UNIVERSITY OF CALIFORNIA, SAN DIEGO

Experimental Tests of the Theory of Poloidal Rotation in the DIII-D Tokamak

A dissertation submitted in partial satisfaction of the requirements for the degree
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in

Physics

by

Colin Chrystal

Committee in charge:

Professor Patrick Diamond, Chair
Professor Robert Cattolica
Professor David Kleinfeld
Professor Clifford Surko
Professor George Tynan

2014
The dissertation of Colin Chrystal is approved, and it is acceptable in quality and form for publication on microfilm and electronically:

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__________________________________________ Chair

University of California, San Diego

2014
EPIGRAPH

Murphy!
—Keith Burrell
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VITA

2007
B. S. in Applied Physics, California Institute of Technology

2009
M. S. in Physics, University of California, San Diego

2014
Ph. D. in Physics, University of California, San Diego

PUBLICATIONS


ABSTRACT OF THE DISSERTATION

Experimental Tests of the Theory of Poloidal Rotation in the DIII-D Tokamak

by

Colin Chrystal

Doctor of Philosophy in Physics

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Professor Patrick Diamond, Chair

The goal of this dissertation was to develop a novel technique for measuring ion poloidal rotation and then using that technique to test poloidal rotation theories. The new poloidal rotation diagnostic has been developed on the DIII-D tokamak. This diagnostic uses charge exchange recombination spectroscopy to measure toroidal rotation on the high- and low-field side of the tokamak midplane to determine the poloidal rotation from a divergence-free description of flow within flux-surfaces. Measurements are made such that no atomic physics calculations are needed to account for the energy dependence of the charge exchange cross section. New techniques for creating magnetic equilibrium reconstructions and performing the spatial calibration have been developed to ensure the accuracy of this new diagnostic. Measurements are made in the core of DIII-D where the spatial
resolution is significantly improved when compared to the direct measurement of
the poloidal rotation.

This diagnostic has been used to investigate impurity poloidal rotation in
the core of a variety of plasmas. For the first time on DIII-D, mean poloidal flow
spin-up coincident with the formation of an internal transport barrier has been
observed. The various measurements of poloidal rotation have been compared
with theoretical predictions. Disagreement with neoclassical calculations have been
found in H-mode, QH-mode, and the core of internal transport barrier plasmas.
The effect of turbulent driven Reynolds stress and fast-ion friction have been
investigated as well, and it has been determined that either of these effects, on
their own, is insufficient to explain the discrepancy with neoclassical predictions.
Modeling results indicate that these effects that are not included in standard
neoclassical calculations are important for calculating the poloidal rotation.
Chapter 1

Introduction

This dissertation describes results of a new poloidal rotation diagnostic that has been developed on the DIII-D tokamak at General Atomics. Poloidal rotation is one of two rotations present in a tokamak and a fundamental element of tokamak plasma physics because charged particles in the tokamak are subject to a \( qV \times B \) force. This diagnostic uses an expansion of an existing measurement capability and innovative analysis to more precisely determine the poloidal rotation. These results are used to test the theory of poloidal rotation, which is an important aspect of understanding tokamak plasma physics. The following chapters detail the theory of poloidal rotation, the principles of this new diagnostic, and the studies of poloidal rotation theory that have been performed.

1.1 Tokamak Fusion

Fusion is a nuclear process whereby separate nuclei collide with each other and form a single nuclei. For light elements, the nucleons that compose the fusion product have less potential energy than when they were in separate nuclei. As a result, energy is released from a fusion reaction in the form of energetic particles and kinetic energy of the fusion product. This process can only occur when the nuclei are able to approach close enough so that the strong nuclear force is able to overcome their Coulomb repulsion. An example fusion reaction that requires a relatively small amount of energy to initiate is the reaction between deuterium
and tritium,

\[ ^2\text{D} + ^3\text{T} \rightarrow ^4\text{He} (3.5 \text{ MeV}) + n \ (14.1 \text{ MeV}). \]

This reaction has practical significance because its cross section is maximized at the relatively low temperature of 70 keV (about 800 million K). The goal of fusion research is to develop a sustained fusion reaction that can produce a net gain in energy by acquiring energy from fusion products at a rate greater than the energy needed to sustain the reaction.

At the temperatures needed to create fusion reactions, electrons are dissociated from their nuclei and the resulting, difficult to control matter is called a plasma. In order to create a viable fusion reaction, this plasma must be contained so that fusion reactions produce energy faster than it is lost to the environment. As a result, the plasma must be sufficiently hot, dense, and insulated from its surroundings. These three qualities of the plasma are typically multiplied together and used as a figure of merit, the triple product,

\[ nT\tau_E. \]

Here, \( n \) is the plasma density, \( T \) is the plasma temperature, and \( \tau_E \) is the energy confinement time. The need to reach a minimum triple product to achieve a self-heating fusion reaction can be referred to as the Lawson criterion. A tokamak is a device that aims to create a viable fusion reaction by containing a relatively hot, low density plasma for a significant amount of time by creating a strong, stable magnetic field. Magnetic fields can contain charged particles by forcing them to execute gyro-orbits around the magnetic field lines, thereby reducing their ability to travel perpendicular to the magnetic field. The size of the gyro-orbit is reduced with increased magnetic field strength, thereby increasing the quality of the confinement. The ability to achieve high values for this triple product is a reason for the success of tokamaks in fusion research.

Tokamak fusion research consists of the study of plasmas that have been confined in a tokamak. The main goal is to understand their dynamics so that performance can be improved and significant fusion power can be achieved. Physically, a tokamak is a toroidal device that contains a plasma with a strong
magnetic field in the toroidal direction (the long way around the torus), and weaker magnetic field in the poloidal direction (the short way around the torus) that is generated by the plasma itself. To generate the poloidal magnetic field, current is driven in the plasma in the toroidal direction. This is accomplished by using the plasma as the secondary coil in a transformer where the primary coil is through the center of the toroidal device (the central solenoid). Driving current in the plasma provides some Ohmic heating, but this amount of heating is not sufficient to achieve a large triple product. External heating systems are therefore used and can increase the temperature and density of the plasma.

The strong toroidal magnetic field is created by external coils that surround the tokamak. Plasma shape and control are accomplished with external coils that add to the plasma-generated poloidal magnetic field. Charged particles easily flow along the magnetic field lines, a path that stays within the tokamak, but have reduced mobility across the field lines, a path that leaves the plasma.

1.2 The DIII-D tokamak

DIII-D is a medium-sized tokamak that has been operated by General Atomics in San Diego, California since 1986. A schematic of the tokamak is shown in Fig. 1.1, and its basic parameters are shown in Table 1.1. An extensive diagnostic suite makes DIII-D well suited to studying the physics of tokamak plasmas. In addition, the ability to greatly modify the basic plasma parameters (e.g. field, current, and shape), have allowed DIII-D to produce world-leading research well into its third decade. DIII-D typically produces deuterium plasmas where carbon is the dominant impurity, due to its graphite tiled walls. In this dissertation, analysis of plasma discharges is accompanied with a record of the basic parameters for that plasma (when applicable): $B_\phi$ (the toroidal field at the center of the vacuum vessel), $I_p$ (the total plasma current), $n_e$ (the average electron density across the vessel midplane), $P$ (the auxiliary heating power which can be from neutral beam injection, electron cyclotron heating, or both), and $\tau_{NBI}$ (the torque on the plasma from the neutral beams).
Figure 1.1: A Schematic of the DIII-D tokamak. The different magnetic field coils are labeled and a person is shown for scale.

Table 1.1: Basic parameters of the DIII-D tokamak. Note that the strength and direction of $B_\varphi$ and $I_p$ can be altered, so their maximum magnitudes are listed here. $R$ is the major radius and $a$ is the minor radius.

<p>| | |</p>
<table>
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<td>$</td>
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<td>$</td>
<td>I_p</td>
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<tr>
<td>$R$</td>
<td>1.7 m</td>
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<tr>
<td>$a$</td>
<td>0.6 m</td>
</tr>
<tr>
<td>$n_e$</td>
<td>$0.1 - 1 \times 10^{20}m^{-3}$</td>
</tr>
<tr>
<td>Pulse Length</td>
<td>10 s</td>
</tr>
<tr>
<td>$P_{aux}$</td>
<td>25 MW</td>
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1.2.1 Magnetic equilibrium reconstruction

Magnetic confinement of fusion is based on the fact that charged particles moving perpendicular to the magnetic field are bent into circles, reducing their ability to relax gradients in this direction. In a tokamak, magnetic field lines are created that are confined to a series of nested toroidal surfaces. Due to toroidal
symmetry and the fact that $\nabla \cdot \mathbf{B} = 0$, the magnetic field of a tokamak can be written as

$$\mathbf{B} = I(\psi) \nabla \hat{\varphi} + \nabla \psi \times \nabla \varphi,$$

(1.1)

where $\varphi$ is the toroidal angle coordinate and $\psi$ is the poloidal flux function, representing the amount of poloidal flux contained within a ring specified by a major radius $R$ (distance from the tokamak’s axis of symmetry) and height $z$ (coordinate measuring position along the tokamak’s axis of symmetry). The magnetic field is represented this way because surfaces of constant $\psi$, called “flux surfaces,” are the nested toroidal surfaces within which transport is along the magnetic field and fast (because $\mathbf{B} \cdot \nabla \psi = 0$), while transport across the surfaces is across the magnetic field and slow. These surfaces are centered around the magnetic axis, the magnetic field line created where the poloidal field inside the plasma is zero. An example contour plot of these surfaces is shown in the plasma cross sections of Fig. 1.1.

Accounting for lowest order force balance in the plasma and using this representation of the magnetic field allows the Grad-Shafranov equation to be written as

$$\Delta^* \psi = -\mu_0 R^2 P'(\psi) + \frac{\mu_0^2 F F' (\psi)}{4\pi^2},$$

(1.2)

where $\Delta^*$ represents the $R^2 \nabla \cdot (\nabla / R^2)$ operator, $P$ is the plasma pressure, and $F = 2\pi R B_\varphi / \mu_0$. Using various diagnostic measurements it is possible to reconstruct the magnetic field throughout the plasma, and thereby determine the shape of the nested toroidal surfaces of constant $\psi$. This reconstruction can be created for DIII-D with an analysis tool called EFIT. Such reconstructions form part of the foundation for a large portion of this dissertation which uses variation of toroidal rotation within a surface of constant $\psi$ to determine poloidal rotation.

Basic reconstructions can be created with measurements of the magnetic field that are made outside the plasma. More sophisticated reconstructions can be made by including measurements of the magnetic field pitch derived from the motional Stark effect (MSE) diagnostic. Still more sophisticated reconstructions can be made by also including measurements of the plasma pressure.

The presence of fast transport along flux surfaces motivates the creation of
magnetic coordinate systems which reduce the description of the tokamak plasma to a single dimension. Symmetry in the toroidal direction is usually taken as a given due to toroidal symmetry being fundamental to tokamak construction, and the ability to quickly equilibrate in surfaces of constant \( \psi \) allows some plasma parameters to be taken as constant within these surfaces. Such parameters are referred to as “flux functions.” For example, temperature is generally taken to be a flux function because heat can easily travel along a field line and equilibrate the temperature within its flux surface.

While a physical coordinate system, e.g. cylindrical coordinates \((R, \varphi, z)\), can be used to describe a tokamak, a magnetic coordinate system, e.g. \((\psi, \theta, \varphi)\) with \(\theta\) being the coordinate that goes around flux surfaces, is very useful. At DIII-D, it is common to use \(\rho\) instead of \(\psi\), as a flux surface label. While \(\psi\) represents the poloidal flux contained within a toroidal ring, \(\rho\) represents the toroidal flux within a flux surface. Typically, \(\rho\) is normalized to the total toroidal flux contained within the plasma boundary, and a one-to-one mapping exists between \(\psi\) and \(\rho\).

Toroidal symmetry and constant conditions within flux surfaces do not allow for a complete description of a tokamak plasma. Nevertheless, magnetic coordinates are useful for exploiting the symmetries of a tokamak plasma where they exist.

### 1.2.2 The DIII-D neutral beam system

The large majority of auxiliary heating on DIII-D is provided by neutral beams. Four beamlines each direct power from two, deuterium fueled, ion sources that each provide \(\approx 2.2\) MW of power. Typical accelerator voltages are 81 kV, but can be as low as 45 kV. However, the neutral beams are not mono-energetic due to the presence of \(D_2^+\) and \(D_3^+\) in addition to \(D^+\) in the ion source. The different beam energy components are referred to as “full”, “half”, and “third” energy components, and they complicate modeling of the interaction between the beam and the plasma.

After the injected neutrals ionize, the resulting fast-ions slow down on
the bulk plasma. This process adds heat and momentum to the plasma. The input momentum is mainly in the toroidal direction because this is the primary direction of neutral beam injection (NBI). In addition, the neutral beams provide a significant source of neutral particles in the core of the plasma, and these neutrals are the basis of important diagnostics such as charge exchange recombination spectroscopy (see Section 1.5).

The versatility of the DIII-D neutral beams is an integral part of many experiments. Most importantly, the eight neutral beams provide a variety of injection directions, allowing the direction of momentum input to be altered. Two of the neutral beams inject in a toroidal direction that is counter to the other six. The injection direction for these two beams is counter to the usual direction of the plasma current and is therefore referred to as counter-current. The six other beams are referred to as co-current. By varying the balance of co- and counter-current beam power, the heating and momentum input to the plasma can be decoupled. This ability is very useful for studying the effects of plasma rotation, discussed in Section 1.3. In addition, one of the neutral beams in each beamline injects more tangentially than the other, meaning it inputs more torque on the plasma for the same momentum input.

The neutral beams can be quickly modulated, with on times as short as 2 ms and off times as short as 10 ms. This allows the total heat and momentum rate to be reduced without changing the accelerator voltage, and it helps beam based diagnostics to separate measurements that are a direct result of the neutral beams from the background.

### 1.2.3 Beam into gas shots

In addition to being able to create plasma discharges, DIII-D is capable of creating beam into gas shots. These consist of filling the vessel with a low density, neutral gas that then serves as the target for neutral beam injection. Collisions between the neutral gas and the energetic neutrals result in atomic emission that can be recorded by certain diagnostics. For beam based diagnostics, this emission is particularly useful because it contains predictable emission that
must pass through all the same instrumentation needed to make measurements during a plasma discharge. For example, an optical diagnostic must use all the same lenses and windows to detect beam induced light for a beam into gas shot or a plasma discharge.

The resulting emission spectra contain well known emission from low-ionization states that is suitable for calibration purposes. Various gases can be used to fill the DIII-D vessel, including deuterium, helium, xenon, and neon. The choice of gas is determined by the atomic lines that result and their usefulness for various calibration needs. The gas pressures are typically $0.05 - 1.0$ mTorr.

The neutral beams exhibit very little attenuation when injecting into the low density gas present during beam into gas shots, and any beam neutrals that ionize are ballistic when no magnetic field coils are energized (as is typically the case). As a result, the total on-time for each neutral beam is limited to 50 ms, so that the neutral beams do not damage the spot on vessel wall where they strike after passing through the gas. This places a practical limit on the amount of signal that can be obtained from beam into gas shots.

### 1.3 Plasma rotation in tokamaks

Rotation is a ubiquitous feature in tokamak plasmas. Tokamak plasmas rotate in both the toroidal and poloidal directions, even in the absence of external momentum input from neutral beams. Ion rotation is of principle concern because the ions contain the large majority of the plasma’s inertia, and the term “rotation” is typically taken to mean just the ion rotation. Note that diagnostic limitations make the measurement of electron rotation much more difficult, but the charge density weighted sum of the ion and electron rotation is the plasma current density.

Rotation has been shown to provide significant benefits to tokamak plasmas: the stabilization of resistive wall modes and neoclassical tearing modes, and the suppression of turbulence through $E \times B$ shear. The effect of $E \times B$ shear in particular is important for understanding phenomena important to tokamaks, such as the formation of transport barriers, and for gyrokinetic modeling of tokamak
transport. Rotation is important for determining the $\mathbf{E} \times \mathbf{B}$ shear, which is done by using force balance to measure the radial electric field according to

$$E_r = \frac{nZ_e \nabla P}{e} - V_\varphi B_\theta + V_\theta B_\varphi,$$

and then calculating the $\mathbf{E} \times \mathbf{B}$ shearing rate as

$$\omega_{\mathbf{E} \times \mathbf{B}} = \frac{R^2 B_\theta^2}{B} \frac{\partial}{\partial \psi} \left( \frac{E_r}{RB_\theta} \right).$$

Measurements of rotation along with measurements of the magnetic field (from, for instance, a magnetic equilibrium reconstruction) are needed to calculate the $\mathbf{V} \times \mathbf{B}$ force present in Eq. 1.3. With the addition of measurements of temperature and density, the radial electric field and its associated shearing rate can be determined.

### 1.4 The study of poloidal rotation

This dissertation concerns the measurement of impurity poloidal rotation. The goal of studying impurity poloidal rotation is to gain a predictive understanding of tokamak plasma rotation, which is important for developing fusion power. Poloidal rotation is typically much lower than toroidal rotation in current tokamaks because of the presence of toroidal neutral beam injection and neoclassical poloidal damping (friction with trapped particles and magnetic pumping slow poloidal rotation, see Chapter 2). However, poloidal rotation is still important to study. This is because, while poloidal rotation is smaller than toroidal rotation, it is multiplied by the toroidal field in Eq. 1.3 while toroidal rotation is multiplied by the smaller poloidal field ($B_\varphi/B_\theta \approx 4 - 10$ in DIII-D). In addition, toroidal rotation is expected to be much smaller in reactor-sized tokamaks (e.g. ITER) in which case the radial electric field will be increasingly determined by the poloidal rotation. Future tokamaks are expected to have low toroidal rotation because the amount of neutral beam torque will not scale with machine size as fast as the plasma’s moment of inertia.

In order to develop a predictive understanding of poloidal rotation, measurements must be compared with theoretical predictions. This dissertation
works on both halves of this comparison by developing a more accurate poloidal rotation diagnostic and also comparing measurements to more advanced theoretical predictions. The new diagnostic was developed through a combination of expanding the DIII-D CER system and developing new innovative analysis techniques.

Studying the validity of the neoclassical formulation of poloidal rotation is important for developing a predictive understanding of rotation, and it could also be important for validating the portion of neoclassical theory that predicts the bootstrap current. Bootstrap current is a fundamental element of all fusion reactor designs, but difficult to diagnose directly. The relation between the bootstrap current and the poloidal rotation in neoclassical theory is due to dependence of both on the viscous forces that arise between trapped and passing particles. In addition, an accurate theory of poloidal rotation within neoclassical theory is required to properly calculate impurity asymmetries within a flux surface.

A significant amount of previous work has compared poloidal rotation measurements to neoclassical theory. Notable disagreements have been seen in DIII-D, JET, and TFTR. Agreement has been seen in TCV, NSTX, MAST, and DIII-D with the latter DIII-D results suggesting that turbulence could be important for driving poloidal rotation. Mixed results have also been observed in DIII-D and ASDEX Upgrade. Though poloidal rotation plays an important role in understanding the edge transport barrier seen in H-mode plasmas (accompanied by a significant amount of literature), the work of this dissertation is focused in the core of the plasma where large discrepancies have been observed and remain unexplained. As explained in Chapter 3, the core is also the measurement region of the newly developed poloidal rotation diagnostic.

1.5 Charge exchange recombination spectroscopy

Ion rotation in tokamaks is typically measured with charge exchange recombination spectroscopy (CER, though many alternate acronyms exist). CER
is a measurement method where atomic emission from ions in the plasma that have
gained an electron via charge exchange with a neutral particle is spectroscopically
analyzed. An example of this process, typical for DIII-D where the ion under
investigation is fully stripped carbon (an intrinsic impurity) and the charge exchange
electron is provided by a neutral deuterium is

\[
D^0 + C^{+6} \rightarrow D^{+1} + C^{+5} + h\nu \quad (n = 8 \rightarrow 7,5291 \ \text{Å})
\]

The Doppler shift of the emission provides information about the velocity
distribution of the fully stripped carbon in the plasma. In addition, the Doppler
broadening and total intensity measure the temperature and density, respectively,
of the fully stripped carbon. Therefore, CER can measure all of the ion properties
in Eq. 1.3. Measurements of emission from fully stripped light impurities are typical
for CER because all light impurities are fully stripped in the core of a hot tokamak
plasma, and the atomic physics governing their emission after charge exchange is
simplified by the atom being in a hydrogen-like state.

The DIII-D CER diagnostic has been operating since the first days of DIII-D
but undergone many upgrades and modifications since then. It is designed to
analyze charge exchange emission that results from fast, neutral beam injected
deuterium charge exchanging with impurity ions in the plasma. This has the
benefit of localizing the measurement location to where the diagnostic view chord
and the neutral beam cross. The angle the view chord makes with the toroidal and
poloidal direction at the point where it crosses the neutral beam determines the
type of rotation it can measure. Tangential view chords measure toroidal rotation
because they are horizontal and mostly aligned with the toroidal direction, and
vertical view chords measure poloidal rotation because they are vertical and mostly
aligned with the poloidal direction. The radial velocity of the plasma is taken to be
small. This is because toroidal and poloidal rotations values of interest are greater
than 1 km/s while a radial velocity one-hundredth this size cannot exist while the
plasma is confined.

A plan view of the DIII-D CER diagnostic is shown in Fig. 1.2. The system
consists of many view chords (64 in total) that observe multiple neutral beams
(including both co- and counter-current neutral beams) and a large portion of the
Figure 1.2: A plan view of the DIII-D CER diagnostic view chords showing the 6 neutral beams used by the CER diagnostic (two more neutral beams in the lower right quadrant are not used by CER and omitted). Standard $B_\phi$ and $I_p$ directions are also depicted. Note that co-current and counter-current neutral beam injection is possible. Tangential view chords added as part of the development of the new poloidal rotation diagnostic are colored green.

plasma. Most of the views are tangential in the midplane of the plasma, and the rest are vertical views that directly measure the poloidal rotation. The vertical views are most easily seen in a cross section view of the CER view chords, shown
in Fig. 1.3. In both views, tangential view chords that intersect co- and counter-current neutral beams inside the magnetic axis are highlighted because they were recently installed as part of the development of the new poloidal rotation diagnostic described in Chapter 3.

![Figure 1.3: A cross section view of the DIII-D CER diagnostic view chords on top of a magnetic equilibrium reconstruction. Tangential view chords added as part of the development of the new poloidal rotation diagnostic are colored green.](image)

**1.5.1 Impurity and main-ion CER**

The CER measurements in this dissertation are almost exclusively measurements of fully stripped carbon, an impurity in the plasma. Measurements of the
main-ion (typically deuterium in DIII-D) properties are more fundamental for describing tokamak plasmas though. CER measurements of the main-ion properties have recently been made robust, but these measurements require significantly more experimental hardware and complicated analysis. Fortunately then, measurements of impurity properties are sufficient for many physics investigations. The main reason for this is because CER can measure all the terms in Eq. 1.3 and therefore determine the radial electric field which is common to all species in the plasma.

