled with the partial sky coverage, means that certain modes are ambiguous in terms of their definition as “E” or “B”. Mis-labeling an underlying B-mode as an E-mode is not a problem, as the latter has a much larger amplitude so the bias is negligible. Leakage of source E into B, on the other hand, can dominate the resulting measured B-mode signal. To avoid this leakage, we employ a “pure” B-mode estimator, which is defined as the estimator containing no E signal by construction [45]. For constructing E maps, we use a basic transform from the Q, U basis to the E basis by a simple linear combination of Q and U maps. The inherent assumption here is that the entire polarization field is E modes, and that the B→E contamination is negligible.

### 4.3.6 Masks

Harmonic decomposition of finalized maps is extremely sensitive to real-space artifacts that have complicated harmonic structure. In particular, sharp features at the edges of the map where array coverage is poor can introduce numerical errors when pixel weights are removed, while celestial point sources cause ringing in the Fourier transformed maps. We address these issues by applying two masks, one for edge ef-
fects and one to remove point sources. The masks are constructed as described below and the multiplied together before being applied to the sky T, Q, and U maps prior to E and B mode estimation. They are also applied to the recovered E and B maps when power spectra are estimated.

To construct the edge mask, the map of pixel weights is first smoothed and then all pixels having a weight above some minimum factor, typically 0.3, of the median weight are given a mask value of 1. All other pixels, which include edge pixels and any inherent padding included when constructing the map, are given a mask value of 0. In maps that include many observations averaged together, this effectively sets the approximately uniform coverage region of the field to a mask value of 1. This mask is then convolved with a 2-d smoothing function, typically a revolved Hann window, to bring the mask from 1 to 0 in a way that induces minimal ringing in harmonic space. The width of this edge apodization can have a considerable impact on the reconstruction B modes, particularly at low $\ell$, when they are generated from Q and U maps, as shown in Figure 4.4.

Observations by other experiments inform a list of bright point sources in our fields that bias the CMB signal, especially at high $\ell$. We remove them by applying a point source mask that essentially zeros pixels close to a known source. We typically select only the brightest sources to mask, as the fainter sources have a minimal impact on CMB measurements, while masking them removes sky and reduces sensitivity. For a chosen list of sources, a mask is generated that begins as all ones. A circle of pixels, typically several beam widths in radius and centered on the pixel containing the point source, is set to zero. Finally, the mask is convolved with a 2-d smoothing function similar to one used in the edge mask.

The smoothing functions applied to the edge and point source masks are defined by their width, which for the Hann profile, is given as the radius of the cosine taper from the center of the function to its edge. For the deep field power spectra, we apply
a 10 arcminute-radius smoothing function to the point source masks, and a 60 (30) arcminute-radius function to the edge mask.

![Figure 4.4: E to B leakage from map edge effects](image)

When using the $\chi_B$ estimator to transform from Q and U maps to a B map, pixels near the edge of the map induce errors that show up as a spurious low-$\ell$ power. The problem is reduced when wide apodization masks are applied to bring the maps smoothly to zero at their edges. The “$f_{\text{sky}}$” term in the legend describes the fractional effective area of the observed field after application of the apodization mask, and is the integral of the mask divided by the integral of the nominal field.
Chapter 5

POLARIZED POWER SPECTRA
WITH SPTPOL

Starting from an accurate measurement of the microwave sky, a range of analyses have been performed with SPT data, including searching for the Sunyaev-Zeldovich signature of massive galaxy clusters [47], and studying submillimeter galaxy (SMG) sources [50]. For analyzing the properties of the CMB itself, the fundamental data product is an estimate of the angular power spectrum. In this chapter I will describe the steps undertaken to produce a power spectrum measurement and how a measured spectrum is used to constrain cosmology.