Investigating the physics of poloidal rotation is also a task well suited to impurity CER. Direct measurements of main-ion poloidal rotation are not currently possible in hydrogenic plasmas with hydrogenic neutral beams. This is because the main-ion charge exchange spectra can only be resolved into its constituents when the emission from high speed beam-neutrals and fast-ions is separated from thermal main-ion charge exchange emission by a significant Doppler shift. This Doppler shift is not seen by vertical views which are nearly perpendicular to horizontal neutral beams (as is the standard in tokamaks). Impurity CER measurements of poloidal rotation are, by comparison, much easier, and they can be used to test neoclassical theory, which can then predict the main-ion poloidal rotation. Accordingly, calculations of the $E \times B$ shearing rate are almost always made with only impurity CER measurements.

Measurements of main-ion poloidal rotation has been made before, but in limited circumstances. Using a plasma composed mainly of helium greatly simplifies the main-ion charge exchange spectra, but this constraint adds experimental complexity and places significant limits on the types of plasmas that can be studied. Combining impurity measurements of $E_r$ with main-ion measurements of pressure and toroidal rotation allows the main-ion poloidal rotation to be determined from the uniqueness of $E_r$, but implicitly depends on measurements of impurity poloidal rotation. Measurements of main-ion poloidal rotation could be accomplished by employing the new measurement method that is explained in Section 3.3. A significant portion of this thesis concerns the development of this method for measurements of impurity poloidal rotation, but
the principles of the method could easily be applied to main-ion measurements.

1.5.2 Cross section effects

The cross section for charge exchange is energy dependent, and it is important to consider the effects of this when making CER measurements. For the C VI \( n = 8 \rightarrow 7 \) transition used for standard CER measurements at DIII-D, more energetic collisions have a larger cross section. As a result, the population of C\(^{+5}\) ions that are observed by CER are a weighted average of the underlying C\(^{+6}\) ions that CER seeks to observe. This process can be approximately thought of as being caused by the fact that carbon ions with velocities that are parallel to the stream of neutral beam particles are less likely to charge exchange with those neutrals than carbon ions that have anti-parallel velocities. As a result, CER measures an amount of C\(^{+5}\) rotation that is into the stream of neutral beam particles but not present in the underlying C\(^{+6}\) population.

This “cross section effect” has been known for some time\(^{23,24}\) and calculations can be used to correct for its effects.\(^{25}\) The cross section effect causes small inaccuracies in temperature measurements in DIII-D but significant inaccuracies in naive rotation measurements. Accounting for all the possible atomic physics that can influence the cross section effect for all plasma parameters is difficult and adds uncertainty to rotation measurements.

Since the DIII-D neutral beams used for CER measurements inject horizontally, it may appear that the cross section effect would not be present for vertical views, but this is not true. As was realized relatively recently,\(^{26}\) the finite excited state lifetime couples cross section effect from horizontal neutral beams into the vertical plane. This process is diagrammed in Fig. 1.4. The dependency of this “gyro-orbit cross section effect” on the excited state lifetime significantly complicates poloidal rotation measurements. This lifetime is \(\approx 1 \text{ ns}\),\(^{27}\) and the gyro-frequency for carbon is \(\approx 0.1 \text{ GHz}\). Accordingly, a carbon ion will execute one-tenth of a gyro-orbit after charge exchanging with a neutral before radiating. However, since magnitude of the cross section effect is usually much larger than the poloidal rotation, accurate accounting for the amount the cross section effect is
coupling into the vertical plane is required. A significant portion of the work in this dissertation is dedicated to making poloidal rotation measurements without the need for corrections derived from atomic physics calculations, so as to remove this uncertainty from the measurements. This process will be explained in Chapter 3.

![Diagram of cross section effect](image)

**Figure 1.4**: A diagram of the cross section effect coupling into the vertical plane as a result of the finite exited state lifetime, $\tau$.

### 1.6 Dissertation organization

This dissertation seeks to further the study of poloidal rotation by developing improved diagnostic techniques and testing more complex theoretical models. To this end, Chapter 2 summarizes the theoretical formulations of poloidal rotation that are tested. In addition to the standard neoclassical formulation of poloidal rotation, Reynolds stress and neutral beam driven poloidal rotation are discussed. Theoretical comparisons made in this dissertation often rely on comparison with computer simulations, and the different simulations are described in this chapter.

Chapter 3 contains a full description of the CER diagnostic at DIII-D, and describes recent modifications that were made as part of this dissertation. These modifications include a newly developed technique for correcting the astigmatism of a Czerny-Turner spectrometer. This chapter also contains the description of the
new poloidal rotation diagnostic that has been developed and the verification of its measurements. Comparisons of these results with the already existing poloidal rotation measurement technique are also presented. The newly developed poloidal rotation diagnostic is used to obtain the majority of the experimental results in this dissertation.

Chapter 4 presents a new method for performing the spatial calibration of neutral beam based diagnostics. This work was developed for the CER diagnostic to ensure the basic spatial calibration procedure produced accurate results. This new method is able to determine the accuracy of neutral beam positions via in-situ measurements of their geometry.

Chapter 5 concerns the measurement of poloidal rotation in ITB plasmas. The new poloidal rotation diagnostic has enabled the CER diagnostic to, for the first time on DIII-D, observe mean poloidal flow spin-up during ITB formation. Poloidal rotation inside the transport barrier is also investigated, and the consistency of these results with the already existing poloidal rotation measurement technique is discussed.

Chapter 6 presents various comparisons between measured and predicted poloidal rotation. Basic comparisons with neoclassical theory are shown throughout, and the possible effects of a turbulence induced Reynolds stress and fast-ion friction are investigated. The measurements show that a more complex description of poloidal rotation is needed in order to make accurate predictions, and it is necessary to consider the combination of many poloidal rotation drive mechanisms.

Chapter 7 contains a summary of the results as well as a discussion of their significance. Possible avenues for future work are presented as well.
Chapter 2

Poloidal rotation theory

This chapter outlines the basic results of theoretical calculations of poloidal rotation in a tokamak. Ultimately, a predictive understanding of poloidal rotation is sought, as this enables improved control and optimization of a tokamak plasma. This is particularly important for poloidal rotation because external control with momentum injection has not been implemented on any tokamak. Though toroidal rotation can be controlled with neutral beam injection, it is expected that future (large) tokamaks will have much less control of rotation because the plasma moment of inertia will scale with the tokamak volume (setting the total plasma mass) and size (setting the value of $R^2$) for a total scaling of $R^5$ while the momentum confinement time is only expected to scale as, approximately, $R^2$. This fact makes the verification of poloidal rotation theories increasingly important for the design and operation of future tokamaks.

A tokamak plasma is an extremely complex system involving charged particles of disparate masses, velocities, and temperatures influenced by static and varying electric and magnetic fields. In addition, important physical properties such as the collision frequency can span orders of magnitude within the same plasma. Some simplifications, such as taking collision frequencies to be much smaller than gyrofrequency ($\nu \ll \omega_c$), are possible and they allow the description to be simplified. Nevertheless, theoretical descriptions of the plasma are often formulated in specific regimes where one or more physical properties are taken to be a certain size. For instance, the gyro-radius is typically assumed to be small
compared to the gradient scale length, i.e. $\rho_g/L \ll 1$. This dissertation focuses on the core of DIII-D, which serves to reduce some of the complexity of theories. In the core of almost all DIII-D discharges, ion and electron temperature exceed 1 keV. This results in much smaller collision frequencies than those seen in the edge of the plasma. It also reduces the number of different species in the plasma by ensuring that light impurities (carbon is the main impurity in DIII-D) are fully ionized.

### 2.1 Basic plasma description

The plasma can be rigorously described by a set of Boltzmann equations, one for each species’ distribution function, $f_a(x, v, t)$,

$$\frac{\partial f_a}{\partial t} + v \cdot \nabla f_a + \frac{e_a}{m_a} (E + v \times B) \cdot \nabla_v f_a = C,$$

where $E$ and $B$ are the large scale fields, and $C$ represents the change in $f_a$ due to sources and collisions. With the distribution functions and Maxwell’s equations (to describe $E$ and $B$), the Boltzmann equations contain all the plasma dynamics. However, this description does not help reduce the complexity of the problem.

By taking moments of these equations, a fluid description can be built which is often sufficiently descriptive. To do this, the distribution function is first integrated over velocity space, yielding definitions for plasma properties that are functions of location and time, a description more suited to basic plasma experiments. These properties are the density, velocity, and temperature, given (respectively) by:

$$n_a(x, t) = \int f_a(x, v, t) d^3v,$$

$$V_a(x, t) = \int v f_a(x, v, t) d^3v,$$

$$T_a(x, t) = \frac{1}{3} \int m_a(v - V_a) f_a(x, v, t) d^3v.$$

For this section, it is convenient to define the $\langle \rangle$ operation as the normalized velocity space average over the distribution function, i.e. $\langle A \rangle \equiv \int A f d^3v/n$. 
Taking similar moments of Eq. (2.1) yields fluid equations for the evolution of these parameters: 

\[
\frac{Dn_a}{Dt} + n_a \nabla \cdot V_a = 0,
\]

(2.5)

\[
m_a n_a \frac{DV_a}{Dt} = -\nabla P_a - \nabla \cdot \pi_a + e_a n_a (E + V_a \times B) + F_a,
\]

(2.6)

\[
\frac{3}{2} n_a \frac{DT_a}{Dt} + p_a \nabla \cdot V_a = -\nabla \cdot q_a - \pi_a : \nabla V_a + Q_a,
\]

(2.7)

where \( D/Dt \equiv \partial/\partial t + V \cdot \nabla \) (known as the convective derivative), \( P = nT \), \( \pi_{jk} = mn \langle (v - V)_j (v - V)_k \rangle - P \delta_{jk} \), \( F_a \) is the total force due to collisions, \( q \equiv n \langle m|v - V|^2 (v - V)/2 \rangle \), \( \pi : \nabla V \equiv \pi_{jk} \partial V_k / \partial x_j \), and the energy flux components are \( Q_j = q_j + 5P V_j / 2 + \pi_{jk} V_k + mn V^2 V_j / 2 \). Though the dynamics of plasma temperature are not tested in this dissertation, they are included for completeness. Each moment equation contains a quantity from the next higher moment equation, so Eqs. (2.5)-(2.7) are not closed.

By taking the cross product of Eq. (2.6) with \( B \), an equation for the perpendicular flow can be obtained,

\[
V_\perp = \frac{E \times B}{B^2} + \frac{B \times \nabla P}{neB^2},
\]

(2.8)

where the \( a \) subscript that differentiated species has now been dropped. To arrive at this equation, plasma equilibrium was assumed (\( \partial/\partial t \approx 0 \)), rotation was taken to be small compared to the thermal speed (\( |V|/V_{th} \approx 0 \), \( V_{th} \equiv \sqrt{2T/m} \)), the distribution function was taken, to lowest order, to be a Maxwellian, and terms that are second order in \( \rho_g / L \), were discarded. This equations shows that the perpendicular velocity is a consequence of the \( E \times B \) and diamagnetic drifts.

Along with the assumption that flow within a flux surface is divergence free, i.e. \( \nabla \cdot (nV) = 0 \), it is possible to determine the parallel return flow that must accompany any perpendicular flow,

\[
nV_\parallel = \frac{RB_c n \omega(\psi)}{B} + k(\psi)B.
\]

(2.9)

Here, \( k \) is an unknown flux function and \( \omega \) is defined by

\[
\omega \equiv -\frac{d\Phi}{d\psi} - \frac{1}{nZe} \frac{dP}{d\psi}.
\]

(2.10)
This allows the form of the total flow, for any species, to be written as
\[ n\mathbf{V} = k(\psi)\mathbf{B} + nR\omega(\psi)\hat{\phi}. \quad (2.11) \]

Note that the assumption of divergence free flow is a result of assuming the time derivative to be second order in \( \rho_g/L \). Eq. (2.11) makes it clear that toroidal rotation is not a flux function, and this forms the basis for the new poloidal rotation diagnostic described in Chapter 3. Ultimately, the parallel projection of Eq. 2.6 is used to determine the poloidal rotation. In the next sections, such calculations will be described. Neoclassical theory (Sec. 2.2) is concerned with the value of \( \nabla \cdot \pi \), turbulence induced Reynolds stress (Sec. 2.3) is concerned with the value of \( \mathbf{V} \cdot \nabla \mathbf{V} \) in the convective derivative, and friction with fast-ions (Sec. 2.4) is concerned with the value of \( F \).

### 2.2 Neoclassical theory of poloidal rotation

Neoclassical theory is principally concerned with effect of a spatially varying magnetic field on classical treatments that considered the magnetic field to be uniform on the scale of a particle’s mean free path. In the core of a tokamak plasma, the mean free path is on the order of, or larger than, the machine size, which contains significant variation in the magnetic field. Though toroidal symmetry is fundamental to the tokamak design, the toroidal magnetic field scales inversely with the major radius. As will be discussed later in this section, this variation in magnetic field strength has two results that are very important for poloidal rotation dynamics: the creation of trapped particle orbits and magnetic pumping. Trapped particle orbits are due to the combination of magnetic mirroring (due to the \( 1/R \) dependence of \( B_\phi \)) and \( \nabla B \) drift. Similarly, the \( 1/R \) dependence of \( B_\phi \) causes magnetic pumping for particles that are not trapped, i.e. passing particles.

Neoclassical transport is separated into different regimes based on plasma collisionality, \( \nu^* \), the normalized collision frequency. This is because the collision frequency determines how particles are able to relax gradients. The typical
normalization factor is the trapped particle orbit frequency, yielding
\[ \nu^* = \frac{\nu_{ii} q R}{\epsilon^{3/2} V_{th,i}}, \]  
(2.12)

where \( q \) is the safety factor and \( \epsilon \) is the inverse aspect ratio (ratio of minor radius to major radius). In this way, the collisionality can differentiate between the regimes where trapped particle orbits complete and those where they do not. To the extent that \( \nu_{ii}/V_{th,i} \approx \nu_{ee}/V_{th,e} \) (typically true for electrons and light ions), this definition applies for ions and electrons.

Frequent collisions, \( \nu^* > \epsilon^{-3/2} \), result in particles being unable to complete trapped or circulating orbits (faster orbits of untrapped particles) before being scattered. This is referred to as the Pfirsch-Schuler regime. Infrequent collisions, \( \nu^* \ll 1 \), allow particles to complete trapped and circulating particle orbits before scattering. This is referred to as the the banana regime (due to the banana-like shape of trapped particle orbits). An intermediate regime is possible when the inverse aspect ratio is small, where particles can complete circulating particle orbits but not trapped particle orbits before scattering, \( 1 < \nu^* < \epsilon^{-3/2} \). This is referred to as the plateau regime.

Equation (2.9) can be rewritten without the unknown flux function \( k(\psi) \) as
\[ nV_{//} = -\frac{nR B \epsilon}{B} \left( \frac{d\Phi}{d\psi} + \frac{1}{nZ_e} \frac{dP}{d\psi} \right) \left( 1 - \frac{B^2}{\langle B^2 \rangle} \right) + \frac{nB \langle BV_{//} \rangle}{\langle B^2 \rangle}. \]  
(2.13)

Here (and for the remainder of this dissertation), \( \langle \rangle \) denotes the flux-surface average,
\[ \langle A \rangle \equiv \frac{\oint A dl_p/B_\theta}{\oint dl_p/B_\theta}. \]  
(2.14)

The value of \( \langle BV_{//} \rangle \) will depend on the ability of particles to complete their orbits, so Eq. 2.13 makes it clear that neoclassical calculations of poloidal rotation will depend on the collisionality. In this dissertation, the focus on core measurements means the collisionality is generally in the banana regime.

Calculations of poloidal rotation for the main-ion and impurity can be made via the moment approach of Hirshman and Sigmar. In this approach, the balance between parallel friction and viscous stress and the equivalent balance for heat friction and heat viscous stress can be used to calculate the poloidal flow. Given
the form of the rotation in Eq. 2.11, and that a similar equation can be written for the heat flux, the viscous stresses are only functions of the poloidal flow. This is because the second term in Eq. 2.11 only produces rigid rotor type motion, which creates no stress.

These calculations were explicitly performed in the trace impurity limit in Ref. 31, and the results are shown here:

\[ V_{\theta,i} = \frac{1}{2} \frac{V_{th,i} \rho_{g,i}}{L_{Ti}} \frac{K_1 BB_\varphi}{\langle B^2 \rangle}, \]  

(2.15)

\[ V_{\theta,I} = \frac{1}{2} \frac{V_{th,i} \rho_{g,i}}{L_{Ti}} \left[ \left( \frac{K_1}{2} + \frac{3K_2}{2} \right) \frac{1}{L_{Ti}} - \frac{1}{L_{Pi}} + \frac{Z_i T_i}{Z_i T_i L_{Pi_i}} \right] \frac{BB_\varphi}{\langle B^2 \rangle}. \]  

(2.16)

In these calculations, \( i \) denotes the main-ion, \( I \) denotes the impurity, \( \rho_g \) is the gyroradius \( (V_{th}/\omega_c) \), \( L_A \) is the gradient scale length \( (\equiv d\ln A/dr) \), and \( K_1 \) and \( K_2 \) are functions of \( \nu^* \) and \( \epsilon \). Typically, all ion temperatures are taken to be the same, due to good collisional coupling between them. Equations (2.15) and (2.16) show that, while the main-ion poloidal rotation scales with the temperature gradient, the impurity poloidal rotation is a more complicated function of the temperature and pressure gradients. For typical DIII-D parameters, the values of \( K_1 \) and \( K_2 \) ensure that the temperature gradient drive of impurity poloidal rotation is opposed by the pressure gradient drive.

Equations (2.15) and (2.16) were derived under the assumption that \( V_{th,I}/V_{th,i} \ll 1 \), for convenience. For DIII-D cases presented in this dissertation, though, \( V_{th,I}/V_{th,i} \approx 0.4 \) (deuterium main-ion, carbon impurity). Nevertheless, these equations are useful for showing the important dependencies of the neoclassical calculation of poloidal rotation. More extensive calculations (discussed in Section 2.2.1) can be made with the moment approach, extending these results to light impurities and plasmas with more than one impurity.

### 2.2.1 Neoclassical simulations

The moment approach forms the basis of the NCLASS simulation code\textsuperscript{32} which is able to simulate multiple light and heavy impurities with continuous coverage across collisionality regimes. As in Ref. 31, rotation is assumed to be
low compared to the thermal speed, but the ratio of the main-ion and impurity thermal speeds is not constrained. The moment approach in unable to determine the absolute toroidal rotation, only the difference between two species. As a result, NCLASS simulations require a measurement of the toroidal rotation to obtain the radial electric field.

On DIII-D, NCLASS is run with front-end code FORCEBAL\textsuperscript{33} where the carbon toroidal rotation is provided as an input. In the NCLASS formalism, poloidal rotation is the result of parallel velocity and thermodynamic forces that arise from gradients in the plasma. This flow generates friction and viscous forces that are balanced against each other so that the poloidal component of the parallel velocity can be calculated. Using friction and viscosity coefficients that depend on density, temperature, charge, and mass, the poloidal rotation for the main-ion and impurity can be calculated. The result of these calculations can be compared with measurements of impurity poloidal rotation as a way to verify the neoclassical theory of poloidal rotation.

Other methods of determining the poloidal rotation within a neoclassical framework exist. For instance, the NEO simulation code\textsuperscript{34} determines the distribution function by solving the drift kinetic equation after it is expanded in powers of $\rho^*$, the ratio of the gyroradius to the system size. The drift kinetic equation describes how collisions change the distribution function along particle trajectories. This $\delta f$ procedure can simulate plasmas with multiple impurities across all collisionality regimes and account for large rotation and plasma shaping effects.

Typically, the poloidal rotation calculations made by NEO are quite similar to those made by FORCEBAL. However, due to the fact that the physics included in the NEO simulation code is more complete than the physics contained in FORCEBAL, NEO results are the results compared to poloidal rotation measurements to determine the accuracy of the neoclassical theory of poloidal rotation. Unfortunately, the additional complexity within NEO makes the time to setup and complete its calculations significantly longer than time needed to acquire FORCEBAL results. For this reason, estimating the error in neoclassical
calculations that result from experimental uncertainties in the measurement inputs is more practically done with FORCEBAL.

In order to estimate this error, the plasma profile inputs to the FORCEBAL code (electron density and temperature, ion density and temperature, impurity toroidal rotation as a function of radius) are randomized according to the experimental errors, and the FORCEBAL calculations based on these profiles are used to determine the error in the FORCEBAL results in a Monte Carlo fashion. In this dissertation, comparisons between calculations of poloidal rotation made with NEO and experimental measurements are used to assess the accuracy of the calculations, and the error in the neoclassical calculations are estimated with this Monte Carlo method.

For both NCLASS and NEO (and neoclassical simulations in general), it is necessary to restrict the calculations to the region where transport is local (see discussion in Ref. 34). Transport can become non-local when $\rho_g/L = O(1)$. This condition is essentially guaranteed in the core where potato orbits (trapped particle orbits with large widths) are present. For this reasons, neoclassical predictions are not shown in the deep core in this dissertation.

### 2.2.2 Accuracy of neoclassical theory

Experiments have shown that neoclassical theory does not contain a complete description of a tokamak plasma. Heat, particle, and momentum transport are frequently larger than neoclassical levels. Discrepancies are usually attributed to the lack of accounting for turbulence in the neoclassical formalism. In the literature, instances where heat and particle transport are as low as neoclassical levels are notable. In particular, in plasmas with reversed magnetic shear (safety factor decreasing with radius) and/or internal transport barriers (ITBs), heat and particle transport have been observed at neoclassical levels. Neoclassical predictions have also been important for understanding plasma resistivity and the self-generated bootstrap current.

Neoclassical predictions of poloidal rotation have shown mixed results. Though the inaccuracies of neoclassical theory can be attributed to the presence
of turbulence, there is reason to believe that the neoclassical theory of poloidal rotation could still be accurate. Fundamentally, this is due to the viscous damping of poloidal flow that serves to drive poloidal rotation back to its neoclassical value.

Physically, neoclassical poloidal damping is due to a combination of two viscous forces: magnetic pumping and friction with trapped particles. Both of these effects are related to the $1/R$ dependence of $B_\varphi$. Poloidally rotating particles lose parallel velocity due to magnetic pumping because of the variation in $B_\varphi$ they experience as they traverse the plasma poloidally. Trapped particles are the result of the variation in $B_\varphi$, and they drag the poloidal rotation of untrapped particles to their poloidal rotation, which is, on average, zero.

Magnetic pumping is more significant at lower collisionality, while friction with trapped particles is more significant at higher collisionality. Ultimately, the damping scales with collisionality because both of these viscous forces are the result of collisions. Though other effects can be theorized to drive poloidal rotation (two such examples are discussed in Sections 2.3 and 2.4), if the neoclassical poloidal damping is sufficiently strong, the poloidal rotation will quickly return to neoclassical levels. Expectations that the damping is strong have motivated continued study of neoclassical poloidal rotation theory even in cases where most every other portion of the theory (transport of heat, particles, momentum) is known to be inaccurate.