5.0.7 Bandpowers and Covariances

The only way to make a true, unbiased measurement of the CMB angular power spectrum $C_\ell$ in the native space of spherical harmonics, $Y_{\ell m}$, is to measure the whole sky. A perfect full-sky measurement would produce an angular spectrum by:

$$C_\ell = \frac{1}{2\ell + 1} \sum_m |a_{\ell m}^* a_{\ell m}|. \quad (5.1)$$
Once data has been removed from the map, i.e. by masking regions contaminated by non-CMB emission (like the galaxy) with a window function $W(\hat{n})$, we use pseudo-$C_\ell$ estimators of the underlying spectrum. Rather than producing a spectrum consisting of measured variances among individual $a_{\ell m}s$, pseudo-$C_\ell$ methods generate estimates of the power, $\tilde{C}_\ell$, related to the source spectrum by a transfer function, $K$,

$$\tilde{C}_\ell = K_{\ell \ell'} C_{\ell'}.$$  \hspace{1cm} (5.2)

As sky coverage decreases for a particular beam size, so too does the inherent independent sampling interval in $\ell$ space, such that a spacing of $\Delta \ell \approx \frac{2\pi}{\theta}$ for a field with linear dimension $\theta$ is required for reasonably independent estimates of the spectral shape [27]. Where the theoretical $C_\ell$ spectrum is expected to be smooth, averaging wider ranges in $\ell$ can improve the signal-to-noise of the measured spectrum. Portions of the measured spectrum that have been averaged over some range in $\ell$ space are referred to as “bandpowers”, and they are the fundamental unit of analysis for CMB power spectrum science. Analyzing the cosmological implications of a set of measured bandpowers typically involves a Monte Carlo analysis of the data compared to models produced by a program like CAMB. Such analysis requires estimates of covariance between the bandpowers, but raw covariances are often plagued by numerical issues, requiring a more careful conditioning of the covariance matrices. At the end of an analysis project, the two final data products are a set of bandpowers and the associated bandpower covariance matrices.

### 5.1 Generating CMB Power Spectra

For power spectrum analyses, SPTpol follows closely the analysis pipeline developed for SPT-SZ, even to the extent of re-using some of the same analysis code. The overall process is based on the MASTER algorithm developed for BOOMERanG [20]. Begin-
ning with partial-sky maps and the window functions describing their weighted sky coverage, we derive a set of so-called bandpowers, which are unbiased estimates of the average sky power spectrum within a particular range of $\ell$. Bandpower widths depend on the data quality and analysis requirements, but are wide enough to be mostly uncorrelated and include some averaging of the noise on individual $C_\ell$ estimates, while narrow enough to avoid averaging over features in the underlying spectrum. CMB spectra vary more slowly in a parametrization that includes the increasing number of modes contributing to each $C_\ell$, so we operate on the corrected spectrum:

$$D_\ell = \frac{\ell(\ell + 1)}{2\pi} C_\ell.$$  

(5.3)

Comparing data-derived bandpowers to a model source spectrum requires disentangling the bias induced by the map processing steps outlined in Section 4.3. The entire process can be represented in linear algebra form as

$$\tilde{D}_b = P_{b\ell}[M_{\ell\ell'}(W)F_{\ell'}B_{\ell'}^2]Q_{\ell'b'},$$  

(5.4)

where $P_{b\ell} = \frac{1}{2\pi} \frac{\ell(\ell + 1)}{\Delta\ell}$ is the binning operator that converts from native $\ell$ spectrum space to bandpowers with a width $\Delta\ell$, $Q_{tb} = \frac{2\pi}{\ell(\ell + 1)}$ an interpolation of a bandpower $b$ back into $\ell$ range it spans. The mode mixing kernel, $M_{\ell\ell'}(W)$ details how power mixes in $\ell$ due to finite sky and apodization. The transfer function $F_{\ell}$ describes the effect of conditioning of the detector data and collapsing it into map pixels, while the beam function $B_{\ell}^2$ encodes the smearing of sky signal due to the detector angular response functions. In practice, the individual components can be calculated independently and combined into a single transformation matrix, the $K$ matrix, that relates a biased measured spectrum, $D^m_b$, to its underlying source source spectrum, $D^s_b$:

$$D^m_b = K_{bb'} D^s_{b'}.$$  

(5.5)
For the unbiased estimates of data bandpowers to be compared to a model, the theoretical CMB model spectrum, $D_{\ell}^{th}$, has to be transformed into the same band-power basis as the data by applying the same operations assumed in their estimation,

$$D_{\ell}^{th} = (K_{W})[P_{\ell\ell^{'}}(W)F_{\ell}B_{\ell^{'}}^2]D_{\ell}^{th}. \quad (5.6)$$

The bandpower window function, $W_{\ell} \equiv (K_{W})[P_{\ell\ell^{'}}(W)F_{\ell}B_{\ell^{'}}^2]$, captures all of the transformations between native $\ell$ resolution and the re-binning of the power spectrum into the averaged $\Delta \ell$ bins. In the following sections, I will describe in greater detail the processing that goes into generating bandpower estimates.

### 5.1.1 Cross Spectrum Analysis

The more observations of a field that are averaged (in the limit that noise is predominantly Gaussian), the greater the map depth, resulting in a higher signal-to-noise power-spectrum estimate. This would naively suggest that all field observations should be combined to achieve a maximally deep map from which to extract band-power measurements. However, taking the autospectrum of a single map includes an inherent noise bias, and each pixel measurement contains both signal and noise. To the extent that the noise spectrum is well-understood it can be removed, but such removal is fraught with systematic and statistical uncertainty.

An alternative means of deriving a power spectrum is to average many cross-spectra of individual maps or small groups of maps added together. Each map contains the same signal, but with nominally uncorrelated noise realizations. If the noise in uncorrelated, the cross-spectrum of two maps is an unbiased measure of their common signal. SPTpol employs a cross-spectrum estimator, similar to the one used for temperature spectrum analysis with SPT-SZ [32], to produce bandpowers. We group maps as described in the next section, such that we have N independent ob-
observations of the sky. We then take the 2-d FFT of each map and derive a spectrum measurement from the real part of their complex product, or

\[
D_i^b = \left\langle W_{kk'} M^a(\vec{k}) M^{b*}(\vec{k'}) \right\rangle_{\vec{k}\in b} \\
D_i^b = S_b + N_i^b
\]  

(5.7)

where \(W_{kk'}\) is the weighting of individual \(\vec{k}\)-mode contributions to each bandpower, \(\frac{|k|(|k|+1)}{2\pi}\). The mean of all individual cross spectrum bandpower measurements is then the final set of raw bandpowers.

The covariance matrix of the bandpowers has two components: sample variance and noise variance. The former is estimated directly by generating a number of simulated, noise-free CMB skies from a fiducial cosmology and then reconstructing the bandpowers from the same sky coverage as our field observations. This captures the result of both the inherent cosmic variance of \(C_\ell\)s on the sky and the effect of partial sky coverage only sampling a portion of all possible modes at each scale. Each measured cross-spectrum has the same signal but a (nominally) independent noise realization, so the variance on the measured bandpowers is simply the power spectrum of the noise, \(\sigma^2 D_b = \sigma^2 N_b\). We add this noise variance component in quadrature with the sample variance to construct the initial bandpower covariance matrix.

We choose windows and bandpower widths to minimize the correlations between \(\ell\) bins, so the final bandpower covariance matrix is dominated by its diagonal and near-diagonal elements. The near-diagonal correlations are a generic result of the partial-sky window function, and must be incorporated into the covariance matrix. However, the finite number of \(D_i^b\) realizations means that the far off-diagonal elements of the matrix are subject to statistical noise. In the limit of infinite simulations \((i \to \infty)\), this noise goes to zero leaving only near-diagonal matrix elements nonzero.
5.1.2 Combining Maps