Formally, the damping is caused by the neoclassical viscosity coefficients that are used to determine the viscous stress that results from poloidal flow. As in Ref. 40, the parallel viscosity can be written as

$$\langle B \cdot \nabla \cdot \pi_{\text{neo}} \rangle = mn\langle B^2 \rangle \mu_k \mu_{\text{neo}},$$

(2.17)

where $\pi_{\text{neo}}$ is the neoclassical viscous stress tensor, $\mu$ represents the viscous damping, and $k_{\text{neo}}$ represents a flux-function that can be used to calculate the poloidal rotation that results from the standard neoclassical treatment. Any additional poloidal rotation present that did not have an origin in the standard neoclassical treatment can be represented in a similar equation,

$$\langle B \cdot \nabla \cdot \pi \rangle = mn\langle B^2 \rangle \mu (k_{\text{anom}} - k_{\text{neo}}),$$

(2.18)
where $k_{\text{anom}}$ represents what is, for now, referred to as the anomalous poloidal rotation.

The essence of Eqs. (2.17) and (2.18) is that poloidal rotation, regardless of origin, creates viscous stress. If anomalous poloidal rotation is hypothesized, Eq. (2.18) can be used in the equation for momentum conservation [Eq. (2.6)], but any drive term will necessarily be reduced by the neoclassical viscous damping coefficient.

### 2.3 Residual Reynolds stress drive

Since discrepancies between neoclassical theory and experimental measurements are often attributed to the presence of turbulence in the plasma, it is logical to consider the effect that turbulence would have on the poloidal rotation. Turbulence can affect momentum balance in the plasma due to a Reynolds stress, which is present in Eq. (2.6) in the advective portion of $D\mathbf{V}/Dt$. Turbulence can create fluctuations in $\mathbf{V}$ that then create a “residual” Reynolds stress which then affects the mean velocity.

Working from a flux surface average of Eq. (2.6) projected onto the parallel direction, a turbulence induced Reynolds stress can be seen to drive poloidal rotation according to

$$mn\langle B^2 \rangle \mu (k_{\text{NBI}} - k_{\text{neo}}) = \langle mn \mathbf{B} \cdot (\tilde{\mathbf{V}} \cdot \nabla \tilde{\mathbf{V}}) \rangle,$$

(2.19)

where $\mathbf{V}$ has been separated into static and fluctuating components, and $\langle \tilde{\mathbf{V}} \rangle \equiv 0$. In order for there to be a finite Reynolds stress drive, there must be some symmetry breaking in the fluctuations. This can be seen clearly in a different formulation of the turbulence induced Reynolds stress,

$$\mu \langle V_\theta \rangle = -\frac{\partial}{\partial r} \langle \tilde{V}_r \tilde{V}_\theta \rangle.$$

(2.20)

This symmetry breaking can be caused by variation in plasma parameters that cause the turbulence intensity to have a radial variation on the order of the size of the fluctuations. As shown by Diamond, this effect can be formulated as
A basic estimation of the size of this effect can be made by determining the poloidal flow driven by turbulence induced radial current via the $\mathbf{J} \times \mathbf{B}$ force,\textsuperscript{41}

$$mn\mu B_0^2 \langle V_\theta \rangle = -J_r B_\theta,$$  \hspace{1cm} (2.21)

where $B_0$ is the field at the magnetic axis and $J_r$ is the radial current density created by the turbulence. In Ref. 41, mixing length estimates were used to determine $J_r$ and Eq. (2.21) yielded large poloidal rotations, up to 30 km/s, in the edge.

In the core of the plasma, larger gradient scale lengths reduce $J_r$ by approximately two orders of magnitude, but lower collisionality reduces $\mu$ by an order of magnitude. Such changes will reduce the turbulence driven poloidal rotation by an order of magnitude, but even 1 km/s is important for the comparisons made in this dissertation. Though initial estimates of this effect made with GYRO\textsuperscript{43} were low,\textsuperscript{44} more recent results corroborate these basic estimations and suggest that turbulence induced Reynolds stress could be a significant drive of poloidal rotation.\textsuperscript{45}

### 2.4 Neutral beam drive

Fast-ions injected by the neutral beams on DIII-D can be highly directional. Another possible mechanism for driving poloidal rotation is then the friction between the thermalized impurities and the neutral beam injected fast-ions. Typically, neutral beams are injected in the tangential direction, so the momentum they carry cannot directly drive poloidal rotation. However, their motion can be parsed as being partially parallel and partially perpendicular to the magnetic field lines, and a difference in viscosity between these two directions would allow the friction between thermalized particles and fast-ions to have a poloidal component. Friction between trapped and passing fast-ions could also be important. Due to the non-Maxwellian nature of the fast-ion distribution function, the proportion of trapped and untrapped fast-ions can be very different from the thermalized ions. Note that fast-ions that result from NBI are focused on in this section and this
dissertation, but the principles outlined could be applied to fast-ion populations created through other means.

The effect of fast-ion friction on momentum conservation is to add a source of momentum to the flux-surface averaged Eq. (2.6) projected onto the parallel direction,

\[ mn(B^2)\mu(k_{\text{NBI}} - k_{\text{neo}}) = \langle B \cdot F_{\text{fast}} \rangle, \]

(2.22)

where \( F_{\text{fast}} \) is the friction from the neutral beam injected fast-ions, and it is assumed the plasma is in equilibrium with no Reynolds stress. As noted before, this additional drive for poloidal rotation will be reduced by the neoclassical viscous damping. Though this mechanism for driving poloidal rotation is not outside the scope of the neoclassical formalism, no current neoclassical simulations calculate this effect. That is, no neoclassical simulations include a separate fast-ion species and determine the interspecies friction when solving parallel force balance.

Work by Newton\(^{46}\) presents a formalism for calculating this effect. In addition, a method for estimating the ratio of fast-ion induced poloidal rotation, \( V_{\theta,\text{NBI}} \), to standard neoclassical poloidal rotation, \( V_{\theta,\text{neo}} \), is provided:

\[ \frac{V_{\theta,\text{NBI}}}{V_{\theta,\text{neo}}} \approx \frac{n_f V_{\text{th},i}}{n_i V_{\text{th},f}} \frac{L}{\rho_{f}^{\text{ban}}}. \]

(2.23)

Here, \( f \) and \( i \) denote the fast and thermalized ions, respectively, \( L \) is the radial scale length, and \( \rho_f^{\text{ban}} \) is the fast-ion banana orbit width. Even when estimating the parameters in Eq. (2.23) so that the potential for fast-ion driven poloidal rotation is minimal, the result of this calculation suggests that this effect is significant. Taking \( L/\rho_f^{\text{ban}} \) to be 0.5, the equivalent temperature of the fast-ions to be 40 keV, and the plasma temperature to be 10 keV, \( V_{\theta,\text{neo}} \) will be only twice as large as \( V_{\theta,\text{NBI}} \) with a 10% fast-ion content. All these parameters are reasonable for a DIII-D discharge with large NBI power.

Note that Eq. (2.23) is based on the assumption that the fast-ions have a net rotation. The calculations in Ref. 46 yield no fast-ion friction when the net rotation of the fast-ions is zero, by construction. This is a relevant point because on DIII-D it is possible, with co-current and counter-current NBI, to create significant fast-ion populations that have no net rotation.
Assuming a basic model of fast-ion friction, where both impurities and main-ions are forced along the field line, the direction of poloidal rotation driven by fast-ion friction can be determined. In DIII-D, the standard toroidal field direction is opposite the standard current direction, which is the injection direction for most of the neutral beams (See Fig. 1.2). As a result, for the standard DIII-D configuration, poloidal rotation is driven in the electron diamagnetic direction. Reversing the direction of the plasma current or the neutral beam torque will switch the direction of this drive to the ion diamagnetic direction. If the direction of the toroidal field is switched, the physical direction of the fast-ion driven poloidal rotation will not change, but its correspondence with the ion or electron diamagnetic direction will.

2.5 Conclusions

Neoclassical theory can predict poloidal rotation for impurities in a tokamak plasma. However, previous results have suggested that the neoclassical theory of poloidal rotation is inaccurate. Additional drives for poloidal rotation, such as a turbulence induced Reynolds stress or friction with fast-ions, will be reduced by the neoclassical viscous damping. This is true for any source of momentum that can be added into the momentum conservation equation. The neoclassical viscous damping scales with collisionality, so it is expected that plasmas with low collisionality will see stronger poloidal rotations that are not predicted by neoclassical simulations. The basic theories of neoclassical, turbulence induced Reynolds stress, and fast-ion induced poloidal rotation discussed in this chapter are tested against improved measurements of poloidal rotation in Chapter 6.
Chapter 3

Diagnostic methods

A large portion of the results in this dissertation were enabled by the development of a new poloidal rotation diagnostic. This diagnostic is based on charge exchange recombination spectroscopy (CER) and was made possible by a combination of significant upgrades to the CER system and innovative analysis. This chapter explains pertinent details of the CER diagnostic, the upgrades needed to enable this new poloidal rotation diagnostic, and the principles of the new poloidal rotation measurements.

3.1 CER on DIII-D

The DIII-D CER system has been upgraded many times since its initial construction in 1985. This section discusses important details of the diagnostic as it has existed since 2010. Flexibility has been a constant priority in the design of the CER system. As a result, wavelength scanning spectrometers are used. This allows emission to be collected throughout the visible spectrum so different atomic transitions from different ions can be analyzed with quick, between discharge tuning. In addition, charge coupled device (CCD) cameras are used that have adjustable exposure times and can take continuous data. This allows the CER system to easily cover an entire plasma discharge with continuous 5 ms exposures or focus on a small portion of the discharge with timing as fast as 0.204 ms.

To accomplish this, twelve scanning, Czerny-Turner type spectrometers\textsuperscript{47}
are used to image sixty-four view chords. These images are detected by CCD cameras that combine properties of high quantum efficiency, low read noise, and high speed. Forty tangential and twenty-four vertical view chords make up the CER system, and coverage of the outboard midplane is quite detailed. View chords in the core are separated by $\approx 4$ cm while view chords in a high resolution edge array are separated by $\approx 0.75$ cm, enabling detailed imaging of edge features such as the edge transport barrier. See Figs. 1.2 and 1.3.

The CER measurements in this dissertation are fairly standard and consist mostly of measurements of the carbon properties via observation of the C VI ($n = 8 \rightarrow 7, 5290.5 \text{ Å}$) emission measured with 2.5 to 5.0 ms exposures. Still, the ability to change wavelength and timing was imperative to various calibrations necessary for performing accurate CER analysis. This section explains the experimental apparatus and method used to collect and analyze CER data. The important calibrations and methods for correcting the effects of the energy dependent charge exchange cross section are also discussed.

### 3.1.1 Optics

Fiber optics are used to transmit light from the tokamak to a laboratory space separate from the region immediately surrounding the tokamak where they are coupled to Czerny-Turner spectrometers. The wavelength resolved light is imaged with CCD cameras. Fused silica fiber optics are used, and these fibers are mounted outside the vacuum vessel at diagnostic ports where lenses are used to create an image of the fiber tip on the desired neutral beam inside the tokamak. The lenses and fibers are separated from the vacuum by a sapphire window, and inside this window (for six of the eight ports used by the CER system), mirrors are used to direct the fiber images to their desired focus locations. The mirrors are coated with Beral\textsuperscript{TM}, a hard metallic coating with reflective properties similar to aluminum that was developed by H. L. Clausing Incorporated. Shutters protect the mirrors and windows from exposure to the glow discharges (low pressure helium plasmas used to clean the walls between each plasma discharge). The fiber optics allow tokamak light to travel to a laboratory space separate from the region.
immediately surrounding the tokamak where they are coupled to Czerny-Turner spectrometers. The wavelength resolved light is imaged with CCD cameras.

There are two type of fiber optics used for the CER system: 1.5 mm core fibers and 0.75 mm core pairs. The 1.5 mm core fibers are used in the core where their larger size increases signal levels, and the 0.75 mm pairs are use in the edge of the plasma where spatial resolution is important. Here, the pairs can be oriented in order to maximize their spatial resolution, i.e. tangential fiber pairs are stacked vertically and vertical fiber pairs are stacked toroidally (see inset diagram in Fig. 1.2). In this way, the two fibers that make up the pair span space with low spatial variation, and separate view chords can be placed as close to one another as possible in the direction that has high spatial variation.

Ten of the twelve spectrometers used are Acton Research Corporation 2/3 m f/4.7 units with toroidal mirrors that correct for the astigmatism inherent in the Czerny-Turner design. The other two spectrometers are Acton Research Corporation 1/2 m f/4.1 units that were not designed with any astigmatism correction. Note that the fiber optics and lenses are \( \approx f/2 \) so it is the spectrometer optics that set the \( f \)-number of the CER system. Astigmatism correction is discussed more in Section 3.2.2. Slit widths range from 0.1 to 0.3 mm and are chosen as a compromise between signal level and wavelength resolution. Eight spectrometers each use two Pixelvision CCD cameras to image four view chords, and four spectrometers each use one Sarnoff CCD camera to image eight view chords. The specifications of these two cameras are shown in Table 3.1.

**Table 3.1:** Specifications of the two types of CCD cameras used to acquire spectroscopic data for the CER diagnostic.

<table>
<thead>
<tr>
<th></th>
<th>Pixelvision Pluto</th>
<th>Sarnoff CAM1M100</th>
</tr>
</thead>
<tbody>
<tr>
<td>Chip size (pixels)</td>
<td>326 × 488</td>
<td>1024 × 1024</td>
</tr>
<tr>
<td>Pixel size (( \mu \text{m} ))</td>
<td>24 × 12</td>
<td>16 × 16</td>
</tr>
<tr>
<td>Quantum efficiency (%)</td>
<td>80</td>
<td>75</td>
</tr>
<tr>
<td>Minimum exposure (ms)</td>
<td>0.272</td>
<td>0.204</td>
</tr>
<tr>
<td>Read noise (e−)</td>
<td>25</td>
<td>30</td>
</tr>
<tr>
<td>Serial read rate (MHz)</td>
<td>2.2</td>
<td>1.6</td>
</tr>
<tr>
<td>Row transfer time (( \mu \text{s} ))</td>
<td>0.6</td>
<td>1.0</td>
</tr>
</tbody>
</table>
In order to observe eight view chords with the Sarnoff cameras, multiple, parallel slits are used on the spectrometer. This allows the fibers to be arranged in a square pattern that best utilizes the square chip in the Sarnoff camera. In order to prevent cross talk between view chords, a mask system was developed which is discussed in Section 3.2.3.

The spectrometers that image the edge CER views have dispersions of \( \approx 0.16 \, \text{Å/pixel} \), and the spectrometers that image the core CER views have dispersions from 0.2 Å/pixel to 0.25 Å/pixel. Note that the Pixelvision cameras are run with on-chip binning so that each effective pixel is actually the sum of two physical pixels on the chip. For simplicity, these effective pixels are referred to as simply pixels. The Pixelvision chip is 652 \( 12 \, \mu\text{m} \) physical pixels wide, but the effective width is 326 pixels in the simplified terminology used for this dissertation. No such binning is done for the Sarnoff cameras.

These cameras are cooled to \(-30^\circ\text{C}\) with thermoelectric coolers to minimize their dark current. At these temperatures, dark current is negligible, even with exposure times of 300 ms. The heat from the thermoelectrics is removed with chilled water (12°C), and nitrogen purges are used when the chip is open to air (to prevent condensation). The Sarnoff cameras are designed to use a nitrogen purge, while the Pixelvision cameras contain their chips in a small vacuum enclosure. For some of the Pixelvision cameras, though, the vacuum seal has degraded over time and necessitated the installation of a custom nitrogen purge. The necessary nitrogen flow is very low and the purges have not reduced the ability of the chips to maintain their low operating temperatures.

### 3.1.2 Analysis

CER data from a plasma discharge is separated into a number of timeslices for each view chord. Each timeslice represents a single camera exposure. When a neutral beam seen by a view chord is on during a timeslice, part of the detected signal will be from beam induced charge exchange. Additional signal from any point along the view chord will also be detected, e.g. visible bremsstrahlung, electron impact excitation from cold ions in the edge, and charge exchange of
the desired ions in the edge of the plasma that are within reach of neutrals coming from outside the plasma boundary. Plume emission is a potential source of charge exchange signal that would not necessarily be eliminated by timeslice subtraction. On DIII-D, the plume emission has been determined to be insignificant with investigations that used used timeslice subtraction to analyze signal from neutral beams seen by view chords that do not intersect that neutral beam. In general, CER analysis is the process by which the charge exchange emission from the neutral beam is fit with a model whose fit parameters can be related to the plasma parameters: temperature, velocity and density.

Though it is possible to develop a fitting model that accounts for all these sources of signal, the different signals can be well blended, making the resulting analysis more uncertain. For this reason, the highest quality DIII-D analysis uses background subtraction to isolate the charge exchange emission associated with the neutral beam. Background subtraction is performed by taking signal from a timeslice when a viewed neutral beam is on and subtracting signal from a nearby time when that neutral beam is off. If the conditions of the plasma, aside from the state of the viewed neutral beam, are relatively constant in both of these timeslices, the resulting signal will only contain charge exchange emission associated with the viewed neutral beam. Figure 3.1 shows an example of the high contrast images that can be obtained with timeslice subtraction.

Ensuring that the proper timeslice subtraction is available can be challenging due to the various beams viewed by all the CER view chords. As part of the work for this dissertation, it was often desirable to have the ability to perform independent timeslice subtraction for views of the 30 degree, 210 degree, and 330 degree neutral beams. Making these measurements while also keeping the plasma in steady state required considerable effort. This required keeping the total neutral beam input torque constant and the input power nearly constant. In hot plasmas, the slowing down time for the fast-ions that result from neutral beam injection is on the order of 100 ms while the time scale for neutral beam modulation is 10 ms. For this reason, the plasma can be thought of as a low-pass filter for neutral beam heat. However, effects on the plasma rotation can occur much faster
Figure 3.1: A charge exchange spectrum measured by the CER system during a single, 2.5 ms camera exposure. The fit is made to just the active portion. Note that this data comes from one of the view chords added in 2011 (see Sec. 3.2) that is on the high-field side, where neutral beam attenuation is greater. As a result, the signal levels seen here are lower than for typical view chords.

than the modulation time, making it important to maintain constant neutral beam torque when possible.

CER spectra are fit under the assumption that the ion distribution function is Maxwellian. The quality of the resulting fits justify this assumption. To fit a distribution function in a spectrum seen by a view chord, the instrumental response for that view chord is convolved with a Gaussian line shape. The instrumental response is itself a sum of Gaussians (finding those Gaussians is covered in Section 3.1.3.3),

$$I_i = \sum_{j=1}^{n} a_j \exp \left( \frac{-(i - p_j)^2}{w_j^2} \right), \quad (3.1)$$

where $n$ Gaussians with amplitudes $a$, locations $p$, and widths $w$ are used to model the intensity $I$ at each pixel, $i$. Future calculations can be simplified if normalized amplitudes, $b_j$, are used so that $\sqrt{\pi} \sum_{j=1}^{n} b_j w_j \equiv 1$. The ability of a sum of Gaussians to fit complex line shapes is excellent, and the use of Gaussians
simplifies the fitting procedure because the convolution of two Gaussians is another Gaussian. CER spectra are then fit according to

\[ I_i = A \sum_{j=1}^{n} \frac{b_j w_j}{W^2 + w_j^2} \exp \left( \frac{-(i - P)^2}{W^2 + w_j^2} \right), \]  

(3.2)

where the intensity at each pixel in a spectrum, \( I_i \), is used to determine the amplitude \( A \), the width \( W \), and the location \( P \) of the observed distribution function.

The speed of the impurity along a view chord’s sightline, \( \hat{s} \) (which points from the neutral beam to the fiber optic), can be determined from the calculation of the Doppler shift,

\[ \mathbf{V} \cdot \hat{s} = -\frac{D(P - p_0)c}{\lambda_0}, \]  

(3.3)

where \( D \) is the dispersion of the view chord, \( c \) is the speed of light, \( \lambda_0 \) is the rest wavelength of the observed atomic transition, and \( p_0 \) is the pixel location of this transition. The temperature can be determined from the calculation of the Doppler broadening,

\[ T = \frac{mc^2D^2W^2}{2\lambda_0^2}, \]  

(3.4)

where \( m \) is the ion mass. Zeeman splitting is not considered in the standard analysis of CER data. This is because its effects are only expected to be important for \( T_i < 100 \) eV, a temperature seen only in a small region of the edge for typical DIII-D discharges. Lastly, the amplitude can be related to the total number of photons measured per timeslice, which can then be related to the impurity density, given information about the neutral beam density, view chord geometry, and the cross section for the charge exchange reaction. This relationship is

\[ \frac{A}{\Delta t} 4\pi f = \sum_{j=1}^{2} \sum_{k=1}^{3} \int_{\Delta L} q_{CX}^{j,k} n_b^{j,k} n_c dl. \]  

(3.5)

Where \( \Delta t \) is the timeslice duration, \( f \) is a calibration factor that converts the camera digitizer count rate to photons/s/sr/m^2, \( \Delta L \) is the view chord’s path length through the neutral beam, \( n_c \) is the impurity density, \( n_b \) is the neutral beam density, \( q_{CX} \) is the rate coefficient of the charge exchange reaction. This calculation is complicated by the presence of multiple energy species in the neutral
beam \((k)\) and different excited states of the electron on the neutral beam particle \((j)\). The values of \(D, p_0, f,\) and \(\Delta L\) are determined with calibrations, as explained in the Section 3.1.3.

### 3.1.3 Calibrations

CER calibrations can be divided into three major groups: vent calibrations, between discharge calibrations, and beam into gas shot calibrations. Vent calibrations are done between experimental campaigns and require entering the vacuum vessel to measure the spatial position of CER view chords and determine each view chord’s absolute intensity response. Between discharge calibrations are largely automated and correct for variation in white light response and wavelength fiducial throughout an experimental day. Beam into gas shot calibrations can be done almost anytime during an experimental campaign, but are typically done only once, at the beginning, with less than ten shots being needed.

#### 3.1.3.1 Vent calibrations

Before and after each experimental campaign, a spatial and intensity calibration is done for the CER system. The measurements from before the campaign are used throughout the campaign itself, and the measurements at the end of the campaign can be used to re-calibrate measurements whose conditions significantly changed during the campaign. No significant changes in the spatial calibration are expected after an experimental campaign. The intensity calibration changes throughout the year due to degradation of mirror surfaces from exposure to the plasma. The typical loss in reflectivity for an experimental year is \(<15\%\). In addition, on the many experimental campaign time scale, neutron browning reduces the transmission of the fiber optics.

The spatial calibration is performed in four steps. First, metal targets are placed at fiducial locations that correspond with the nominal locations of the various neutral beams viewed by CER view chords. Second, the view chords that view each neutral beam are backlit with small bulbs inside their spectrometers, making images of their fiber tips on the metal targets. Third, those images are
marked on the target. The last step is to measure the locations of those images on the target and convert the measurements to the machine coordinate system based on the fiducial locations used to place the targets. A coordinate measuring machine (CMM) has been used to verify the locations of the in-vessel fiducials. Neutral beam locations are determined by mounting a laser in place of the ion source and tracing its path through the vessel.

This spatial calibration can easily provide information about the major radius and height of the intersection of view chords and the neutral beam. The alignment of the view chord with the toroidal, vertical, and radial directions is important for determining the plasma rotation, allowing the view chord trajectory to be represented as a vector, $\mathbf{s}$. The view chord trajectory is defined by the location of the lens center that the view chord passes through and the position of the view chord on a target.