Cross-spectrum analyses, and in particular the estimates of bin-bin covariances, are inherently error-on-the-mean calculations. Thus, the bandpower estimate improves as the square root of the number of realizations. Ideally, the number of measurements, \( N_m \), would far exceed the number of parameters being measured, \( N_p \), where in this case \( N_m \) is the number of individual cross spectra and \( N_p \) is the number of bandpowers being estimated. Too few measurements relative to the number of parameters being estimated results in excess uncertainty on the covariance estimate [13]. This suggests using individual observations as the map units \((\alpha, \beta)\) in Equation 5.7. However, because of our scan strategy and the separation of individual pixel beams on the sky, the coverage of individual maps is poor. The construction of maps as arrays of weighted pixels means that any operations (such as combining maps) that involve removing the weight means that divide-by-zero errors can become problematic. Furthermore, the cross spectrum analysis for \( N \) maps involves \( \frac{N(N-1)}{2} \) multiplications between their FFTs which severely taxes computing resources.

The map coverage issue is addressed by grouping individual maps together into “bundles” that have improved coverage relative to single observations. An ideal bundle has a single observation performed at each dither step to ensure maximum field coverage. We also want to have a total of \( N_{\text{band}} \) bundles where \( N_{\text{band}} \sim N_b \). The bandpower covariance smoothing described above reduces the number of independent values being estimated for each bandpower/covariance combination, so having the same number of map bundles as bandpowers results in a less than 5\% excess variance on the bandpower values.

For the 2013 500 deg\(^2\) field, there were a total of 680 good observations in the lead-trail configuration, with 18 distinct dither steps. The natural bundling scheme only produces 37 maps, which is too few to ensure minimal excess uncertainty in the bandpower estimates. We avoid this by re-grouping the dither steps by every third
step, so the new dither step 1 is built from original steps 1-3, new step 2 from original steps 4-6, etc. Bundles are then built from a single map from each new dither step, for a total of 113 map bundles.

5.1.3 Mode-mixing

The dominant effect of sampling only part of the sky is the mixing of power between angular-spectrum bins that results from incomplete sky coverage. In the flat-sky limit, this reduces simply to the case of bin-to-bin leakage in the FFTs, which then propagate into the power spectrum estimate. The effect of a particular sky window function (the shape of the map itself, after masks have been applied) can be calculated analytically or by simulation. The analytical calculation requires some well-motivated simplifications that can be checked by simply creating maps with power in a single spatial mode and then propagating them through the analysis pipeline to track where the final power is distributed. In either case, the mode mixing is measured at a finer $\ell$ bin resolution than the eventual bandpowers to ensure that any subtle effects are captured. For bandpower spacing of $\Delta\ell = 50$, we calculate the analytical mode mixing at $\Delta\ell = 5$. Figure 5.1 shows a comparison of the analytical approximation with the simulated version, for a single bin of input power.

5.1.4 Simulations

Using the metadata stubs generated during IDF creation, as described in Section 4.2.6, we can generate simulated timestreams with the same properties as our real data. We have measurements of our average detector beam and the polarization angles of individual detectors, as well as calibration factors relating detector response to sky brightness. Beginning with a particular source CMB power spectrum, typically taken as the best-fit $\Lambda$CDM model (supplied by the Planck satellite at the time of this writing), we generate a series of sky realizations. Using our measured detector
5.1.5 Transfer Function

Timestream filtering operations described in Section 4.3.2 are necessary to mitigate low frequency noise, but CMB signal is suppressed along with the noise. In the
final bandpower estimate, the effects of filtering are accounted for by the transfer function, $F_\ell$ in Equation 5.4. To construct $F_\ell$, we begin with sky realizations described in Section 5.1.4, sample them into simulated timestreams, then push the simulated data through the bandpower estimation pipeline. Spectra recovered from the filtered simulated maps are then compared to the input spectra, and their ratio as a function of angular scale is $F_\ell$. The final bandpowers and error bars are divided by the transfer function to correct for the effects of filtering. Where the transfer function is low, the errors on bandpowers are inflated, while a value near 1 in $F_\ell$ corresponds to a minimal impact of data processing on the power spectrum estimation. Removing data when filtering low-frequency noise reduces power across the spectrum, so the transfer function never reaches 1. A plot of the transfer function for the deep field EE spectrum is shown in Figure 5.2.