A important concern for target based calibrations is whether the true neutral beam position is the same as the position of the target. That is, there is uncertainty as to whether the neutral beams are in the nominal locations when their ion source is installed and they are injecting power. This question is addressed in Chapter 4. The results shown there demonstrate that sufficient accuracy is obtained for CER analysis with the target calibrations.

The intensity calibration is performed in much the same way as the spatial calibration, but instead of drawing a fiber image on a target, a fiber image is placed on the emitting portion of a calibrated integrating sphere. The apparatus used is a Optronic Laboratories model OL 455-12-1 with an OL 450 controller. This is an 18 inch sphere with a 4 inch aperture that can easily accommodate multiple view chords. The radiance of the sphere (photons/s/m$^2$/sr/Å) is calibrated and is used to convert the measured signal rate (counts/s from the camera) to photons/s [as is needed in Eq. (3.5)]. This measurement is then used to determine density. By performing this calibration from inside the tokamak, all the apparatuses that are involved in the detection of light, from any in-vessel obstruction to the CCD electronics, are included.
3.1.3.2 Between discharge calibrations

Small variations in the CER optics throughout the day have prompted the implementation of two automatic calibrations that occur immediately before and after each plasma discharge. Before each discharge, a white light correction is used to determine the flat field response of each view chord. After each discharge, a wavelength calibration is performed with a calibration pen lamp. These measurements are made without perturbing any of the CER optics by using integrating spheres that are inserted between the fiber tips and their lenses at the tokamak ports with pneumatic actuators. The integrating spheres are custom designs made entirely from the Spectralon material available from the Labsphere Corporation.

The white light correction is provided by observing light from a quartz tungsten halogen bulb, whose intensity is essentially constant over the wavelength range observed (< 60 Å). This calibration corrects for differences in the intensity response across the detector region used for each view chord by determining correction values for each pixel that result in measuring a flat white light intensity spectrum. The intensity response can differ across the detector as a result of: wavelength dependence of CCD detected sensitivity, wavelength dependence of optics, vignetting of spectrometer and camera optics, differences in gain between multiple digitizers used to read data for a single view chord, sensitivity between detector pixels, and variations in sensitivity between pixels (e.g. dust on the detector). Typical corrections are < 5%, but there are cases where corrections near the edge of a view chord’s detector region are large due to significant vignetting inherent to the camera optics (see Section 3.2.3). Figure 3.2 shows two representative measurements of the white light spectrum. Applying the white light correction flattens these spectra.

White light spectra are essentially constant throughout an experimental campaign, but performing the calibration before each discharge automates the detection of the changes that do occur. For example, when the Pixelvision cameras change their gain setting (high or low gain are available, their difference in gain being a factor of five), the small misalignment between the two digitizers that read
data for a single view chord [shown in Fig. 3.2(a)] can change.

Figure 3.2: Normalized Measurements of white light spectra for two view chords. The white light spectra is needed to correct the difference in gain between the two halves of the detector, as seen in (a), and significant vignetting due to camera optics, as seen in (b).

The wavelength fiducial \( p_0 \) in Eq. (3.3) is measured by observing multiple Ne I emission lines from a capillary discharge lamp. The emission from these lamps is cold and designed for calibration purposes. The emission observed has wavelengths of 5274.0, 5280.1, 5298.2, and 5304.8 Å.\(^{53}\) These emission lines bracket the 5290.5 Å C VI emission and can be fit so that the pixel location of the C VI wavelength fiducial can be interpolated. The accuracy of these interpolations are on the order of 0.01 pixels, corresponding to a velocity on the order of 0.1 km/s. The wavelength fiducial is observed to change significantly over the course of a day, and this change is mainly correlated with changes in temperature within the CER laboratory. Changes between morning and evening on the order of a pixel have been observed, corresponding to a velocity of 10 km/s, illustrating the need for making wavelength fiducial measurements throughout the experimental day. The rate of change of fiducial between discharges is small enough that the \( \approx 90 \) s delay between the start of the discharge and the acquisition of wavecal fiducial data is insignificant.
3.1.3.3 Beam into gas calibrations

Beam into gas discharges are used to create instrumental responses, wavelength calibration offsets, beam path length normalization factors, and dispersions for every CER view chord. In addition, beam into gas discharges provide information of the geometry of the view chords and the neutral beams they view. This geometric information is used for the work in Chapter 4.

By injecting neutral beams into xenon gas (50 ms pulses per neutral beam), cold xenon atomic emission (no Doppler shift or Doppler broadening) is created. This emission is primarily useful because of the relatively bright Xe II 52922 Å emission\(^{54}\) that is very close to the typical C VI emission (5290.5 Å) observed by CER. This emission is used to construct instrumental responses by fitting multiple Gaussians to the resulting shape. The resulting instrumental response represents the line broadening inherent to the view chord’s optics, so it is convolved with a shifted, broadened Maxwellian to fit the CER data. Using other emission lines (potentially from other gasses) can be used to acquire the same information for CER analysis of other impurity emission at different wavelengths. Beam into helium shots are routine and enable good measurements of the He I 4713 Å emission which is very near to the He II (n= 4 \rightarrow 3) 4686 Å emission that can be observed with charge exchange. An example of the multi Gaussian fit to a Xe II beam into gas emission line is shown in Fig. 3.3

Beam into xenon shots are also used to determine any offset between the post-discharge wavelength calibration and the true wavelength when observed from tokamak illumination. This is because Xe II 52922 Å emission is bright and near to the 5290.5 Å C VI emission (≈ 8 pixels for typical dispersions). By comparing the measured wavelength of the Xe II emission from the post shot wavelength fiducial with the true wavelength, an offset in the wavelength fiducial is obtained. This offset is interpreted as being the result of the difference in fiber illumination when using the post-shot calibration integration spheres\(^{52}\) and the neutral beams themselves. Typical differences are less than 0.1 pixels, corresponding to a velocity of about 1 km/s.

Beam into gas discharges also provide a method for measuring relative
Figure 3.3: Xe II 52922 Å emission seen during a beam into gas shot that has been fit with multiple Gaussians in order to obtain an analytic model for the instrumental response for this view chord.

differences in intensity between view chords that result from the shape of the neutral beam and the different angles each view chord makes with the neutral beams. By using the beam into gas as a uniform light source, measurements of the intensity of neutral beam stimulated emission for every view chord can be normalized to one another. These relative factors are important because measurements of density need to know the observed neutral beam density that is dependent on the view chord’s path through the finite extent of the neutral beam. That is, to measure density, the value of \( \int n_b(l)dl \), where the integral is done over the view chord’s path, is needed. The effective path length, \( \Delta L \) can be defined as \( \hat{n}_b\Delta L \equiv \int n_b(l)dl \), where \( \hat{n}_b \) is the beam density on the beam centerline. The value of the \( \Delta L \) is found for a single view chord based on assuming the beam density to be Gaussian in the vertical and horizontal directions perpendicular to its propagation direction and using nominal values for the 1/e widths. With this calculation, the intensity of beam into gas measurements can be used to create a self-consistent set of effective path lengths for every view chord. Both He I 4713 Å
and Xe II 5292 Å emission have been used for this purpose.

The uniformity of the beam into gas light source is a result of the uniform distribution of the neutral gas in the vessel and negligible attenuation of the neutral beam as it passes through the gas. Uniform distribution of the neutral gas is accomplished by minimizing sources of pumping in the vacuum vessel (turbo pump valves closed, cryogenic pumps warmed), and injecting the gas long enough before the neutral beams turn on that its pressure stabilizes. Since the neutral beam pulses are short compared to the remaining rate that the gas is pumped out (from the neutral beam’s turbo pumps and cryogenic panels) this stabilization is relatively easy to achieve. The attenuation of the neutral beam is negligible because the gas is cold and its density is low. Furthermore, even when a neutral beam particle ionizes due to a collision, it does not leave the stream of the neutral beam because magnetic field coils are not energized during the discharge. The resulting uniformity of the neutral gas has been previous verified.\textsuperscript{51}

Beam into neon shots are used to determine view chord dispersions. The same neon emission used for the post-discharge wavelength calibration is used to calculate the dispersion. The dispersion is determined from the slope of a fit to the fit emission locations versus their true wavelengths. The dispersion is assumed to be constant across the detector region used for each view chord, and there is no indication in the linear fits that a wavelength dependent representation should be used. Errors in the fits propagated through the linear fit suggest that the dispersion is determined with an error of at most 1\% with the large majority of dispersions determined to less than 0.5\%. These dispersions are used to determine the geometric factors in the grating equation:

\[ D = \frac{l_p d}{f} \left[ \sqrt{\cos^2 \phi - \frac{\lambda^2}{4d^2}} - \frac{\lambda}{2d} \tan \phi \right]. \]

(3.6)

where \( l_p \) is the pixel size in the wavelength dispersion direction, \( d \) is the grating groove spacing, \( f \) is the focal length of the focusing mirror, \( \phi \) is the half angle between the incident and diffracted rays from the grating (a design parameter of the spectrometer), and \( \lambda \) is the wavelength. Equation (3.6) is used to calculate the dispersion at other spectrometer wavelength settings.
Beam into deuterium shots are used to observe the shift between $D_\alpha$ from a stationary gas and the high speed neutral beam particles. The Doppler shift contains geometric information about the angle between the view chord and the neutral beam. In addition, if the same shot is done with a static toroidal field, the beam neutral $D_\alpha$ light will be Stark split. The Stark splitting contains information about the major radius of the point where view chords intersect the neutral beam because the major radius sets the magnitude of the $\mathbf{V} \times \mathbf{B}$ electric field felt by the beam neutrals. Measurements from these two types of beam into deuterium shots are used to form an in-situ spatial calibration for the CER system. This process is described in Chapter 4. Note that the beam into deuterium shot that measures the Stark splitting is the only beam into gas shot where magnetic field coils are energized.

### 3.1.4 Accounting for the cross section effect

As introduced in Section 1.5.2, rotation measurements made with CER are affected by the energy dependence of the charge exchange cross section. On DIII-D, the effect this has on carbon temperature measurements is small, but these effects must be accounted for in order to obtain accurate rotation measurements. Atomic physics calculations can be used to correct the measurements, but it is also possible to measure the effect directly. Measuring the cross section effect is much preferred because it removes uncertainty in the complicated atomic physics calculations that require knowledge of the charge exchange cross section, beam density fractions, and excited state fractions.

Measuring the cross section effect is made possible by the presence of the co- and counter-current neutral beams on DIII-D. For carbon, the cross section effect results in measured velocities that are more into the stream of the neutral beams than the true velocity. Therefore, measurements of toroidal rotation on the co-current neutral beam are more counter-current than the true velocity, and vice versa. By combining the two measurements, it is possible to measure the cross section effect and the true toroidal rotation, i.e. the toroidal rotation of the $C^{+6}$ ions as opposed to the biased population of $C^{+5}$ ions that come from charge
exchange events.

Let each CER view chord be represented as a normalized vector
\[ \hat{s} = s_\phi \hat{\phi} + s_R \hat{R} + s_z \hat{z}, \]
the direction of injection for a neutral beam be represented by
\[ \hat{V}_b = \cos \gamma \hat{\phi} - \sin \gamma \hat{R}, \]
and the magnitude of the cross section effect on velocity to be \( \alpha \). It is assumed that the cross section effect is aligned with the neutral beam because neutral beam velocities are large (2800 km/s for typical full energy beam components) compared to plasma rotations (between 0 and 500 km/s), and their velocity therefore dominates the energy dependence of charge exchange events. In this case, the velocity of the \( C^{+5} \) population that has undergone charge exchange can be represented as
\[ V' = V + \alpha \hat{V}_b. \]
Projecting this onto the measurement view for a tangential view chord (\( s_z \approx 0 \)) gives the apparent velocity,
\[ V_{\text{app}} = \hat{s} \cdot V' = s_\phi (V_\phi + \alpha \cos \gamma) - s_R \alpha \sin \gamma. \tag{3.7} \]

If two measurements of the apparent velocity, \( V_1 \) and \( V_2 \), are made at the same position in the plasma (same \( R \) and \( z \)) with two different viewing geometries (\( \hat{s}_1 \) and \( \hat{s}_2 \)), it is possible to solve for the two unknowns in Eq. (3.7), \( \alpha \) and \( V_\phi \):
\[ \alpha = \frac{s_2 \varphi V_1 - s_1 \varphi V_2}{s_2 R s_1 \varphi \sin \gamma_2 - s_1 R s_2 \varphi \sin \gamma_1 + s_1 \varphi s_2 \varphi (\cos \gamma_1 - \cos \gamma_2)}, \tag{3.8} \]
\[ V_\varphi = \frac{V_1 + \alpha s_1 R \sin \gamma_1}{s_1 \varphi} - \alpha \cos \gamma_1 = \frac{V_2 + \alpha s_2 R \sin \gamma_2}{s_2 \varphi} - \alpha \cos \gamma_2. \tag{3.9} \]

When the two view geometries are of neutral beams with opposite signs of \( \cos \gamma \), experimental uncertainties in the result are significantly reduced. This is the case for the co- and counter-current neutral beams, so the CER view chords of these beams can be used to determine \( V_\varphi \) and \( \alpha \).

As an illustration of this measurement, consider Fig. 3.4. In Fig. 3.4(a), measurements of apparent velocity from two tangential CER views are shown. Both views see their respective neutral beam at \( R = 1.898 \) m, and their \( z \)-positions are separated by only 3 mm (an insignificant separation on the plasma midplane). The beam sources are run with the nearly the same accelerator voltage. Taking the magnitude of the atomic physics effects to be the same then enables the measurement of \( V_\varphi \) and \( \alpha \) from Eqs. (3.8) and (3.9). Those results are shown in Fig. 3.4(b). Note that, because of the dependence of the cross section effect
on beam properties that are not flux functions (e.g. beam component density fractions), the cross section effect is not a flux function.

![Graphs showing velocity measurements](image)

**Figure 3.4:** (a) Measurements of apparent velocity from two view chords with equal major radius positions on co-current and counter-current neutral beams and the true toroidal rotation that is determined from these measurements. (b) the measured value of the velocity induced along the neutral beam injection direction (the cross section effect).

Note that the cross section effect on the direct measurement of poloidal rotation cannot be measured with the method outlined here because of the need to know the excited state lifetime. Though it is possible to design CER view chords that could determine the gyro-orbit cross section effect, these are not available on DIII-D. In Section 3.3, a method for measuring poloidal rotation without atomic physics calculations via a combination of multiple co- and counter-current toroidal rotation view chords is described.

### 3.2 Upgrades to DIII-D CER

New measurements present in this dissertation were made possible by several upgrades to the CER system that were made before the 2011 DIII-D experimental campaign. The basic goal of the upgrades was to enable CER measurements to be made on the high-field side of the magnetic axis. This required
various changes to be made to the CER system. The important elements of these efforts are discussed in this section.

### 3.2.1 Addition of low major radius views

The addition of low major radius tangential CER views was done so that a more complete measurement of toroidal rotation could be made. It was important that these views be able to measure toroidal rotation on the high-field side of the magnetic axis \( R_{\text{axis}} \approx 1.75 \text{ m} \) because these measurements can be combined with low-field side measurements to determine the poloidal rotation (see Section 3.3). In order to exploit the ability to measure the cross section effect and determine the toroidal rotation without any atomic physics calculations (see Section 3.1.4), these high-field side tangential views were added to both the co-current and counter-current neutral beams. The added view chords are shown in green in Figs. 1.2 and 1.3. Before the 2011 experimental campaign, four view chords of the co- and counter-current neutral beams (eight total) were added, and an additional view chord for each set of beams was added before the 2014 experimental campaign (ten total). In Fig. 1.2, the two green view chords on the co- and counter-current neutral beams with the highest major radii were the most recently added.

Ideally, as much coverage of the low-field side of the plasma is desirable, but it was not possible to obtain views with \( R < 1.3 \text{ m} \) without extensive modification of existing port designs. These modifications required more resources than were available when these upgrades to the CER system were being done. The two CER ports where the new high-field views were added are designated 165R0 and 315T0 (the lower-right and upper-left ports in Fig. 1.2, respectively). The lowest major radius of the high-field side view chords that were added is 1.28 m.

Tangential views of co-current neutral beams on the high-field side were already available at 315T0, but not coupled to detectors. Views of the counter-current neutral beams were chosen so that their radii on the more perpendicular beam would match the views of the more perpendicular co-current neutral beam. This choice aids the measurement of the cross section effect and was also made with the expectation that signal levels will be better on the perpendicular sources.
because their beam attenuation is less than the tangential sources (which must travel through more plasma to reach the same major radius position). The original views of the counter-current neutral beams were designed to match views of the co-current neutral beams on the more tangential beam.

As a result of using existing optics and fiber mounting at the 315T0 port, only minor modifications were needed there. The most significant modification was an increase in the size of the integrating sphere used for between discharge calibrations (see Section 3.1.3). More extensive modifications of the 165R0 were required. The major difficulty arose from the fact that the large angular spread of CER views needed (to retain previous coverage of the counter-current beam while adding low major radius views) was too large for the re-entrant tube used on the port. To solve this problem, a shorter focal length lens was used so that the fiber optics could be fit within the re-entrant tube. In addition, the lens was tilted so that the coma resulting from off axis lens illumination could be more evenly distributed between the different view chords. After the lens tilt, the coma is a relatively minor aberration because the spectrometer optics ($\approx f/4$) are much slower than the lens ($\approx f/2$).

Accommodating these new high-field side views at 315T0 and 165R0 on existing spectrometers and cameras was desired so that hardware cost could be minimized. This was done by increasing the view chord density on existing detectors. Two relatively inexpensive hardware modifications were developed for this purpose and are described in Sections 3.2.2 and 3.2.3.

An unfortunate consequence of the modifications to the 165R0 port that were necessary to obtain the desired views was the lack of room for an integrating sphere that could be used for the between discharge calibrations (see Section 3.1.3.2). As a result, a different method for obtaining this calibration data was developed. This method uses the multiple spectrometer entrance slits (present on all the spectrometers that image view chords from 165R0) to obtain the white light and wavelength calibration data between each discharge. For each view chord, a calibration fiber optic is mounted on an adjacent slit and coupled to a standard integrating sphere inside the CER laboratory. This allows calibration light to
be observed by 165R0 view chords with a wavelength shift due to the horizontal separation between the slits.

The wavelength shift has no effect on the white light calibration, but different calibration lamp emission must be used to create the wavelength fiducial. The 5222.4, 5234.0, and 5400.6 Å Ne I emission lines\textsuperscript{53} are used for these adjacent slit wavelength calibrations. By installing the same adjacent slit calibration capability for view chords that are not at the 165R0 port, it was possible to compare the wavelength calibrations between the adjacent slit and on-vessel integration sphere measurement methods. In Fig. 3.5, the two measurements of the wavelength fiducial throughout an experimental week are shown. It can be seen that both methods give comparable results.

\begin{figure}
\centering
\includegraphics[width=\textwidth]{figure35.png}
\caption{Changes in measurements of wavelength fiducial for a view chord as determined with calibration light from the on-vessel integrating sphere and an adjacent slit fiber coupled to a CER laboratory integrating sphere for a typical experimental week. The small differences are appropriate given the uncertainties in determining the wavelength fiducial.}
\end{figure}
3.2.2 Novel Czerny-Turner Astigmatism Correction

Basic Czerny-Turner spectrometers\textsuperscript{47} suffer from astigmatism. This is because light from the object to be analyzed is incident on the first and second spherical mirror from off their optical axes. As a result, the focal distance for the sagittal image is different from the focal distance for the tangential image. The object and the optical axis of the first spherical mirror define the tangential plane, and the plane perpendicular to this plane and containing the principal ray from the object define the sagittal plane.\textsuperscript{56} The usual Czerny-Turner spectrometer is oriented on a bench such that the sagittal plane is vertical and the tangential plane is horizontal. This terminology is more convenient and will be used for the remainder of this section. A schematic showing the astigmatism of a Czerny-Turner spectrometer is shown in Fig. 3.6.

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{fig3_6.png}
\caption{A schematic of the (a) horizontal (tangential) and (b) vertical (sagital) planes of a Czerny-Turner spectrometer, showing that the position where the light rays which diverge horizontally (red) are re-focused is different from the position where the light rays which diverge vertically (green) are re-focused.}
\end{figure}

This blurring of the focal image is detrimental to performing spectroscopy. The detector cannot be placed at the position of best vertical focus because the object’s wavelength is dispersed in the horizontal direction, and wavelength resolution is fundamental to CER analysis. However, when the detector is placed at the position of best horizontal focus, the image height is larger than necessary, forcing more detector area to be used to analyze the light. The ability to use
detector area more efficiently reduces the need for expensive spectrometers and detectors. Example images from the two focus locations are shown in Fig. 3.7

![Images](image.png)

**Figure 3.7:** Images of 5401 Å Ne I emission from a calibration lamp created by an astigmatic Czerny-Turner spectrometer at positions of (a) best horizontal focus and (b) best vertical focus. Red indicates higher intensity.

In order to correct this astigmatism, a technique was developed for bringing these two focus locations together. The first step to doing this is determining the initial separation between the focal planes. The horizontal and vertical focal distances can be calculated as:

\[
Q_h = \frac{f \cos(\vartheta)(1 + X/p) - X}{2 - f \cos(\vartheta)/p - X/[f \cos(\vartheta)] + X/p},
\]

\[
Q_v = \frac{f - X \cos(\vartheta) + X f/p}{2 \cos(\vartheta) - X \cos^2(\vartheta)/f - f/p + X \cos(\vartheta)/p}.
\]

Here, \(f\) is the focal length of the spectrometer mirrors, \(p\) is the object distance from the collimating mirror (typically the same as the focal length), \(X\) is the length of the path between the two mirrors, and \(\vartheta\) is the angle the incident light makes with the optical axis of a mirror. The value of \(Q_v - Q_h\) represents how much farther the vertical focal distance is than the horizontal focal distance. When \(Q_v = Q_h\), there is no astigmatism.
Many Czerny-Turner spectrometers are manufactured with non-spherical mirrors that can correct for the astigmatism. For instance, a toroidal collimating mirror can be designed so that its horizontal and vertical focal lengths are different from each other in a way that makes \( Q_v - Q_h = 0 \). Ten of the twelve Czerny-Turner spectrometers used by the CER system have toroidal collimating mirrors. However, two Czerny-Turner spectrometers were not manufactured with a solution for their astigmatism. As part of the 2011 upgrade to the CER system, it was necessary to correct their astigmatism. These spectrometers were 1/2 m Acton Research Corporation \( f/4.1 \) spectrometers. Using toroidal mirrors would have required expensive optics, precise design, and presented significant installation challenges. Instead an inexpensive, easy to design, easy to install solution was developed.

This solution is based on using the entrance slit of the spectrometer to shift the position where light from the object diverges vertically from the position where it diverges horizontally. Since the slit is narrow when compared to the fiber optic that brings light to the spectrometer (\( \approx 0.2 \) mm versus a fiber optic diameter of 1.5 mm), it defines the horizontal object to be imaged, even when the fiber is pulled back from the slit. Conversely, when the fiber is pulled back from the slit, the light continues to diverge when traveling between the fiber and the slit. If the fiber is pulled back the correct distance, the horizontal and vertical images can be made to focus on top of each other, correcting the astigmatism.

The distance the fiber should be pulled back to correct the astigmatism can be determined by using Eqs. (3.10) and (3.11) and allowing the vertical and horizontal object distances to be different. Using \( p_v \) as the vertical object distance and \( p_h \) as the horizontal object distance, \( p_v \) can be found from \( Q_v(p_v) - Q_h(p_h) = 0 \) after taking \( p_h \) to be the distance between the collimating mirror and the slit. For typical parameters, though, it can be verified that pulling the fiber back a distance equation \( Q_v(p_v) - Q_h(p_h) \) gives the desired result to good approximation. When the fiber is pulled back this amount, though, light will diverge in all directions and only a small amount will make it through the slit. The solution developed as part of this dissertation places a thin glass plate between the fiber and the slit so that the light is confined by total internal reflection in a one-dimensional waveguide.
but allowed to diverge vertically before it is incident on the slit. This method is advantageous not necessarily for its extreme accuracy, but instead for its cost. The glass piece required to implement this solution can be acquired for less than $100, and only requires a simple mounting system.

A schematic showing the fiber and glass piece mounting to the spectrometer is shown in Fig. 3.8. In this figure, angles and distances are labeled that will be used to explain the method for determining the appropriate glass length. As when the fiber was simply pulled back from the slit, the goal is to make the glass plate length so that $Q_v(p_h + x) - Q_h(p_h) = 0$. This is because light from the fiber, after refracting from the end of the glass plate appears to come from a position that is a distance $x$ behind the slit.

![Diagram of the fiber optic, glass plate, and spectrometer optics coupling. The most off-axis light-ray accepted by the spectrometer is shown. The $f$-number determines the angle $\alpha'$ which can be used to determine the optimal glass length for correcting Czerny-Turner astigmatism.](image)

**Figure 3.8:** Schematic of the fiber optic, glass plate, and spectrometer optics coupling. The most off-axis light-ray accepted by the spectrometer is shown. The $f$-number determines the angle $\alpha'$ which can be used to determine the optimal glass length for correcting Czerny-Turner astigmatism.

Consider light existing the glass plate at angle $\alpha$, with respect to the normal. The largest this angle can be is determined by the spectrometer $f$-number, $N$, i.e.

$$\tan \alpha' = \frac{1}{2N}. \quad (3.12)$$

Using Snell’s Law, the distance behind the slit that these light rays appear to
diverge from is
\[ x = \frac{h}{\tan \alpha'} = \frac{L \tan \alpha}{\tan \alpha'}. \tag{3.13} \]
Taking \( x = Q_v - Q_h \) to correct the astigmatism, the desired glass length is
\[ L = (Q_v - Q_h) \frac{\sqrt{n^2 - \sin^2 \alpha'}}{\cos \alpha'}. \tag{3.14} \]

Using Eq. (3.14) to design a glass plate, the astigmatism of a Czerny-Turner spectrometer can be corrected. An example of the difference between the best possible image obtainable before and after the correction is shown in Fig. 3.9. Here the dramatic improvement in image height can be seen, effectively doubling the data density on a given detector.

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{spectroscopic_image.png}
\caption{Best spectroscopic image of the 5401 Å Ne I emission from a Czerny-Turner spectrometer before and after astigmatism correction with a thin glass plate. Image height has been approximately halved. Red indicates higher intensity.}
\end{figure}

Note that, in Eq. (3.14), no approximation was made for the smallness of the angles, so this analysis holds for fast spectrometers. Also note that for typical parameters, the dependence of \( L \) on \( \alpha' \) is weak. For this reason, sufficient accuracy is obtained by using the last light ray to perform the length calculation even though the light entering the spectrometer is distributed between \( \alpha' = 0 \) and...
the last accepted ray. Also note that the length depends on the index of refraction, so there can be chromatic blurring of the correction due to changes in the glass piece’s refractive index.

The issue of chromatic blurring was investigated during initial tests of the glass plate design on a McPherson Corporation Model 207 2/3 m f/4.7 spectrometer. A custom B270 glass plate with a length determined by Eq. (3.14) and polished faces for the interface between the fiber and the plate and the plate and the slit (see Fig. 3.8) was used for these tests. It is relatively easy to acquire glass whose length precision is ±0.25 mm, and this limit was kept in mind when determining the accuracy needed for the glass plate length. The width of the glass plate should be significantly wider than the slit (to ease alignment), but not any wider than the fiber diameter (which would decrease light throughput).

![Figure 3.10](image)

**Figure 3.10:** Horizontal and vertical image widths as a function of camera position for a glass plate corrected Czerny-Turner spectrometer at (a) 5400 (b) 3600 (c) 6700, and (d) 8990 Å. Locating the camera at 22.25 mm results in good horizontal and vertical focusing for all wavelengths.

The chromatic blurring was investigated by measuring the horizontal and vertical focus at wavelengths other than the design wavelength (5401 Å). To
determine the focus, measurements were made of the width of the image when summed vertically (giving a measure of the horizontal width) and summed horizontally (giving a measure of the vertical width). Focus is optimized when the width is minimized. The width metric used was taken to be the full width about the image center (determined by an intensity weighted average of the image position) within which 90% of the intensity resided. This metric was chosen because it tends to minimize the effect of coma, which can result in low intensity, large width tails forming on the side of images. The result of these scans at 5400, 3600, 6700, and 8990 Å are shown in Fig. 3.10. Here it can be seen that the chromatic blurring is minimal, and designing a solution for one wavelength yields good correction for all wavelengths in the visible spectrum.

It has been experimentally determined that use of the glass plates reduces the total light throughput by $\approx 10\%$. This is an acceptable tradeoff for the increase in image density. Finally, it is important to point out that, aside from correcting the astigmatism, there is no major change in the spectrometer optics. This can be seen in Fig. 3.11 where the vertical and horizontal instrumental response are shown for the 5401 Å calibration lamp Ne I emission. The instrumental response is compared to another form of correction available for the McPherson Corporation Model 207 spectrometers, a cylindrical mirror that can be placed in front of the entrance port. Note that this cylindrical mirror adds a vertical magnification to the spectrometer, which is why the total heights in Fig 3.11(b) are different. Aside from some slight rounding of the image borders near the edges, the glass plate instrumental responses are the same as the instrumental response that is created with a cylindrical mirror correction.

3.2.3 Masking charge coupled device cameras

Densely packing spectroscopic data from multiple CER view chords onto a single charge coupled device (CCD) detector is necessary for any large CER diagnostic because it dramatically reduces the hardware costs. Though a CCD camera chip itself may be small (see Table 3.1), the camera form factors are much larger, typically a few inches to a side. As a result, even when using complex optics
Figure 3.11: Horizontal (a) and vertical (b) instrumental responses for a Czerny-Turner spectrometer with astigmatism corrected by either a glass plate or a cylindrical mirror. The two corrections yield very similar results.

At the spectrometer exit to separate the different spectra, it is difficult to use more than two cameras to analyze data from a single Czerny-Turner spectrometer. It is therefore desirable to develop methods for imaging multiple view chords on a single camera. Any such multi-view design must operate without cross talk between different view chords.

Taking the direction that wavelengths are dispersed in a Czerny-Turner spectrometer to be the horizontal direction, the prevention of horizontal cross talk necessitates various design considerations. Cross talk in the vertical direction is generally associated with cross talk in the time domain. This is because CCD cameras are oriented to read out in the vertical direction so that pixels at the same wavelength can be summed before being digitized, thereby increasing their signal to noise ratio.

If two fibers containing light from two view chords are placed horizontally adjacent to one another at the entrance of the spectrometer, good horizontal focusing ensures that their images of the same wavelength will be well separated. However, light of a different wavelength than the spectrometer setting will appear to either side of the fiber images and can result in cross talk between the two view chords. If the same two fibers are placed on top of one another at the entrance of the spectrometer, good vertical focusing ensures that any vertical cross talk will
be minimal. However, depending on the readout speed, it is possible for these view chords to cross talk with one another. This is especially true when a CCD camera is running at its fastest speed and it is possible that the entirety of the chip is not read out at once. In this case, data from a view chord is only shifted toward the read out register and then remains in that position, being exposed to light from a different view chord, until the next frame is acquired.

These two considerations motivated the design of a masking system for Sarnoff Corporation CAM1M100 CCD cameras. This design enables eight view chords to be analyzed on a single detector without any shifted wavelength or time domain cross talk, even when the camera is acquiring exposures as fast as possible and only reading a quarter of its full frame at a time. A schematic showing this process is shown in Fig. 3.12. This mask setup was originally developed on a McPherson Corporation Model 207 2/3 m f/4.7 spectrometer, but ultimately implemented on two Acton Research Corporation 1/2 m f/4.1 and two Acton Research Corporation 2/3 m f/4.7 spectrometers, all with Sarnoff Corporation CAM1M100 CCD cameras. The increased density of view chords on detectors that the implementation of these masks enabled was needed to increase the number of CER view chords without the need for significant new detector equipment and the associated cost.

### 3.3 New poloidal rotation measurement technique

The development of a new poloidal rotation diagnostic is motivated by the desire for high quality measurements that can be used to rigorously test poloidal rotation theories. As will be described in this section, the new diagnostic that has been developed does not require atomic physics calculations to account for the cross section effect. This feature is the main advantage that this new diagnostic has when compared to direct measurements of poloidal rotation on DIII-D.
Figure 3.12: A schematic of the CCD mask system developed to prevent cross talk between multiple view chords on a single spectrometer. Light from view chord X is imaged in region A1 and light from view chord Y is imaged in region B2. When the first frame is acquired, the mask prevents wavelength shifted light from X from landing on B1, where the next frame for Y is currently residing. When the second frame is acquired, the mask prevents wavelength shifted light from Y from landing on A2, where X’s first frame data is residing until the next frame read out.

3.3.1 Principles of new measurement technique

The new diagnostic is based on the functional form of the lowest order flow equations in a tokamak. For each plasma species, the simplest representation of the toroidal and poloidal flow is

\[ V_\psi = k(\psi)B_\psi + R\omega(\psi), \]  
\[ V_\theta = k(\psi)B_\theta, \]

where \( R \) is the major radius, \( k \) and \( \omega \) are functions of the poloidal flux function \( \psi \), and \( \theta \) is the poloidal coordinate [also see Eq. (2.11)]. This diagnostic uses Eq. (3.15) and measurements of toroidal rotation at two positions on a flux surface to determine \( k \) and \( \omega \). The value of \( k \) is then used to determine the poloidal rotation with Eq. (3.16). Since this method is based on measuring the poloidal asymmetry of the angular rotation, it is referred to as PAAR. The new view chords that enable PAAR analysis are shown in green in Figs. 1.2 and 1.3.
In the formulation of Eqs. (3.15) and (3.16), toroidal rotation is taken to be small compared to the ion thermal speed ($|V_\phi|/\sqrt{2T/m} \ll 1$) and, hence, density is a flux function. This assumption is not valid for the carbon impurity in many DIII-D plasmas. All PAAR results use a formulation that does not make this assumption, which is explained in Section 3.3.3. However, the basic mechanics of the diagnostic are unchanged from those described up to this point. For clarity, the full description is not presented until the preliminary description of the diagnostic is complete.

Since the new view chords observe the co- and counter-current neutral beams, measurements of toroidal rotation can be made without atomic physics calculations on both the high- and the low-field side of the tokamaks midplane. As a result, the PAAR method is able to determine the poloidal rotation without atomic physics calculations. Another advantage of the PAAR method is the fact that the poloidal rotation is present in the toroidal rotation with a multiplicative factor of $B_\phi/B_\theta$ (4 to 10 is typically in DIII-D), making the measurement of this typically small value easier and more certain. Poloidal rotation can typically be determined with the PAAR method with 1-σ errors of $\approx 1$ km/s. An additional advantage of the PAAR method is its high spatial resolution in the core of the plasma. This is a result of the fact that the vertical dimension of the DIII-D neutral beams is larger than the horizontal dimension by a factor of $\approx 3$. Finite spatial averaging inherent to the vertical CER measurements is exacerbated in the core where flux surface heights are reduced. Even though the vertical CER view chords are nearly tangential to the flux surfaces at the midplane, charge exchange signal will be acquired from a significant range of different flux surfaces. This effect is less important when plasma gradients in the core are small, but can be debilitating when considering very steep core gradients, e.g. as in plasmas with an internal transport barrier (ITB). This is because direct measurement of poloidal rotation needs a good measurement of the vertical velocity near the axis in order to determine the excited state lifetime.\(^{27}\)

There are some difficulties associated with the PAAR method. Most notably, an accurate magnetic equilibrium reconstruction is needed to map the
toroidal rotation measurements on the high-field side to the low-field side so that
the poloidal asymmetry can be measured. It is fortunate that the accuracy of
an equilibrium reconstruction can be easily corroborated by measurements of ion
temperature symmetry. The ion temperatures measured by the same view chords
that measure the toroidal rotation used in the PAAR method should form a single
profile when plotted as a function of a magnetic coordinate such as \( \rho \). This is
because, in the core of DIII-D, rapid transport parallel to the magnetic field ensures
that ion temperature is constant on flux surfaces.

As an example of this phenomenon, consider the ion temperature profiles
shown in Fig. 3.13(a). The good alignment of the high- and low-field side
temperature measurements corroborate the accuracy of the magnetic equilibrium
reconstruction for these low-confinement mode (L-mode) and high-confinement
mode (H-mode) plasmas. In Fig. 3.13(b), the toroidal angular rotation, \( \frac{V_\phi}{R} \),
plotted with the same magnetic equilibrium reconstructions is nearly constant for
the L-mode plasma and clearly varying around the flux surface for the H-mode
plasma. According to Eq. (3.15) then, it is expected that poloidal rotation is near
zero for the L-mode plasma and significant for the H-mode plasma.

A limitation of the PAAR method, as it is implemented on DIII-D, is
the limited radial range. Neutral beam attenuation is large enough that charge
exchange signal for the high-field side view chords can become weak when density is
moderate or large for DIII-D, \( n_e > 5 \times 10^{19} \text{m}^{-3} \). Signal levels and current diagnostic
port optics determine the lowest major radius where CER measurements can be
made. For typical plasma shapes and magnetic axis locations, poloidal rotation
can be determined with the PAAR method inside \( r/a \approx 0.6 \).

Independent of the first development of the PAAR diagnostic on DIII-D,\(^{59}\)
the same method for measuring poloidal rotation was also developed on TCV.\(^{14}\)
In addition, a similar diagnostic has been developed on the TJ-II stellarator.\(^{60}\)
Figure 3.13: (a) Ion temperature profiles for a L-mode (144948) and H-mode (149468) plasma. Good overlap of these CER measurements on the high- and low-field side of the magnetic axis ($R_o$) indicates that the magnetic equilibrium reconstruction is accurate. In (b), the corresponding profiles of toroidal angular rotation show that it is not necessarily a flux function.

3.3.2 Determining poloidal rotation from fits to toroidal rotation

In order to fit the raw measurements of tangential line of sight velocity to the functional form of Eq. (3.15), it is necessary to account for the cross section effect. Taking $\hat{s} = s_\varphi \hat{\varphi} + s_R \hat{R} + s_z \hat{z}$ to be the normalized vector representing a CER view chord’s sightline, $\beta$ to be the angle between the view chord’s sightline and the neutral beam injection direction, and $\alpha$ to be the magnitude of the cross section effect in the direction of neutral beam injection (the same as in Section 3.1.4), Eq. (3.15) can be rewritten to describe the line of sight velocity as

$$V_{\text{LOS}} = (kB_\varphi + R\omega)s_\varphi + \alpha \cos\beta.$$ \hfill (3.17)

Equation (3.17) can be used to determine the value of $k$ from the measurements of tangential line of sight velocity from co- and counter-current neutral beam views on the high- and low-field sides of the plasma. A fitting algorithm optimizes cubic spline representation of $k(\rho)$, $\omega(\rho)$, and $\alpha(R)$ in order
to generate the best fit to the $V_{\text{LOS}}$ measurements. The $B_\phi$ factor is determined from the magnetic equilibrium reconstruction, and the $R$, $s_\phi$, and $\cos \beta$ factors are determined from the spatial calibration (see Section 3.1.3.1 and Chapter 4). The spline fit for $\alpha$ is the simplest to generate because it can be constrained by direct measurements of the cross section effect. The resulting spline fit for $k$ is then used to calculate the poloidal rotation with Eq. (3.16). Note that the poloidal rotation is not a flux function. By convention, PAAR poloidal rotation is determined on the outboard midplane of the tokamak vessel (that is $z = 0$, not the height of the plasma magnetic axis, though the two are usually similar). At this location, the positive poloidal direction is downward, regardless of the plasma current direction, so that the $(r, \theta, \varphi)$ coordinate system is right handed.

![Figure 3.14: Poloidal rotation from matched discharges with oppositely directed toroidal fields with the positive poloidal direction being (a) downwards on the outboard midplane and (b) the ion-diamagnetic drift direction.](image)

A basic check that can be performed with any poloidal rotation diagnostic is the analysis of similar plasmas that differ only by the sign of their toroidal field. In such plasmas, all drift directions change signs, and it is expected that the poloidal rotation will change direction as a result (this effect can also be seen formally in neoclassical predictions of poloidal rotation\textsuperscript{31}). This phenomena was observed with the PAAR method by observing poloidal rotation for two matched
Figure 3.15: Poloidal rotation determined with the PAAR method for a (a) L-mode and (b) H-mode plasma. The poloidal rotation is small in (a) and significant in (b), as was expected from the profiles of $V_r/R$ shown in Fig. 3.13.

L-mode discharges that have oppositely direction poloidal fields ($|B_\varphi| = 2.1 \, T$, $I_p = 1.2 \, \text{MA}$, $n_e = 2 \times 10^{19} \, \text{m}^{-3}$, $P_{\text{NBI}} = 2 \, \text{MW}$, $\tau_{\text{NBI}} = 1.8 \, \text{Nm}$). In Fig. 3.14(a), poloidal rotation plotted for these two discharges against a physical coordinate system is shown, and in Fig. 3.14(b), it can be seen that the poloidal rotation measurements match when plotted against a magnetic coordinate system. The agreement between these measurements in Fig. 3.14(b) corroborates the accuracy of the PAAR diagnostic.

Error bars shown for PAAR results are generated with Monte Carlo analysis that takes into account uncertainty in the rotation measurements and the position (in major radius) of those measurements. The results of this analysis are dominated by the error in the rotation measurements. Typically, the error in any single CER measurement, calculated with photon statistics, is lower than the observed random variation in the data. For this reason, measurement errors are determined by taking the standard deviation of all the measurements within a time window where the plasma conditions are static.

Poloidal rotation results from PAAR analysis for the plasmas in Fig. 3.13 are shown in Fig. 3.15. As expected, the poloidal rotation for the L-mode plasma is
small, while the poloidal rotation for the H-mode plasma is significant. Note that the measurement errors are quite low for these measurements, allowing poloidal rotation to be accurately determined within 1 km/s.

In order to better illustrate the method for measuring poloidal rotation by fitting toroidal rotation data on the high- and low-field side of the tokamak midplane, the fit that generates the result shown in Fig. 3.15(b) is shown in Fig. 3.16. In Fig. 3.16(a), the fit to the toroidal rotation measurements on the co- and counter-current neutral beams is shown to be good. The splines for $\omega$ and $k$ that yield this fit are shown in Fig. 3.16(b), and the value of $k$ is used to determine the poloidal rotation with the value of $B_\theta$ from the magnetic equilibrium reconstruction. Note that, due to the cross section effect, and the presence of measurements on the co- and counter-current neutral beams, $V_\varphi$, appears doubly valued. The basic mechanics of this fit can be understood with Eq. (3.17), but, as explained in Section 3.3.3, the full equation used to create the fit is Eq. (3.33).

**Figure 3.16:** (a) A fit to the measurements of $V_\varphi$ for high- and low-field side view chords on co- and counter-current neutral beams, and (b) the spline representations of $\omega$ and $k$ that produce this fit. The poloidal rotation that results is shown in Fig. 3.15(b).

Generating the high quality equilibrium reconstruction with EFIT needed to determine $B_\varphi$, $B_\theta$, and the mapping of real space onto a magnetic coordinate
system can be time consuming. Ultimately, the PAAR method is based on measuring asymmetries within flux surfaces, so the generation of those flux surfaces is of paramount importance. A basic magnetic reconstruction can be created with only information from magnetic coils external to the plasma, but these reconstructions lack accuracy in the core where the results are most important for PAAR analysis. By adding measurements of the magnetic field line pitch from the motional Stark effect (MSE) diagnostic, an improved magnetic equilibrium reconstruction can be obtained. In addition, constraints to the total pressure profile (including fast-ion pressure calculated by NUBEAM) can be added to EFIT, and a “kinetic reconstruction” can be generated that is usually the most accurate.

Kinetic magnetic equilibrium reconstructions are typically needed for PAAR analysis. It has been found that, with sufficient skill in the operation of EFIT, accurate magnetic equilibrium reconstructions that yield single ion temperature profiles can be readily obtained. As an example, three temperature profiles with good overlap between the high- and low-field CER measurements are shown in Fig. 3.17. The plasmas shown in this figure include L-mode (144569), H-mode (149661), and quiescent high-confinement mode (QH-mode, 149094).

Confidence in the spatial calibration and consistent success in achieving high- and low-field side temperature overlap has prompted an upgrade to EFIT that allows high- and low-field side temperature measurements to be added as a constraint to the reconstruction. Since the lowest major radius of MSE measurements made on DIII-D is 1.5 m, the high-field side ion temperature measurements provide EFIT with new information about the internal structure of the flux-surfaces. The procedure for adding this constraint to EFIT is to specify the major radii of the high-field side view chords and the major radii where the temperature should be the same on the low-field side. This is determined with a natural spline fit to the low-field side measurements, which have higher radial resolution than the high-field side measurements. To demonstrate the efficacy of this constraint, a magnetic equilibrium reconstruction that uses only magnetic probe and MSE data is compared to a reconstruction that also has ion temperature
Figure 3.17: Ion temperature profiles with high- \((R < R_o)\) and low-field side \((R > R_o)\) measurements, demonstrating that good temperature overlap can be achieved for low, moderate, and high temperature plasmas.

constraints and a kinetic reconstruction that does not have ion temperature constraints. The temperature profiles that result from these three EFIT results are shown in Fig. 3.18. The plasma used for this example was specifically chosen because the basic reconstruction (containing only magnetic probe and MSE data) shows an unusual amount of temperature mismatch [Fig. 3.18(a)]. As expected, the addition of ion temperature constraints significantly improves the temperature overlap [Fig. 3.18(b)], and this overlap is nearly replicated by the temperature profile that is created by the kinetic reconstruction [Fig. 3.18(c)].