![EE Filtering Transfer Function](image)

**Figure 5.2: EE Spectrum Filtering Transfer Function**

Average transfer function between input EE spectra from ΛCDM best fit and the recovered spectra from observing simulated skies. Where the transfer function approaches zero, the sensitivity to CMB signals is suppressed by the removal of modes dominated by 1/f noise, so dividing it out magnifies the errors in those bins.
5.1.6 Temperature Deprojection

A dominant source of contamination is leakage of temperature signal into polarized maps. There are various mechanisms that can cause $T \rightarrow P$ leakage, but for our purposes the exact origin is less important than our ability to identify and remove it. The simplest form of leakage is the monopole component, which is just a scaled copy of the temperature map in the Q and U maps. We identify the monopole leakage by cross-correlating TQ and TU maps to generate their cross-spectra, then scaling the results by the TT autospectrum. Since the average correlation between Q/U maps and T maps should be zero any residual correlation indicates the presence of leaked T in the Q or U map. For the SPTpol deep field data, we find that the $T \rightarrow Q$ and $T \rightarrow U$ coefficients are $-0.0062$ $(0.0045)$ and $-0.0001$ $(-0.0073)$ for 90 (150) GHz maps. Using the measured leakage fractions, we remove an appropriately scaled version of the T map from each Q and U map in a particular bundle.

5.1.7 Jacknives

As a check against likely sources of systematic errors we perform a series of “jackknife” tests. To perform a jackknife, we first rank observations according to a metric that tracks the systematic being tested, such that the first and second halves of the ranked list of observations are maximally different with respect to the systematic under test. We then select a map from each half of the list and difference them to get a “difference map”, which is then processed the same as a data map to generate a set of bandpowers. The bandpowers for each difference map are then cross-correlated to produce an unbiased jackknife spectrum. In the simplest case, the jackknife spectrum is consistent with zero, with variances on each bandpower corresponding to the data noise covariances. However, various effects such as changes in receiver tuning, alterations to the observation strategy, or different map weightings can result in a residual signal that does not indicate systematic error, but rather an uneven weighting of the
two halves of the data, for example. We account for this by generating “expectation bandpowers” for each jackknife, using simulated skies sampled with IDF stub metadata so that any residual not attributable to the systematic being tested will persist.

Jackknives can be performed on nearly any separation of the data, provided an appropriate ranking metric can be devised. In the absence of reasonable concern for a particular source of systematic contamination, we perform jackknives based on three metrics:

1. Date

Here the observations are simply ordered by date and split in half chronologically to ensure that the instrument was not subject to a monotonic change in performance over the course of the season.

2. Azimuth

A target CMB field is located at different telescope (az, el) coordinates depending on the time of day when the observation begins. Test maps made in ground-centered coordinates indicate variable ground contamination as a function of azimuth. For this test, maps are grouped by az ranges with the least vs. most expected ground contamination. This test checks the assumption that ground pickup is suppressed by the random distribution of azimuth ranges.

3. Left-right

Maps are processed as leftgoing-only and rightgoing-only pairs, and are then typically recombined as the first step of subsequent analysis. For the this test, we instead generate difference maps between leftgoing- and rightgoing-only maps. Nominally, the sky signal in the two sets of maps should be completely correlated (and subtract out), while the noise is completely uncorrelated. This test checks in particular for any scan-synchronous effect that correlates with
telescope motion.

We evaluate the results of a jackknife by measuring the $\chi^2$ of the $N_{jk}$ data difference-map bandpowers relative to the expectation spectrum and calculating the statistical probability-to-exceed (PTE) the resulting $\chi^2$ for a $N_{jk}$ degrees of freedom. The PTE values of the individual tests are then checked for their consistency with a uniform distribution to support the conclusion of no systematic contamination of the reported bandpowers.