3.3.3 Extension of PAAR analysis to high toroidal rotation regime

As mentioned in Section 3.3.1, Eqs. (3.15) and (3.16) come from the standard form of the equilibrium flow in a tokamak,$^{29}$

\[
V = \frac{K(\psi)B}{n} + R\Omega\hat{\phi},
\]  

(3.18)
Figure 3.18: Ion temperature profiles with high- \((R < R_o)\) and low-field side \((R > R_o)\) measurements for three different magnetic equilibrium reconstructions: (a) a basic reconstruction with magnetic probe and MSE constraints, (b) a reconstruction with magnetic probe, MSE, and ion temperature constraints, and (c) a kinetic reconstruction that does not have ion temperature constraints. The temperature overlap achieved in (b) and (c) is similar, demonstrating the efficacy of the ion temperature constraints.

which is derived under the assumption that toroidal rotation is small, \(|V_\varphi| \ll V_{th}\), and hence that density is a flux function, \(n = n(\psi)\). For an impurity ion, this assumption is often violated. Even when considering plasmas without significant neutral beam injection (i.e. plasmas that exhibit intrinsic rotation), the ratio of the toroidal rotation to the thermal speed for carbon can approach 20% (See Fig. 1 of Ref.11). To ensure that PAAR measurements are accurate when the rotation speed approaches the thermal speed, the functional form of tokamak flow was derived without assuming the toroidal rotation to be low, i.e. \(|V_\varphi|/V_{th} = O(1)\). The results of that work\(^{62}\) are outlined in this section.

Using the first two moments of the kinetic equation,

\[
\nabla \cdot n \mathbf{V} = 0, \tag{3.19}
\]

\[
m \mathbf{V} \cdot \nabla \mathbf{V} + \frac{\nabla P}{n} - Ze(\mathbf{E} + \mathbf{V} \times \mathbf{B}) = 0, \tag{3.20}
\]

where \(n\) is the ion density and \(Ze\) is the ion charge. The time derivative,
particle source, friction force and source terms are second order in the ratio of the gyroradius to the gradient scale length, \( \rho_g/L \) and, therefore, ignored. With interspecies friction being small, these equations are valid for a multi-ion species plasma. As in Hinton and Wong,\(^6^3\) the divergence of the stress tensor is neglected. Equations (3.19) and (3.20) will be expanded in orders of \( \rho_g/L \) and solved.

With \( \mathbf{B} \) [given in Eq. (1.1)] being the \(-1\)th order, \( n^{(0)}V^{(0)} = 0 \) and \( \nabla \Phi^{(-1)} = V^{(0)} \times \mathbf{B} \), with \( -\nabla \Phi \) being the electric field. Taking \( \Phi^{(-1)} = \Phi^{(-1)}(\psi) \) and \( V^{(0)} = R\Omega^{(0)}(\psi)\hat{\psi} \), as in Hinton and Wong,

\[
\Omega^{(0)}(\psi) = \frac{\partial \Phi^{(-1)}}{\partial \psi}.
\]  

In the next order, the continuity equation is \( \nabla \cdot n^{(0)}V^{(1)} = 0 \), and the momentum equation is

\[
-\frac{1}{2}m\Omega^{(0)}R^2 \nabla T + \frac{\nabla P^{(0)}}{n^{(0)}} + Ze\nabla \Phi^{(0)} = ZeV^{(1)} \times \mathbf{B}.
\]  

Using \( T = T(\psi) \) and integrating the vector dot product of \( \mathbf{B} \) and Eq. (3.22) yields

\[
\ln n^{(0)} + \frac{Ze\Phi^{(0)}}{T} - \frac{m\Omega^{(0)}R^2}{2T} = H(\psi),
\]  

where \( H(\psi) \) is the integration constant. By taking the flux surface average of Eq. (3.23), \( H(\psi) \) can be related to the other variables in the problem. After defining \( N(\psi) \) from \( \ln N(\psi) \equiv \ln n^{(0)}(\psi) \), the zeroth order density can be written as

\[
n^{(0)} = N(\psi) \exp \left[ -\frac{Ze}{T}(\Phi^{(0)} - \langle \Phi^{(0)} \rangle) + \frac{m\Omega^{(0)}R^2}{2T}(R^2 - \langle R^2 \rangle) \right].
\]  

Substituting Eq. (3.24) into Eq. (3.22) yields

\[
V^{(1)} \times \mathbf{B} = \nabla \langle \Phi^{(0)} \rangle + \frac{1}{Ze} \nabla T + \frac{T}{Ze} \nabla \ln N(\psi) - \frac{m\Omega^{(0)}R^2}{2Ze} \nabla \langle R^2 \rangle
\]

\[- \frac{1}{Ze} \ln \left( \frac{n^{(0)}}{N(\psi)} \right) \nabla T + \frac{m\Omega^{(0)}}{Ze} (R^2 - \langle R^2 \rangle) \nabla \Omega^{(0)}.
\]  

Note that all gradients in Eq. (3.25) operate on functions of \( \psi \) so the direction of \( V^{(1)} \times \mathbf{B} \) is \( \nabla \psi \).
Using the continuity equation and axisymmetry, a general solution for the first order particle flux can be written as

\[ n^{(0)} V^{(1)} = \nabla G \times \nabla \varphi + F \hat{\varphi}, \quad (3.26) \]

where \( G \) and \( F \) are scalar functions. Equation (3.26) can be used to write

\[ n^{(0)} \nabla \times B = - \frac{I R^2}{R^2} \nabla G - \left[ (\nabla G \times \nabla \varphi) \cdot \nabla \psi \right] \nabla \varphi + \frac{F}{R} \nabla \psi. \quad (3.27) \]

In order for Eqs. (3.25) and (3.27) to be consistent, we must have \( G = G(\psi) \). This means that the \( \nabla G \times \nabla \varphi \) term in Eq. (3.26) represents the poloidal velocity and can be written as \( K(\psi) B_{\theta} \). Representing \( F \) so that the remaining algebra is simplified, the first order particle flux can be rewritten as

\[ n^{(0)} V^{(1)} = K(\psi) B + n^{(0)} R \Omega^{(1)} \hat{\varphi}. \quad (3.28) \]

Substituting this result into Eq. (3.25) gives

\[
\Omega^{(1)} = \frac{\partial \langle \Phi^{(0)} \rangle}{\partial \psi} + \frac{1}{Ze} \frac{\partial T}{\partial \psi} + \frac{T}{Ze N} \frac{\partial N}{\partial \psi} - \frac{m \Omega^{(0)2}}{2Ze} \frac{\partial \langle R^2 \rangle}{\partial \psi} \\
- \frac{1}{Ze} \ln \left( \frac{n^{(0)}}{N} \right) \frac{\partial T}{\partial \psi} + \frac{m \Omega^{(0)}}{Ze} (R^2 - \langle R^2 \rangle) \frac{\partial \Omega^{(0)}}{\partial \psi}. \quad (3.29)
\]

Note that the first four terms on the right hand side of Eq. (3.29) are flux functions while the last two vary poloidally.

Adding the results for the two orders, the final form for the velocity is

\[ V = \frac{K(\psi)}{n^{(0)}} B + R \left[ \Omega^{(0)}(\psi) + \Omega^{(1)} \right] \hat{\varphi}. \quad (3.30) \]

This result is accurate to first order in \( \rho_g / L \).

In order to recast these results in a form more applicable to PAAR analysis, \( \Omega^\dagger \) is defined as \( \Omega^{(0)} + \langle \Omega^{(1)} \rangle \), and Eq. (3.30) is rewritten as

\[ V = \frac{K(\psi)}{n^{(0)}} B + R \left[ \Omega^\dagger(\psi) - \frac{1}{Ze} \ln \left( \frac{n^{(0)}}{N(\psi)} \right) \frac{\partial T}{\partial \psi} + \frac{m}{2Ze} (R^2 - \langle R^2 \rangle) \frac{\partial \Omega^{(0)2}(\psi)}{\partial \psi} \right] \hat{\varphi}. \quad (3.31) \]

Note that \( \Omega^\dagger \) and \( \Omega^{(0)} \) differ by terms that are first order in \( \rho_g / L \). Since the term containing \( \partial \Omega^{(0)2} / \partial \psi \) is already first order in \( \rho_g / L \), first order accuracy
can be maintained by using $\Omega^\dagger$ instead of $\Omega^{(0)}$. Furthermore, an equivalent form can be written by changing the integration constant in Eq. (3.25) to be based on values of position and density at the outboard midplane, $n_p$ and $R_p$, respectively. Making these two modifications, an equation for the first order tokamak flow that can be used for PAAR analysis is

$$V = \frac{K(\psi)}{n^{(0)}} B + R \left[ \Omega^*(\psi) - \frac{1}{Ze} \ln \left( \frac{n^{(0)}}{n_p} \right) \frac{\partial T}{\partial \psi} + \frac{m}{2Ze} \left( R^2 - R_p^2 \right) \frac{\partial \Omega^*}{\partial \psi} \right] \hat{\phi}. \quad (3.32)$$

Equation (3.32) reduces to Eq. (3.18) when $|V|/V_{th} \ll 1$, but can be used for plasmas with a whole range of $|V|/V_{th}$. In addition, the inclusion of the measurements of $n^{(0)}$ enable PAAR analysis when density is not a flux function, whether that be a result of centrifugal effects or poloidal variation in the electrostatic potential.\textsuperscript{64}

Locating $n_p$ and $R_p$ on the outboard midplane is purposely done because this is the region where the highest density of CER measurements are made. By construction then, the new terms on the right-hand side of Eq. (3.32) are zero at the outboard midplane, simplifying the calculation of poloidal rotation once $K$ is found. The fitting of line of sight data is much the same as when using the zeroth order flow equations. Explicitly reforming Eq. (3.32) so that it describes the line of sight measurements and includes the cross section effect, results in

$$V_{\text{LOS}} = \frac{K(\psi)}{n^{(0)}} B_\phi s_\phi +$$

$$R \left[ \Omega^*(\psi) - \frac{1}{Ze} \ln \left( \frac{n^{(0)}}{n_p} \right) \frac{\partial T}{\partial \psi} + \frac{m}{2Ze} \left( R^2 - R_p^2 \right) \frac{\partial \Omega^*}{\partial \psi} \right] s_\phi + \alpha \cos \beta. \quad (3.33)$$

Using the first order flow equation to perform PAAR analysis is routine and uses the temperature and density measurements from the same CER view chords that provide the measurements of toroidal rotation. The fitting procedure is similar to the procedure used to fit data with Eq. (3.17), with $K$ and $\Omega^*$ taking the place of $k$ and $\omega$, respectively. The procedure is slightly more complicated, though, because $\Omega^*$ enters into the equation in two different ways. The term that depends on $\partial \Omega^*/\partial \psi$ is smaller order than the term that is linear in $\Omega^*$. As a
result, it has been found that it is sufficient to fit rotation data in an iterative way. After using a fit to the toroidal rotation measurements on the outboard midplane to generate an initial guess for $\Omega^*$, the $\partial \Omega^*/\partial \psi$ term is held constant and a new value of $\Omega^*$ is determined. This value is used to generate an updated value for the derivative term, and the process repeats until the value of both terms has converged. All PAAR results shown in this dissertation are analyzed with this method, aside from the following figure where PAAR results using the zeroth and first order flow equations are compared.

In Fig. 3.19, the varying values of $|V_\varphi|/V_{th}$ on the inferred $V_\theta$ are compared for three cases. The first analysis presented is from a H-mode plasma that has balanced neutral beam injection ($B_\varphi = -1.7$ T, $I_p = 0.7$ MA, $n_e = 3 \times 10^{19}$ m$^{-3}$, $P_{ECH} = 0.9$ MW, $\tau_{NBI} = 0$ Nm). Though the net neutral beam torque is negligible, $|V_\varphi|/V_{th} \approx 0.2$ for this plasma because of intrinsic rotation and a relatively low core ion temperature of 2 keV. The density asymmetry [ratio of high-field side (H) to low-field side (L) density] measured in this example is relatively extreme; its largest deviation from unity is $n_H/n_L \approx 0.5$. Figure 3.19(a) shows the PAAR results using the low and high rotation formulations of the flow for this example, and, despite this large density asymmetry, it can be seen that the difference between the two analysis techniques is within the measurement errors and less than 0.5 km/s across the whole measurement range. A second H-mode plasma that has unbalanced neutral beam injection is presented next. For this plasma ($B_\varphi = -2.15$ T, $I_p = 1.2$ MA, $n_e = 4 \times 10^{19}$ m$^{-3}$, $P_{NBI} = 6$ MW, $\tau_{NBI} = 4$ Nm), $|V_\varphi|/V_{th} \approx 0.6$ and the density asymmetry is as large as 0.7. As seen in Fig. 3.19(c), the difference between the two analysis techniques is larger than in the previous case, but still not large enough to exceed the random measurement error. Finally, a QH-mode example ($B_\varphi = -2.0$ T, $I_p = 1.3$ MA, $n_e = 1 \times 10^{19}$ m$^{-3}$, $P_{NBI} = 3$ MW, $\tau_{NBI} = 2$ Nm) is presented for which $|V_\varphi|/V_{th} \approx 0.8$ and the density asymmetry is as large as 0.6. As seen in Fig. 3.19(c), the difference between the two analysis techniques is as large as $\approx 2$ km/s across the measurement range and exceeds the $\approx 1$ km/s measurement error.

The results on the left side of Fig. 3.19 show that substantial differences
Figure 3.19: Differences between the full PAAR analysis [using Eq. (3.32)] and PAAR analysis that assumes toroidal rotation to be low [using Eq. (3.18)] and the carbon Mach number for three different plasmas: (a-b) a low-rotation H-mode, (c-d) a high rotation H-mode, and (e-f) a high rotation QH-mode. On the left side the two different analysis results are shown, and on the right side the full PAAR analysis is done while ignoring the $\partial \Omega^*/\partial \psi$ or $\ln(n^{(0)}/n_p)$ terms.
in the results of the full analysis and the low rotation formulation exist in cases where the rotation speed approaches the thermal speed. To understand what produces these differences, the effects of the various differences between Eqs. (3.18) and (3.32) are considered. These are the presence in Eq. (3.32) of the poloidally varying density \( n(0) \) in the denominator of the first term and the presence of the two poloidally varying terms in the brackets in the second term.

In order to see which of these is most important, modifications of the full analysis are shown on the right side of Fig. 3.19. First the term containing the logarithm of the density ratio is ignored, and second the term containing the derivative of \( \Omega^* \) is ignored. As can be seen there, dropping each of these terms makes a comparable change in the result. Interestingly, the signs of the effects of the two terms are opposite (as a result of the fact that density is larger on the low-field side for these cases). However, neither of these terms by itself explains all of the differences shown on the left side of Fig. 3.19. Accordingly, it must be true that the presence of \( n(0) \) in the denominator of the first term in Eq. (3.32) is somewhat more important than those other two terms for these cases. The poloidal variation of the density is related to the centrifugal term in Eq. (3.20); however, the connection between toroidal rotation and density asymmetry is complicated by the creation of a poloidally varying electrostatic potential due to trapped particles.\(^ {64} \)

Though this analysis is used for cases where the carbon rotation is near its thermal speed, the magnetic equilibrium reconstructions used do not include rotation effects. This is because the rotational kinetic energy is too small to affect the overall equilibrium.\(^ {65} \) Even when the carbon rotation speed is close to thermal, the deuterium rotation speed is not due to the mass difference between the two species. The equilibria used satisfy the condition \( T = T(\psi) \), indicating that neglect of the effect of rotation on the equilibrium is appropriate.

### 3.3.3.1 Effect of high rotation on the measurement of the cross section effect

High plasma rotation can also affect the determination of the cross section effect with co- and counter-current neutral beams, as described in Sec. 3.1.4. This
is due to the fact that the interaction energy between the neutral beam particles and the plasma ions is different for the co- and counter-current neutral beams when the plasma rotation is a significant fraction of the beam neutral velocity (typically 2800 km/s for the full-energy component). In order to arrive at Eq. (3.8), it was necessary to assume that the cross section effect was the same for the matched views of the co- and counter-current neutral beams. However, when the rotation is large enough, the value of $\alpha$, as well as the direction of the cross section effect change for these two views. The direction of the cross section effect is $V - V_b/|V - V_b|$, but, if it is assumed that $|V| \ll |V_b|$ [as done in Eq. (3.7)], this direction is simply $\hat{V}_b$.

To determine the significance of this effect for highly rotating DIII-D plasmas, an atomic physics simulation capable of accounting for the effect of the plasma rotation was used. This simulation is a collisional radiative model whose atomic physics is fully resolved in quantum numbers $n$ and $l$. Given a measurement of the line of sight velocity and the plasma parameters necessary for modeling the neutral beam and atomic physics (electron density, electron temperature, ion temperature, and impurity density), this simulation is able to determine what value of $C^{+6}$ rotation would produce this measurement of the $C^{+5}$ velocity. The difference between these measurements is due to the cross section effect. By artificially scaling the input measurement of line of sight velocity, a dependence of the cross section effect on the $C^{+6}$ rotation can be found.

For this investigation, a high temperature QH-mode discharge was chosen (149094, the same as was used in the analysis shown in Fig. 3.19). The high ion temperature (10 keV) in this discharge ensures that the cross section effect will be significant for the measurement of the large rotation present (350 km/s). Results from the collisional radiative model showed that determining $\alpha$ with co- and counter-current neutral beams has a maximum error of $\approx 14\%$ for this discharge. This is certainly significant for the determination of $\alpha$, but, as seen in, this has a small effect on the determination of the $C^{+6}$ rotation. This is because the 14% error is in a correction term that is 10-20% of the measurement for high rotation cases. As a result, the effect on the measurement of $C^{+6}$ rotation is only 1-3%.
To determine the effect that this error can have on PAAR measurements, it is instructive to consider a basic representation of how poloidal rotation is determined with the PAAR method. Fundamentally, the measurement is made by determining the difference between angular rotation on the high- and low-field side (H and L) of the tokamak. Working from Eqs. (3.16) and (3.17),

\[
\frac{V_{\theta,L}}{B_{\theta,L}} \left( \frac{B_{\phi}}{R} \right)^L_H = \left( \frac{V_{\text{LOS}}}{s_{\phi,R}} \right)^L_H - \frac{\alpha \cos \beta}{s_{\phi}}^L_H. \tag{3.34}
\]

Here it can be seen that it is the difference between the cross section effect on the high- and low-field side that is important. Note that this is true even if all the extra terms in the high rotation formulation of the flow equations [Eq. (3.33)] were included. If all cross section effect calculations were off by a constant amount, there would be no effect on the PAAR results. The difference between the high- and low-field side cross section effects is most dramatic for the PAAR measurements that are furthest from the core. These measurements can have significantly different values of \( \alpha \) because of the changing beam density fractions, while differences in \( \alpha \) in the region immediately surrounding the magnetic axis is minimal because the beam density fractions are very similar.

Using the same QH-mode discharge as a test case, the collisional radiative model calculates a 1 km/s error in the measurement of the value of the second term on the right-hand side of Eq. (3.34). This value is deemed insignificant because the difference between any two toroidal rotation measurements cannot credibly be determined to 1 km/s, since this is the level of error in one source of rotation measurement error, the determination of the wavelength fiducial (see Sec. 3.1.3.1). Note that the error in PAAR measurements can be less than this because of toroidal rotation is multiplied by a factor of \( B_{\theta}/B_{\phi} \) (typically 0.1-0.25) in the equation for poloidal rotation. Said another way, a 1 km/s difference in toroidal rotation measurements on the high- and low-field sides of a flux surface produces a poloidal rotation much smaller than 1 km/s.
3.3.4 Comparison of PAAR and vertical CER measurements

Insights into the advantages of the PAAR method can be found by comparing results from the PAAR method and the vertical CER measurements of poloidal rotation. As an example, comparisons between the PAAR and the vertical CER measurements for an L-mode and H-mode plasma are shown in Fig. 3.20. Here it can be seen that, while the measurements are consistent with each other, the PAAR results have significantly smaller uncertainty.

![Figure 3.20: Poloidal rotation measured with vertical CER and the PAAR method for an L-mode and H-mode discharge. The agreement between the two measurements is good, but the uncertainty in the PAAR measurement is lower.](image)

However, agreement between the two measurement techniques is not constant. Nevertheless, the PAAR results are consistent with the vertical CER measurements. To see this, consider a comparison between the two techniques for a H-mode plasma at a slightly different time than the one shown in Fig. 3.20(b). As shown in Fig. 3.21, there is a small region of clear disagreement between the two measurements. However, since there is some uncertainty in the determination of the finite excited state lifetime, $\tau$, with the technique for measuring poloidal rotation with the vertical CER measurements, the best way to compare PAAR
results with vertical CER measurements is to see if the PAAR results are consistent with the vertical line of sight data. That is, the simplest comparison is one which determines if there is a $\tau$ for which the PAAR method and the vertical CER measurements are consistent.

In order to perform this comparison, PAAR results are decreed to be the true poloidal rotation in the vertical CER analysis routine that can produce a measurement of the poloidal rotation. Rather than being allowed to vary $\tau$ and the poloidal rotation in order to get the best fit to the vertical CER data, this routine must keep poloidal rotation constant while varying $\tau$ and the fit to the tangential rotation data. The difference between the measured and predicted vertical CER measurements then provides information about the consistency of the two measurements. If the pattern of these residuals is consistent with random error, the two measurements of poloidal rotation are in agreement, when the value of $\tau$ is not the value determined by the vertical CER analysis routine. The result of this analysis for the PAAR results shown in Fig. 3.21 are shown in Fig. 3.22,
where it can be seen that the residuals show no significant pattern.

Figure 3.22: Residuals of a fit to the vertical CER measurements when the poloidal rotation is fixed at the result from the PAAR method. Though the vertical CER analysis routine yields a slightly different measurement of poloidal rotation, the PAAR results are consistent with the vertical CER measurements when the value of $\tau$ is reduced by 39%.

These results show that, even for a case where the PAAR and vertical CER analysis appear to disagree, the PAAR results are still consistent with the vertical CER data. Due to the fact that the way the PAAR method accounts for the cross section effect is superior to the way it is accounted for with the vertical CER analysis, the PAAR results are considered more accurate in these cases. In this particular case, the PAAR results, after being used in the modified analysis of the vertical CER data, suggest that $\tau$ has a value of $1.4 \pm 0.3$ ns while the original vertical CER analysis determined $\tau$ to be $2.3 \pm 0.3$ ns.

The vertical CER analysis is aided by vertical CER measurements near the magnetic axis. At this location, the poloidal rotation must be zero, and this aides the ability of the vertical CER analysis to determine the value of $\tau$. In addition, the vertical CER analysis does not account for the line integrated nature of the CER measurement. Vertical CER view chords on DIII-D average over more flux
surfaces than tangential view chords because the neutral beams are approximately 2.5 times as tall as they are wide. For this reason, PAAR analysis has better spatial resolution in the core of the plasma than the vertical CER analysis. Also note, as seen in Fig. 3.22, even the vertical view chord that is closest to the axis is a significant distance away for a case where PAAR and vertical CER analysis disagree.

The spatial averaging of vertical CER measurements can prevent good measurements of rotation at the magnetic axis when gradients in the core are large, as is the case for plasmas with an internal transport barrier (ITB). Measuring poloidal rotation in ITB plasmas is discussed in Chapter 5, and in Section 5.3 the comparison between PAAR and vertical CER analysis is revisited for this special case.

Due to the increased spatial resolution and lack of atomic physics calculations to account for the cross section effect, the PAAR method is the preferred method for measuring poloidal rotation in the core of the plasma. However, the experimental conditions required for PAAR analysis are more complicated than those required for the vertical CER measurements. This is mostly due to the neutral beam requirements. PAAR analysis requires co- and counter-current neutral beam injection that enables good subtraction to obtain high quality analysis, but most DIII-D experiments use a single direction of neutral beam injection.

As a result, there are instances where a DIII-D experiment contains interesting poloidal rotation data, but only direct measurements of poloidal rotation are available. Such instances are always noted in the accompanying text, and only two such instances exist (both in Chapter 6). Even in these instances, though, an attempt to verify the vertical CER measurements with PAAR analysis is made. For one of these cases where vertical CER measurements are needed to make the comparison to theory, it is possible to compare the two methods at a time that is not instructive for the theoretical comparison, but near to the times that are. This comparison is presented here as another case where agreement is seen between the two methods for a H-mode plasma. This agreement can be seen in Fig. 3.23 where it is again clear that the uncertainty in the vertical CER analysis
is larger than the uncertainty in PAAR analysis.

![Graph](image)

**Figure 3.23:** Poloidal rotation measured with vertical CER and the PAAR method at a time when a discharge is transitioning from being primarily heated with NBI to primarily heated with ECH. This is a H-mode discharge, and the agreement between the two methods is good. The PAAR method has superior accuracy.

### 3.4 Acknowledgements

Chapter 3 contains material from Review of Scientific Instruments vol. 81. Chrystal, Colin; Burrell, Keith H.; Pablant, Novimir A., American Institute of Physics, 2010. The dissertation author was the primary investigator and author of this paper.

Chapter 3 contains material from Review of Scientific Instruments vol. 83. Chrystal, Colin; Burrell, Keith H.; Grierson, Brian A.; Groebner, Richard J.; Kaplan, David H., American Institute of Physics, 2012. The dissertation author was the primary investigator and author of this paper.

was a co-author on this paper.

Chapter 3 contains material from Physics of Plasmas vol. 21. Chrystal, Colin; Burrell, Keith H.; Grierson, Brian A.; Staebler, Gary M.; Solomon, Wayne M.; Wang, Weixing X.; Rhodes, Terry L.; Schmitz, Lothar; Kinsey, Jon E.; Lao, Lang. L; deGrassie, John S.; Mordijck, Saskia; Meneghini, Orso, American Institute of Physics, 2014. The dissertation author was the primary investigator of this paper.
Chapter 4

In-situ measurements of diagnostic geometry

This chapter describes a new method for determining the viewing geometry of the CER diagnostic. A basic implementation of this method has been done before, but this work significantly expands the data set and the number of geometric parameters that are determined. This new method adds complexity to the spatial calibration with the aim of removing uncertainty from important geometric parameters that are difficult to determine with standard spatial calibration methods. This work is motivated by the fundamental importance of the geometry of the CER view chords when determining poloidal rotation with the PAAR method.

The PAAR method is based upon measuring asymmetry in toroidal rotation. As a result, the position of CER view chords does not just need to be known well in an absolute sense, but in a relative sense. A small error in, for instance, the position of a neutral beam may only cause a small error in the absolute position of the intersection between the view chords and the beam but can cause a larger error in the relative position of two view chords on either side of the magnetic axis. The geometric dependencies of the PAAR method can be seen by examining Eq. (3.17) or (3.33). The major radius ($R$), relation between the view chord and the toroidal direction ($s_\phi$), and the projection of the neutral beam injection direction onto the view chord ($\cos \beta$) are all geometric parameters...
embedded in PAAR analysis.

4.1 Basic spatial calibration technique

The basic spatial calibration procedure for the CER system was discussed in Section 3.1.3.1, and this section presents more details of this procedure that are relevant to the current chapter. For clarity, some terminology of the machine design is presented. The eight neutral beams on DIII-D are located in four neutral beamlines that are labeled according to the name of the port on the machine through which they pass. As a result, there exist “30”, “150”, “210”, and “330” neutral beams on DIII-D. For each beamline there is a neutral beam whose injection direction is more perpendicular to the toroidal direction and a neutral beam whose injection direction is more tangential to the toroidal direction. These two sources are labeled as “LT” (for left) and “RT” (for right) according to the following convention: when viewing the neutral beamline from inside the tokamak, the neutral beam that is to the left is the “LT” source and the neutral beam that is to the right is the “RT” source. Hence there are “30LT” and “210RT” neutral beams. The “LT” neutral beams are the more tangential source for every beamline except the 210 beamline, which is the counter-current beamline. For reference, see the annotated plan view of the tokamak in Fig. 4.1.

The primary measurements needed for PAAR analysis are the tangential CER measurements on the 30 and 210 neutral beams. The spatial calibration developed in this chapter concerns only these views, which originate from two ports, 315T0 and 165R0. Each of these ports has a single lens, and it is assumed that each view chord that uses a lens passes through a single point that is the lens center. The lens center is one point used to define a line that represents each CER view chord. The other point is derived from the measurements made on the calibration targets. For determining the CER geometry, the effect of mirrors is ignored, so, as seen in Fig. 4.1, the position indicated for the 165R0 lens is where the view chords would converge if there were no mirror.

The calibration targets are metal targets that are placed in the machine
Figure 4.1: An plan view of the DIII-D tokamak with various objects important for the spatial calibration labeled.

at the nominal locations of their respective neutral beams. These locations are determined by two fiducial markings: a mark in the drift duct that is directly below where the two neutral beams in one beamline cross, and a mark on the centerpost where the perpendicular neutral beams hits the centerpost (30RT and 210LT for the relevant neutral beams). The angle between the two neutral beams is a design parameter based on the location of the beam cross over and the location of the ion sources at the far end of the neutral beamlines. This angle is built into the targets so that, once the perpendicular portion of the target is placed according to the fiducials, the location of the tangential portion of the target is set as well.

When light is passed through the fiber optics for each view chord into the
vessel, an image is formed on the target, and this image is then marked on the target. Once removed from the vessel, very accurate measurements of the position of each view chord on the target can be made. Taking $t_c$ to be the position of the target crossover point (same as beam crossover point), $x$ to be the distance between the crossover and the center of the view chord’s image, and $\hat{V}_b$ to be the neutral beam injection direction, the position of a view chord on a neutral beam can be written as

$$r = t_c + x\hat{V}_b.$$  

For this work, cylindrical coordinates $(R, \varphi, z)$ are used, and it is elementary to obtain $R$ and $z$ for a view chord from $r$.

With the lens position and the position of a view chord on a target, it is possible to construct the viewing vector,

$$\hat{s} = \frac{l - r}{|l - r|},$$  

where $l$ is the position of the lens center. By convention, the viewing vector points in the direction that photons travel from the tokamak to the lens. From $\hat{s}$, $\hat{s}_\varphi$ is readily obtained, and $\cos \beta \equiv \hat{s} \cdot \hat{V}_b$.

4.1.1 Uncertainty in basic spatial calibration

Though the method described above is sufficient to completely define the geometry of each CER view chord, uncertainties in key parts of the calculation deserve to be scrutinized. First, it is assumed that the position of the neutral beams, $(t_c)$ and the beam injection direction $(\hat{V}_b)$ are well known. Though extensive measurements can be made of the neutral beam ion sources and their relationship to the tokamak, measuring the neutral beam position when it is actually injecting power is much more difficult. Of course, the CER view chord geometry is ultimately interested in the position of the neutral beam when it is injecting power.

The second uncertainty is related to the difficulty in determining the position of the CER lenses $(l)$. On DIII-D, this is a result of the fact that the CER lenses are outside the vacuum vessel and cannot be accessed directly from within the vessel. Lens positions can be determined from design specifications
of ports, but this process tends to accumulate uncertainty as the information on the scale of the machine size is propagated to the mounting hardware for a 2” diameter lens. It is also possible to determine the lens location from multiple target measurements for each view chord. With this procedure, more than one measurement of a view chord’s position can be used to determine their sightlines, and the position that best matches the location where these sightlines converge is taken to be the lens location. This method has a fundamental limitation though: when view chord positions are measured near to the neutral beam they view, significant extrapolation is required to determine the location of the lens some distance away, and when view chord positions are measured away from the neutral beam they view, their images are unfocused.

In Fig. 4.2, these two uncertainties are sketched. Uncertainty in the lens location affects $\hat{s}_\varphi$ and $\cos \beta$ while uncertainty in the beam location affects these parameters as well as $R$. The reason for developing the new spatial calibration technique described in this chapter is to reduce these uncertainties with complementary measurements.

**Figure 4.2:** An plan view of a portion of the DIII-D tokamak depicting uncertainties (in black, exaggerated) in the CER lens location, view chord angle, neutral beam crossover location, and neutral beam injection direction.
4.2 Integrated spatial calibration

In order to remove the uncertainties in the basic spatial calibration, new geometric information was added and then the uncertain parameters were allowed to vary. Essentially, if the lens and beam geometry are no longer fixed to their nominal values, new measurements that contain information about the beam and view chord geometry must be added so that the CER geometry can be completely constrained. This new method for performing the spatial calibration is referred to as an “integrated” approach because of the inclusion of multiple types of geometric measurement within one spatial calibration.

This integrated approach parameterizes the view chord and beam geometry and then uses a least-squares fit to determine what set of parameters best fits the provided measurements. The coordinate system of the fit is Cartesian \((u, v)\), and final results are converted to cylindrical. In the integrated approach, only tangential view chords are used so the elevation of the CER lenses and neutral beams is the same as the tokamak midplane. In this case, it is reasonable to ignore the near zero \(z\) components of \(\hat{s}\) and \(\hat{V}_b\). Each neutral beam is parameterized by its intercept with the \(v\)-axis \((v_0)\) and an angle between the \(u\)-axis and the beam injection direction \((\theta_b)\). Each CER view chord is parameterized by the position of its lens center \((l_u, l_v)\) and an angle in the tokamak midplane \((\theta_c)\). This results in each pair of neutral beams having 4 parameters, and, working just with view chords from 315T0 and 165R0, \(n + 4\) parameters for the \(n\) view chords.

Three types of geometric information are used in the integrated approach: target measurements, measurements of Doppler shifted beam emission, and measurements of Stark split beam emission. The latter two are discussed in detail in the next section. Most importantly, these two measurements provide \textit{in-situ} measurements of the neutral beams, i.e. measurements of the beam positions when they are injecting power. In the integrated approach, the target positions are no longer tied to the beam positions, which are allowed to vary.

The target positions are set by measurements made by a coordinate measurement machine, and the target measurements provide information about points in the plasma that each view chord should pass through. For every view
chord that is part of this analysis, target measurements are available at two positions, either the nominal positions of the 30RT and 30LT neutral beams (for 315T0 view chords), or the nominal positions of the 210RT and 210LT neutral beams (for 165R0 chords). Note that for each view chord, a lens center and angle could be constructed that would match the two target measurements exactly. However, when all view chords that use the same lens are constrained to use the same lens center, this no longer becomes possible. In fact, with reasonable errors attributed to the target measurements, the lens center has a significant amount of freedom due to the fact that it is $\approx 2$ m away from the targets which are only $\approx 0.1$ m apart.

### 4.2.1 Doppler shift measurements

When neutral beams are injected into a tokamak filled with cold gas (a beam into gas shot, see Sec. 1.2.3), such as neutral deuterium, four clear sources of emission near the $D_\alpha$ spectral line ($6561 \text{ Å}$) are visible. Excitation of the cold gas yields an unshifted $D_\alpha$ emission, and the three energy components of neutral beam also emit $D_\alpha$ emission, but their emission is Doppler shifted due to their significant velocity (for typical accelerator voltages of 81 kV, the full, half, and third energy beam components have velocities of 2800, 2000, and 1600 km/s respectively).

By measuring the wavelength shifts between the beam neutral emission and background gas emission, information about the angle between a view chord and the beam injection direction can be determined. This can be seen in the equation for the Doppler shift for the full energy beam component:

$$\Delta \lambda_{\text{Doppler}} = \lambda_0 \sqrt{\frac{2eE_{\text{beam}}}{m_Dc^2}} \cos \beta,$$

(4.3)

where $\lambda_0$ is the rest wavelength, $E_{\text{beam}}$ is the accelerator voltage and $m_D$ is the mass of deuterium.

In principle, each view chord can make three Doppler shift measurements for each neutral beam, one for each beam species. In practice though, the half and third energy components can be well blended for some view chords due to small values of $\cos \beta$. This blending increases uncertainty in their location. To avoid
placing undue weight on the Doppler shift measurements for view chords that can make measurements of each beam species, only measurements of the full energy component are used in the integrated approach.

An example of the observed spectra, and the fit to the spectra is shown in Fig. 4.3. The fit quality is high, so the precision of these measurements is good. For a typical beam into gas shot, a single 50 ms beam blip can be observed, and signal levels dictate that five 10 ms CCD exposures can be used to observe the resulting light. Assuming the beam output is relatively constant, the population standard deviation in these Doppler shift measurements are taken to be the error.

**Figure 4.3**: Example of a fit to the observed spectra near $D_\alpha$ (6561 Å) for a beam into deuterium shot (fit is in red and overlays data in black). The shorter wavelength emission is Doppler shifted $D_\alpha$ emission from the beam neutrals. Measurements of the wavelength shift of this emission contains information about the angle between the neutral beam and the view chord.
4.2.2 Stark splitting measurements

If the same measurements outlined in Sec. 4.2.1 are performed when the toroidal magnetic field is turned on, the Doppler shifted beam emission will also be Stark split. This stark splitting is a result of the electric field felt by the beam neutrals which is, in the lab frame, a \( \mathbf{V} \times \mathbf{B} \) force. Because the strength of the magnetic field scales as \( 1/R \), measurements of the Stark splitting contains information about the major radius where a view chord and neutral beam cross.

The dependence on major radius can be seen in the equation for the Stark splitting:

\[
\Delta \lambda_{\text{Stark}} = \frac{3}{2} \frac{e a_0}{h c} \left| \sqrt{\frac{2 e E_{\text{beam}} I_{\text{TF}} V_b \times \hat{\phi}}{m_D c^2}} \right| R, \tag{4.4}
\]

where \( a_0 \) is the Bohr radius, \( h \) is Planck’s constant, and \( I_{\text{TF}} \) is the total current through the toroidal field coils. A fit to the Stark split spectra is shown in Fig. 4.4. In this case, the measurement being made is of the splitting of the emission full energy beam component emission (compare to Fig. 4.3). As with the measurements of Doppler shifts, only information from the full energy component is used, and the population standard deviation for a set of measurement during a single beam blip is used as the error in the measurement.

The ability of the Stark splitting measurements to determine \( R \) and the Doppler shift measurements to measure \( \cos \beta \) constrain the geometry enough that the neutral beam and lens locations can be free parameters in the fit. Due to the fine structure of the Stark manifold, only the main-ion CER system\(^{22} \) can make accurate measurements of this effect. Main-ion view chords exist at both the 315T0 and 165R0 ports, and their measurement capabilities are an indispensable part of this integrated spatial calibration.

4.2.3 Results

The fit routine for the integrated method determines: \( \theta_c \) for 19 view chords that use the 165R0 lens (8 from the main-ion system) and 18 view chords that use the 315T0 lens (7 from the main-ion system), \( l_u \) and \( l_v \) for the 165R0 and 315T0 lenses, and \( v_0, \theta_b \), and \( E_{\text{beam}} \) for 5 neutral beams. To do this it uses: target
Figure 4.4: Example of a fit to the observed spectra near $D_\alpha$ (6561 Å) for a beam into deuterium shot with a strong toroidal magnetic field (fit is in red and overlays data in black). The Doppler shifted beam emission is Stark split due to the beam particle velocity and the magnetic field. Measurements of this splitting contains information about the major radius where the neutral beam and the view chord intersect.

data measurements from 30LT, 30RT, 210LT, and 210RT, CMM measurements of target positions $t_u$, and $t_v$, Doppler shift measurements of emission from 30LT, 30RT, 210LT, 210RT, and 330LT (which can be seen by 165R0 chords due to insignificant attenuation during beam into gas shots) beam neutrals, and Stark splitting measurements of emission from 30LT, 210LT, and 210RT beam neutrals. The Stark splitting measurements are the most challenging, so they are not available for 30RT and 330LT because the angles the view chord make with these beams is closer to perpendicular, reducing the Doppler shift of the beam emission and blending the stark manifolds of each beam species.

Results of the integrated method fit demonstrate that there is a consistent set of geometric parameters that is compatible with these various measurements. Minor adjustments to the free parameters yield a reasonable fit with reduced
$\chi^2 = 4.8$ (the fit has 64 free parameters and 140 degrees of freedom). Given the somewhat imprecise method for obtaining the errors for each of the complementary geometry measurements, this result is quite good. The biggest parameter change produced by this fit was in the 165R0 lens position. Cursory examination of beam emission Doppler shifts had previously suggested this lens location was inaccurate, and those suspicions were confirmed by the integrated spatial calibration. The lens moved by $\approx 12$ cm in the integrated method result. Though this is a significant distance, its movement was toward the beam, roughly aligned with the view chords. As a result, this change resulted in only relatively minor adjustments to the main important geometric parameters.

To understand the effect of the integrated method’s adjustment to the view chord and beam parameters, the original and geometric parameters and the geometric parameters produced by the integrated method are compared. The comparison of these original and “new” parameters is shown in Fig. 4.5. All the observed changes are relatively minor, showing that the original spatial calibration is consistent with the new Doppler shift and Stark splitting measurements.

The most important test of these changes is the effect they have on PAAR measurements. To test this, PAAR analysis was performed with the geometric parameters from the original and new spatial calibration and then compared. In Fig. 4.6, this comparison is shown for two discharges with large, oppositely directed poloidal rotation. The changes in the PAAR analysis associated with the new spatial calibration make little difference in the results. From the perspective of PAAR analysis then, these two spatial calibrations are equivalent. Note that PAAR analysis, among all types of CER analysis, is the most sensitive to the spatial calibration.

### 4.3 Conclusion

In this chapter, a new method for spatially calibrating the CER diagnostic was described. This method uses in-situ measurements of the neutral beams to remove uncertainty in the spatial calibration associated with the neutral beam
Figure 4.5: Comparisons of geometric parameters important to PAAR analysis determined by the original spatial calibration and the “new” integrated calibration method. Parameters from view chords that use both the 165R0 (Counter-views) and 315T0 (Co-views) lens are shown. Note that major radius changes in (a) are less than 1 cm, the range shown for the $s_\varphi$ changes in (b) is reduced and no change is larger than 2%, and the changes in $\cos \beta$ are largest near zero, where their effect on PAAR analysis is minimized.

Figure 4.6: PAAR poloidal rotation measurements made when using the original and new spatial calibration results. For strong positive and negative poloidal rotations, no significant change is seen in the results.
position and CER lens location. The results of this new spatial calibration are very similar to the original spatial calibration results, indicating that the nominal positions of the neutral beams was sufficiently accurate. Though these results were useful for ensuring that inaccuracy in the spatial calibration was not affecting the PAAR measurements of poloidal rotation, all PAAR analysis presented in this dissertation uses the original spatial calibration.

The accuracy of the spatial calibration is also very important for the analysis of vertical CER measurements. This is a result of the fact that the values of $s_\phi$ and $s_R$ are needed to determine the toroidal rotation pickup and cross section effect on the measurements of vertical rotation [see Eq. (5.2)], and both of these effects can be larger than the true vertical rotation. Unfortunately, even if this new spatial calibration approach were expanded to three dimensions, it could not be used to determine the geometry of the vertical view chords. For these view chords, both the Doppler shift and Stark splitting measurements are not available because the vertical view chords and neutral beam are nearly perpendicular to each other. The blending of the beam emission that results makes the Doppler shift and Stark splitting measurements effectively impossible.

4.4 Acknowledgements

Chapter 4 contains material from Review of Scientific Instruments volume 85. Chrystal, Colin; Burrell, Keith H.; Grierson, Brian A.; Lang, Lao L.; Pace, David C., American Institute of Physics, 2014. The dissertation author was the primary investigator of this paper.
Chapter 5

Poloidal rotation in the presence of a transport barrier

The poloidal rotation can be a critical part of the calculation of the $E \times B$ shear that enables the formation of transport barriers.\(^6\) Edge transport barriers, which define H-mode plasmas, are critical for tokamak fusion performance while internal transport barriers (ITBs) are typically not included in reactor scenarios because of their relatively low stability. Nevertheless, they provide an interesting plasma to study because the physics that controls the internal transport barrier is the same as the physics that controls the edge transport barrier.\(^6\) This chapter concerns measurements of poloidal rotation in plasmas with internal transport barriers. The consistency between the measurements of poloidal rotation described in this chapter and theories of poloidal rotation is discussed in Chapter 6.

5.1 ITB formation and description

ITB formation is usually accomplished by injecting a relatively large amount of power early in the discharge, during the current ramp. This increases the core plasma temperature and therefore delays current diffusion into the core. As a result, ITB discharges tend to have reversed magnetic shear, i.e. hollow current profiles where the poloidal magnetic field increases with radius in the region immediately around the axis. This is the opposite of standard discharges where
the current profile is peaked on the magnetic axis, and the poloidal magnetic field decreases with radius.

The reversed magnetic shear has a stabilizing effect\(^{68}\) that allows an \(\mathbf{E} \times \mathbf{B}\) shear feedback loop to initiate. On DIII-D, this feedback is typically propelled by the rotation. This distinction is important because, for a typical plasma configuration, the \(\nabla P\) and \(\mathbf{V} \times \mathbf{B}\) components oppose each other, so one of these radial forces must dominate in order for a significant \(E_r\) to develop. As seen in Eq. (1.3), the radial forces that results from co-current neutral beam injection, and ion-diamagnetic poloidal flow are anti-parallel with the pressure gradient force. Effectively, this means that the standard recipe for forming an ITB on DIII-D is to inject significant power and torque early in the discharge to create both the reversed magnetic shear and a significant gradient in \(V_\psi\) that can then lead to the formation of a transport barrier.

To maximize NBI torque in DIII-D, the six co-current neutral beams must be used as much as possible. However, in order to make PAAR measurements, counter-current beams are needed as well. To minimize the effect amount counter-current torque needed for PAAR measurements in an ITB discharge, the counter-current beams only inject “blips” of power. For the discharges shown in this chapter, a counter-current beam is run with a 25% duty cycle and a 40 ms period. Unfortunately, this prevents continuous measurements of poloidal rotation with the PAAR method. When the counter-current beam is turned on, a balancing co-current beam is also turned on to prevent a time varying torque from affecting the measurements.

An example of the development of an ITB plasma analyzed in this chapter is shown in Fig. 5.1. The formation of the ITB is near 1200 ms, and is marked by the sudden decrease in the gradient scale length for ion temperature and toroidal rotation in the core. These increases are due to the formation of a transport barrier. The spikes seen in \(P_{\text{NBI}}\) are due to the counter-current (and balancing co-current) beam blips needed for the PAAR measurements. For this discharge, \(B_\psi = -2.0\ T\), \(P_{\text{NBI}} = 5\ MW\), and \(\tau_{\text{NBI}} = 3.2\ Nm\).

The formation of this ITB can be seen more clearly in the evolution of the
Figure 5.1: Progression of a plasma discharge where an ITB forms. Near 1200 ms, the core gradients of $T_e$, $T_i$, and $V_\varphi$ suddenly increase, signifying the that the ITB has formed. This transport barrier suddenly relaxes before 1300 ms, but quickly reforms, as seen in the $T_i$ measurements. Different color traces represent measurements at different radii.