5.2 SPTpol Power Spectrum Results

The first science result from SPTpol was an indirect measurement of lensing B-modes in polarized CMB maps [17]. A lensing potential map is derived from a measurement of the CIB at 500 $\mu$m from the SPIRE instrument on the Herschel satellite, along with previous work relating the CIB to the lensing potential, $\phi$. SPTpol E-mode polarization maps are then taken as the source CMB field and cross-correlated with the CIB-derived maps of $\phi$. A result of the cross-correlation is a map of residual B-modes resulting from the deflection of the source E-modes by the weak lensing from $\phi$, using an appropriate kernel to describe the lensing interaction. This template B-mode map, which is the product of two high signal-to-noise products, is then cross-correlated with the much noisier SPTpol measured B-mode maps. An autospectrum of the B-mode data itself is subject to a considerable noise bias, which instead becomes an unbiased variance on the cross-correlation measurement with the template B-modes. The resulting cross power spectrum bandpowers of B-modes from the template maps crossed with SPTpol data maps is shown in Figure 5.3.

The signal-to-noise in E-mode maps from SPTpol data was high enough at both 90 and 150 GHz to produce B-mode templates in cross-correlation with the CIB-derived $\phi$ maps, and both templates produced similar results as part of testing the spectral
dependence of the result. However, only the SPTpol 150 GHz measured B-mode maps were correlated with the template maps because of their greater depth and, at the time the result was published, a more thorough understanding and control of systematics.

![Figure 5.3: SPTpol Lensing BB Detection](image)

Derived BB bandpowers from a cross correlation of SPTpol data with a B-mode map template. The template was formed by cross-correlating an SPTpol 150 GHz E-mode map with a map of the lensing potential, $\phi$, from Herschel CIB measurements and previously-established correlations between the CIB and $\phi$ [17]. Only 150 GHz SPTpol B-mode maps were used in the correlations owing to their greater depth than the 90 GHz maps.

At the time of writing, polarized power spectra from deep field observations are in preparation, while analysis and generation of data products described in the previous section are ongoing for the first season of survey field observations. We are currently observing the survey field for a second season, and plan to have a total of three seasons devoted primarily to survey field observations.

Analysis choices described in previous sections are what we use for the initial
power spectra from deep field observations. A larger field is necessary to have any signal-to-noise in the $\ell$ range of gravity wave B-modes, so the focus of the deep field B-mode analysis is the lensing spectrum. The E-mode measurements for both the deep field and survey field represent a marked improvement over the first generation of CMB polarimeters, but will soon be surpassed in overall sensitivity by data from the Planck satellite, which has a much lower cosmic variance term in its noise owing to greater $f_{\text{sky}}$.

In March 2014, the BICEP2 collaboration announced a detection of inflationary B-modes with a best-fit $r = 0.20^{+0.07}_{-0.05}$, from observing the same region of sky as the SPTpol wide field [1]. The result generated considerable excitement in the CMB community, as previous temperature-only data suggested a lower limit of $r = 0.11$ [2]. Deep field observations with SPTpol do not yield enough sensitivity at low $\ell$ to follow up on the BICEP result, but the SPTpol survey overlaps with the BICEP2 field, offering the possibility of cross-correlation between BICEP2 and SPTpol maps.

The primary ongoing challenge with the SPTpol survey field data is mitigating 1/f noise in the data to maintain our sensitivity to low $\ell$ modes. For the 1.09 $\deg/s$ scan speed used for the survey field observations during the 2013 season, a 1/f knee in the timestreams of 0.1 Hz corresponds to $\ell \sim 70$ on the sky. Low-frequency noise in a single detector is expected to be dominated by variations in unpolarized atmospheric opacity. Thus, during the construction of polarized maps, the atmospheric 1/f component should be suppressed by the implicit differencing of the timestreams in the detectors within a pixel. Preliminary noise data analyzed prior to the development of the full calibration pipeline described in the previous sections suggested a 1/f knee near 0.1 Hz. An active effort is presently underway to identify the origin of 1/f noise in our data and ensure that the analysis pipeline is tuned to suppress it as much as possible without also suppressing polarized signal.

A nominal forecast for the combined 90 and 150 GHz BB spectrum bandpowers
Figure 5.4: SPTpol Preliminary EE Power Spectra

Preliminary EE power spectra from SPTpol observations of the deep field (top) and survey field (bottom). The sample variance term for the deep field data is inflated relative to the survey field by a lower $f_{sky}$, as shown by the grey shaded regions, but in both cases it dominates the error bars on band powers below multipoles of $\sim 1500$. 

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Figure 5.5: SPTpol Sensitivity Forecast

Bandpower estimates from SPTpol, using the target map depth from 3 years of survey field observations with 90 and 150 GHz data combined, compared with the 2014 BICEP2 B-mode data. The sample variance range shown is calculated with an $f_{\text{sky}}$ roughly equivalent to the sky observed by both experiments, and is interpolated between the limits for each nominal SPTpol bandpower.

from the full SPTpol survey is shown in Figure 5.5, with the published BICEP 2 bandpowers for comparison. At the multipoles where the gravity wave component of the $r = 0.2$ theory spectrum peaks, both BICEP 2 and SPTpol are near the sample variance limit. However at higher multipoles, where the lensing spectrum amplitude comes to dominate, the higher resolution of SPTpol affords a more precise measurement of the spectrum.

While the bandpowers shown are from the combined 90 and 150 GHz instrumental sensitivity, SPTpol’s multi-band data allows for a test of the spectral dependence of any B-mode signal, which is crucial for separating CMB signals from poorly-understood polarized galactic foregrounds. At the time of this writing, the raw SPTpol 90 GHz maps are the deepest in a complimentary band to the BICEP 2 data, and because both experiments observed the same field, a multi-band cross-
spectrum analysis is possible. In addition to clarifying the potential contribution of polarized foregrounds, such an analysis provides a strong check on the systematics of each experiment, which are uncorrelated in the respective maps.
Chapter 6

Conclusion

The SPTpol instrument was designed as a high-resolution, polarization-sensitive mm-wave camera using TES detectors to achieve background-limited noise performance while observing from the South Pole. In order to probe outstanding questions about the earliest moments of the Universe and its subsequent evolution, we seek precise measurements of CMB polarization and its B-mode component in particular. Such measurements require low-noise detectors and tight control on a wide range of systematics.

The detector development program yielded a dichroic focal plane comprised of 1536 polarization sensitive bolometers, using distinct architectures for detectors in the two different bands. Thermal stability of the TES sensors was tuned by adding bling to increase the TES heat capacity, while superconducting Nb structures were deposited directly onto the TES material to reduce the slope of the superconducting transition. We measured a $< 1\%$ coupling to cross-polar radiation, along with $90 \%$ detector optical coupling efficiency and combined dark detector white noise below the expected optical loading for a range of bias points in the transition.

In the deployed instrument, alignment angles for individual detectors are measured to better than $2^\circ$ statistical precision, resulting in an array-averaged statistical error
of less than 0.1° on the focal plane alignment. A series of systematic tests on the polarization calibration measurement yields systematic uncertainties on the mean angle of 1.22° (0.78°) for 90 (150) GHz detectors. We replaced optical filters and detector feeds after the first season of observations in 2012, reducing atmospheric loading and improving the NET values for the entire focal plane.

We developed a data processing pipeline that transforms raw detector outputs into polarized CMB maps that have random glitches and anomalously noisy data excised prior to any subsequent processing. Persistent, array-correlated noise from the telescope cooling system is identified and precisely removed with a minimal loss of bandwidth, eliminating features that would otherwise be present in our final maps.

Data from the first season of observations with SPTpol produced the first significant detection of B-modes in the CMB in 2013. At the time of writing, SPTpol is analyzing both 90 and 150 GHz data on the same field the BICEP 2 collaboration observed for their reported detection of primordial gravity wave B-modes. In addition to testing the BICEP 2 result with an independent data set, SPTpol data will yield a precise measurement of the lensing B-mode spectrum.
Bibliography