Figure 5.2: Profiles of (a) ion temperature and (b) toroidal rotation (measured on the outboard midplane) before and after ITB formation, showing a sharp rise in their gradients near $\rho = 0.3$. 
Figure 5.3: Profiles of (a) electron temperature, (b) electron density, and (c) carbon density before and after ITB formation, demonstrating that these parameters are also affected by the ITB.

$T_i$ and $V_\varphi$ profiles. In Fig. 5.2, it can be seen that the temperature and toroidal rotation profiles have steepened significantly after the ITB formation time for this discharge, $\approx 1200$ ms. The core ion parameters are quite large for DIII-D, peaking at 12 keV and 500 km/s, and only ITB plasmas can achieve these parameters while maintaining an L-mode edge. The formation time coincides with the minimum safety factor reaching 2, as has been previously observed.\textsuperscript{69}

Internal transport barriers on DIII-D can also affect electron and particle transport as well. To demonstrate this, profiles of electron temperature, density, and carbon density are shown in Fig. 5.3 for the same ITB discharge for which the ion temperature and rotation were shown. Clear rises in core gradients after the transport barrier forms can be observed. For clarity, only the spline fit to the respective measurements are shown.

Additional ITB discharges similar to the discharge depicted in Fig. 5.2 are discussed in this chapter. They represent relatively weak ITB plasmas, not obtaining the extreme parameters seen in some previous DIII-D ITB discharges.\textsuperscript{70} This is a result of experimental design considerations and neutral beam availability at the time of the experiment. Additional, large radius ITB discharges were created with significantly more NBI torque and power ($B_\varphi = -2.0 \ T$, $I_p = 1.3 \ MA$, $n_e = 2.5 \times 10^{19} \ m^{-3}$, $P_{NBI} = 7.0 \ MW$, $\tau_{NBI} = 5.6 \ Nm$). These discharges formed a transport barrier at an unusually large radius, $\rho \approx 0.65$. Though this prevented PAAR measurement from being made at the transport barrier, it did allow the
Poloidal rotation in the core of a transport barrier to be investigated. The $T_i$ and $V_\phi$ profiles before and after the transport barrier formation for discharges of this type are shown in Fig. 5.4. Nearly 60% of the plasma volume is contained within the barrier for these discharges, where $T_i \approx 10$ keV.

![Figure 5.4](image)

**Figure 5.4**: Profiles of (a) ion temperature and (b) toroidal rotation (measured on the outboard midplane) before and after ITB formation, showing a sharp rise in their gradients at a relatively large radius, $\rho \approx 0.65$.

## 5.2 PAAR measurements during ITB formation

The study of poloidal rotation for ITB plasmas has led to the expectation that poloidal rotation increases in the presence of an ITB$^{13,69,71}$ (though Ref. 13 predates the accounting for the gyro-orbit cross section effect,$^{26}$ the resulting radial electric field measurements are corroborated by MSE). The PAAR method is well suited to studying these plasmas because of its good spatial resolution in the core, and ability to measure the cross section effect, which becomes increasingly important for the high temperature core of ITB plasmas. Due to the combination of co- and counter-current neutral beams used to make the PAAR measurements, no prompt-torque effects$^{49}$ were observed in the toroidal rotation measurements. PAAR measurements have shown, for the first time on DIII-D, mean poloidal flow...
spin-up at the formation of an internal transport barrier.

Profiles of core poloidal rotation before and after the formation of an ITB are shown in Fig. 5.5. These results are for the same times in the same discharge as presented in Fig. 5.2. The measured poloidal rotation peaks in the same region where the transport barrier forms. The peak poloidal rotation is \(\approx 5\) km/s, which is much larger than the measurement uncertainty. The time history of the poloidal rotation is shown in Fig. 5.6. Though details of the neutral beam requirements set the minimum time between PAAR analysis, these results clearly show the change in the poloidal rotation being coincident with the dramatic changes in the core plasma parameters.

![Graph showing poloidal rotation before and after ITB formation](image)

**Figure 5.5**: Poloidal rotation before and after the formation of an ITB. Successive measurements show the poloidal rotation spinning-up in the region of the transport barrier.

This poloidal spin-up seen in 149472 was also observed for two similar ITB discharges: 149469 \((B_\varphi = -2.0\ T,\ I_p = 1.2\ MA,\ n_e = 2 \times 10^{19}\ m^{-3},\ P_{\text{NBI}} = 3.8\ MW,\ \tau_{\text{NBI}} = 2.4\ \text{Nm})\) and 149470 \((B_\varphi = -2.0\ T,\ I_p = 1.2\ MA,\ n_e = 2 \times 10^{19}\ m^{-3},\ P_{\text{NBI}} = 5.3\ MW,\ \tau_{\text{NBI}} = 3.3\ \text{Nm})\). Profiles of poloidal rotation before and after the ITB formation for these cases are shown in Fig. 5.7. Discharge
Figure 5.6: (a) Poloidal rotation at selected radii inside and outside an ITB ($\rho_{\text{foot}} \approx 0.4$), showing the spin-up and increase of shear during formation, and (b) ion temperature measurements made in the same time window, showing the barrier formation with a higher time resolution.

149469 exhibited a weaker transport barrier at a lower radius, and discharge 149470 exhibited ITB properties similar to 149472.
Figure 5.7: Poloidal rotation for two ITB discharges that are similar to 149472 (whose poloidal rotation is shown in Fig. 5.5). The amount of spin-up and location of maximum poloidal rotation track with the location of the transport barrier.

Figure 5.8: PAAR measurements of poloidal rotation before and after the formation of a large radius ITB plasma. The measurements show an increase in the magnitude and shear of the poloidal rotation after the transport barrier forms.
As described in Section 5.1, poloidal rotation results from the large radius ITB cases (Fig. 5.4) cannot determine history of poloidal rotation at the transport barrier. Nevertheless, interesting changes in the poloidal rotation are observed as the transport barrier forms. As shown in Fig. 5.8, when the ITB forms, poloidal rotation inside the transport barrier increases significantly in magnitude. The largest acceleration is in the electron diamagnetic direction at the mid radius region. In the deep core, ion diamagnetic acceleration is observed, as was also seen for the low-radius ITB plasma. These large radius ITB plasmas enable high resolution measurements of poloidal rotation in a large region of reversed magnetic shear.

5.2.1 \( E \times B \) shear during ITB formation

With the PAAR measurements of poloidal rotation in these ITB plasmas, the \( E \times B \) shear can be measured. It is expected that there will be an increase in the shearing rate that reduces turbulent transport and allows the transport barrier to form.\textsuperscript{13,71} For the low-radius ITB plasmas, there is a clear increase in

![Figure 5.9](image)

**Figure 5.9:** (a) Measurements of the \( E \times B \) shearing rate before and after the formation of an ITB, and (b) different components of the \( E_r \) measurement. Interestingly, the increase of \( E \times B \) shear is coincident with the ITB formation, and the poloidal rotation makes a significant contribution to that result.
the $\mathbf{E} \times \mathbf{B}$ shearing rate, as seen in Fig. 5.9(a). This increase is due to an increased contribution of the $\mathbf{V} \times \mathbf{B}$ terms. The different components of the $\mathbf{E} \times \mathbf{B}$ shearing rate calculation are shown in Fig. 5.9(b), where the poloidal rotation is shown to make a significant contribution to this increase in stability.

In the core of a transport barrier, the $\mathbf{E} \times \mathbf{B}$ shearing rate is also seen to increase after barrier formation, as seen in Fig. 5.10(a), which shows results from a large radius ITB plasma. This shearing rate increases particularly in the region just inside the transport barrier, where the $T_i$ and $V_\phi$ gradients are still relatively steep (see Fig. 5.4). The poloidal rotation is acting to decrease stability in this case though, as seen in Fig. 5.10(b).

**Figure 5.10**: (a) Measurements of the $\mathbf{E} \times \mathbf{B}$ shearing rate before and after the formation of a large radius ITB, and (b) the different components of $E_r$ that contribute to that measurement. Poloidal rotation is only significant near to the transport barrier, where its effect is to decrease the magnitude of the $\mathbf{E} \times \mathbf{B}$ shear.

### 5.3 Vertical CER measurements during ITB formation

As discussed in Section 3.3.4, the vertical CER system is not well suited for measurements of core poloidal rotation in an ITB plasma. In this section,
the observation of poloidal spin-up during ITB formation is investigated with a modified analysis of the vertical CER measurements. For simplicity, the analysis considers measurements from two view chords: one inside the transport barrier, and one outside the transport barrier. The difference in poloidal rotation from these two view chords is observed to diverge when the transport barrier forms and yields similar shear to the PAAR results, provided reasonable accounting is made for the lack of spatial localization in the position of the vertical view chords.

The standard vertical CER analysis requires measurements across the plasma profile, and uses the condition that $V_b(\rho = 0) = 0$ to help determine a value of the excited state lifetime, $\tau$, that yields the best fit to the vertical CER measurements. However, when the spatial resolution of the core vertical CER view chords is insufficient to obtain a good core measurement, this analysis breaks down. The potential for a lack of core spatial resolution for the vertical view chords was first discussed in Section 3.3.4. For the ITB discharge 149472, a relatively low major radius magnetic axis and high gradients in the core accentuate this lack of spatial resolution. The result is easily seen by comparing temperature measurements from vertical view chords with those from tangential view chords. The vertical view chord measurements are 2 keV (20%) lower than the tangential view chord measurements once the ITB has formed in discharge 149472.

The modified analysis uses co- and counter-current neutral beam measurements to determine the true toroidal rotation and the magnitude of the cross section effect (see Section 3.1.4). The form of the cross section effect on the measured velocities is taken from Ref. 72,

$$V_{CVI} = V_{CVII} + \alpha \dot{V}_b + \frac{\omega_c \tau}{1 + \omega_c^2 \tau^2} \left( \frac{V_{CVII} + \alpha \dot{V}_b}{B} \right),$$  

(5.1)

where $V_{CVI}$ is the velocity of the population of $C^{+5}$ ions, $(V_{CVII})$ is the velocity of the $C^{+6}$ ions, $\alpha \dot{V}_b$ is the velocity of the cross section effect (as in Ref. 55), and $\omega_c$ is the ion cyclotron frequency. Only the zeroth and first order terms in $\omega_c \tau$ have been retained. By working on the midplane, where vertical CER view chords intersect the neutral beam, and taking $\mathbf{B} \cdot \mathbf{R}$ and $V_{CVII} \cdot \mathbf{R}$ to be small, the value of $V_{CVII} \cdot \mathbf{z}$ can be found. This allows the poloidal rotation to be determined from
the vertical line of sight measurements ($V_{\text{LOS}}$),

$$\left(1 + \frac{s_R B \hat{\varphi}}{s_z B} G\right) V_{CV II} \cdot \hat{z} = \frac{V_{\text{LOS}}}{s_z} - \frac{s_R}{s_z} (V_{CV II} + \alpha \hat{V}_b) \cdot \hat{\varphi} - \frac{s_R}{s_z} \alpha \hat{V}_b \cdot \hat{R} - \frac{B \hat{\varphi}}{B} G \alpha \hat{V}_b \cdot \hat{R} - \frac{s_R B}{s_z B} G (V_{CV II} + \alpha \hat{V}_b) \cdot \hat{\varphi}, \quad (5.2)$$

where ($s_R$, $s_z$, $s_\varphi$) is the normalized viewing vector and $G \equiv \omega_c \tau/(1 + \omega_c^2 \tau^2)$. On the right-hand side of Eq. (5.2), the first term represents the raw measurement of the CER view chord, the second term represents the pickup of the toroidal rotation (which is finite for DIII-D vertical CER views), the third term represents the radial portion of the cross section effect, the fourth term represents the gyro-orbit cross section effect, and the fifth term represents the gyro-orbit correction to the radial portion of the cross section effect.

Using this equation, it is possible to determine the poloidal rotation for each vertical view chord based on values of $V_{CV II}$ and $\alpha$ from tangential view chords on co- and counter-current beams, and an ad hoc value for $\tau$ (keeping in mind that $\tau$ is typically less than a few nanoseconds). In addition, the amount of toroidal pickup and cross section effect can be adjusted to be more consistent with the lower measurements of temperature. Using the nominal positions of two vertical view chords, one just inside (at tokamak midplane, $\rho = 0.19$) and one just outside (at tokamak midplane, $\rho = 0.42$) the transport barrier, for discharge 149472 and a 1 ns excited state lifetime, the time history of the poloidal rotation on the outboard midplane, can be found. The result is shown in Fig. 5.11 where a toroidal spin-up approximately the same size as observed by the PAAR measurement (compare to Fig. 5.5) is seen.

Though a 1 ns excited state lifetime is on the relatively low end, the effect of increasing $\tau$ (and therefore the gyro-orbit cross section effect) can be compensated for with modest accounting for the spatial averaging. In Fig. 5.12, the excited state lifetime is doubled to 2 ns, and the spatial position of the view chord inside the transport barrier is adjusted outward, in $\rho$, by 0.05. The resulting spin-up of poloidal rotation is preserved. Note that this shift in $\rho$ is just enough to compensate for the under-measurement of the ion temperature. If $\tau$ is reduced for this case, the observed spin-up is increased. Therefore, it can be safely said that poloidal
Figure 5.11: Measurements of poloidal rotation at the outboard midplane when using the tangential measurements on co- and counter-current neutral beams and an excited state lifetime of 1 ns on either side of an ITB. After \( \approx 1200 \) ms, when the transport barrier forms, poloidal spin-up similar to PAAR measurements is observed.

Spin-up is consistent with the vertical CER measurements, given the value of \( \tau \) falls within the expected range.

The analysis presented here is imperfect, and a true correction of the vertical CER measurements would require performing an emissivity weighted integral across the lines of sight of the view chords and inverting the measurement to obtain local parameters, similar to the work done by Bell. Another possibility would be to take the PAAR measurements as the true poloidal rotation and use a synthetic diagnostic in a calculation that can calculate the excited state lifetime to simulate the spectra that should be observed by the vertical view chords. The method outlined here is meant only to roughly demonstrate that, even in a case where the vertical CER view chords are ill-suited to measure the poloidal rotation, modified analysis that uses reasonable assumptions suggests that the vertical CER measurements are consistent with the PAAR results.
Figure 5.12: Measurements of poloidal rotation at the outboard midplane when using the tangential measurements on co- and counter-current neutral beams and an excited state lifetime of 2 ns on either side of an ITB. The position of these view chords is increased in $\rho$ by 0.05 to compensate for the line integrated nature of the vertical CER measurement. Observed poloidal spin-up similar to Fig. 5.11.

5.4 Acknowledgements

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Chapter 6

Tests of poloidal rotation theories

In this chapter, poloidal rotation results are compared to theories of poloidal rotation that were described in Chapter 2. The most basic comparison is to neoclassical theory. Poloidal rotation driven by a turbulence induced Reynolds stress or friction with the fast-ions adds to the neoclassical results. As a result, all comparisons in this chapter begin with a neoclassical simulation. The simulations used are NCLASS\textsuperscript{32} (called with FORCEBAL\textsuperscript{33}) and NEO.\textsuperscript{34} The diversity of disagreement presented in this chapter demonstrates the complexity of poloidal rotation dynamics, and suggests that an accurate theory of poloidal rotation must involve more effects than are included in current models.

6.1 Basic tests of neoclassical predictions

The simplest test of poloidal rotation theories that can be performed is a comparison of measurements and results from neoclassical simulations. As discussed in Section 2.2.1, the more accurate neoclassical calculations are made by NEO while the NCLASS results are presented so that the error that results from uncertainty in the measurements that are inputs to the simulation can be seen. A comparison of the L- and H-mode poloidal rotation measurements in Fig. 3.15 to neoclassical simulations is shown in Fig. 6.1. The important issue to resolve is why disagreement is observed in some cases but not others.

One possible path of investigation is to use a dimensionless parameter
Figure 6.1: Measurements and neoclassical predictions of poloidal rotation for an (a) L-mode and (b) H-mode plasma. As seen for the H-mode case, not all measurements of poloidal rotation are neoclassical.

By creating matched plasmas that differ only in the value of one dimensionless parameter, it is possible to test the neoclassical scaling of poloidal rotation against the measurements for that parameter. A scan in $\nu^*$ was attempted for this dissertation because it is fundamental to the neoclassical calculation of poloidal rotation. Such a scan can be performed by creating two plasmas with different $B_\phi$, the same shape and density, and scaling $I_p$ linearly, and $T$ quadratically with $B_\phi$. Unfortunately, a subtlety of these types of scans prevented this analysis from yielding usable results. Though the value of $T$ was appropriately scaled, within error bars, $\nabla \ln T$ varied. As a result, it was not possible to separate out the $\nu^*$ and $L_T$ effects on poloidal rotation [see Eq. 2.16].

The study of indirect measurements of main-ion poloidal rotation has suggested that, when other dimensionless parameters are allowed to vary, decreasing $\nu^*$ results in larger poloidal rotation than is predicted by neoclassical calculations. However, no simple dependence like this has been observed with impurity poloidal rotation measurements. Even in cases of low collisionality, agreement has been observed. In Fig. 6.2, a comparison is shown for a H-mode discharge for which $\nu^* \approx 0.01$, and the agreement (within the error bars) shows that neoclassical pre-
Figure 6.2: The measurement and neoclassical prediction of poloidal rotation for a low collisionality H-mode. The impurity poloidal rotation is neoclassical even though $\nu^* \approx 0.01$.

Predictions can be accurate even when collisionality is low. Results such as these motivated more sophisticated lines of inquiry.

6.2 Tests of theory for ITB plasmas

In this section, the measurements of poloidal rotation during ITB plasmas from Chapter 5 are compared with neoclassical calculations. As in that chapter, the main dynamics of the poloidal rotation that is of interest is the spin-up that occurs as the transport barrier forms. Comparisons with neoclassical calculations before and after the ITB formation for the plasma shown in Fig. 5.5 are shown in Fig. 6.3.

Before the ITB forms [Fig. 6.3(a)], it is clear that the neoclassical calculation is inaccurate. The simulation predicts strong poloidal rotation due to the core temperature gradient, but this result is not seen in the plasma. Interestingly, after the ITB has formed, [Fig. 6.3(b)], the measurement and the neoclassical
Figure 6.3: Measurements and neoclassical predictions of poloidal rotation (a) before and (b) after the formation of the transport barrier in an ITB plasma. Disagreement is clear in (a), but agreement is seen in (b) in the transport barrier region (see Fig. 5.2).

Here, the temperature and main-ion density gradient have increased. As a result, the neoclassical predictions before and after the ITB forms show strong ion-diamagnetic poloidal rotation where the temperature gradient is largest. However, the measurements suggest that the temperature gradient drive of poloidal rotation is only strong after the ITB has formed. A similar phenomenology was observed for similar ITB discharges, as shown in Fig. 6.4.

Due to these observations, a hypothesis was formed that the presence of turbulence was the cause of disagreement with neoclassical predictions, and the suppression of turbulence in the transport barrier region allowed the rotation to reduce to just the neoclassical value (similar to the observations of heat transport in an ITB\textsuperscript{35–37}). In order to investigate this hypothesis, stronger internal transport barriers were sought, so that a correlation between turbulence suppression and neoclassical prediction accuracy could be obtained. These efforts resulted in the large radius ITB plasmas presented in Chapter 5.

Unfortunately, the extreme radii of the transport barrier for these plasmas
Figure 6.4: Measurements and neoclassical predictions of poloidal rotation before [(a) and (c)] and after [(b) and (d)] the formation of a transport barrier in two different ITB plasmas. Poloidal rotation spins up as the barrier forms and a narrow region of agreement with neoclassical predictions is observed where the transport barrier has formed.

precluded the measurement of poloidal rotation in the steep gradient region. This result highlights the reason ITB plasmas are ultimately undesirable for fusion reactor design: though they can achieve extreme plasma parameters, their stability can be marginal. In this case, though many attempts were made to change the
development of the ITB, its location was locked in. Even 50 ms changes in the timing of the NBI power resulted in a plasma disruption. This was probably a result of the discharge developing along a fine line in stability where the current diffusion was particularly important. Nevertheless, interesting results were obtained in the large region inside the transport barrier where the magnetic shear is reversed.

Comparisons of measurements and neoclassical predictions before and after the formation of this large radius ITB are shown in Fig. 6.5. At both times, significant disagreement is seen in the core of the plasma. This result is similar to the observations seen for the lower radius ITB plasmas, where disagreement inside the region where the transport barrier forms was seen before and after its formation as well. This disagreement in this case though is quite extreme. After the ITB forms, the measurement and neoclassical prediction are separated by as much as 8 km/s, and rotating in opposite directions. Qualitatively, the disagreement in this region of the discharge is unexpected because, after the transport barrier forms, the core gradient scale lengths are reduced (see Fig. 5.4), so the free energy

Figure 6.5: Measurements and neoclassical predictions of poloidal rotation (a) before and (b) after the formation of a large radius ITB plasma. Both cases show significant disagreement. This transport barrier is much stronger than the cases shown in Figs. 6.3 and 6.4, but poloidal rotation measurements can only be made inside the transport barrier.
source necessary for creating turbulence that could produce a Reynolds stress (see Section 2.3) should be reduced.

To investigate this more rigorously, linear stability calculations with GKS are used to assess changes in turbulence associated with the ITB formation. These results are shown in Fig. 6.6. For both times, the GKS calculations predict no instability is present for $\rho < 0.4$. Though the maximum growth rate increases after the ITB forms, the $E \times B$ shearing rate increases even more. These results show that predicted turbulence intensity is decreasing as the disagreement between measurements and neoclassical predictions is increasing. Therefore, this case shows qualitative disagreement with the hypothesis that agreement with neoclassical theory is due to the lack of turbulence induced Reynolds stress.

A complication for these discharges is the negative magnetic shear in the core and the potential this has to result in complex electron transport. Note that, though the GKS results do not predict instability, the propagation direction is the electron diamagnetic direction, the same as the direction of disagreement.
between the measurement and the neoclassical predictions.

6.3 Tests of theory with turbulence regime changes

Further investigations of turbulence driven poloidal rotation used a different approach. Rather than trying to suppress the turbulence in the core with strong $E \times B$ shear, the character of the turbulence was significantly altered with a change in heating source. Plasmas heated with NBI typically have very different gradients in density and temperature than plasmas heated with ECH. This is partially a result of the very different deposition profiles of the two heating methods, and partially a result of changes in transport that are clear once ECH is applied to the plasma (e.g. ECH pump-out). It is expected that a significant change in the plasma gradients can cause a significant change in the turbulence such as a turbulence regime dominated by ion turbulence (e.g. ion temperature gradient driven turbulence) to a regime dominated by electron turbulence (e.g. trapped electron mode or electron temperature gradient driven turbulence). If turbulence induced Reynolds stress is a significant drive of poloidal rotation, then such a change in the turbulence should be observable in the disagreement between measurements and neoclassical predictions of poloidal rotation.

Consider a quiescent H-mode (QH-mode) discharge that begins with NBI being the only auxiliary heating method and then adds ECH power approximately halfway through the discharge (see Fig. 6.7). Though the majority of the input power is still from NBI in the second half of the discharge, the added ECH causes a significant decrease in the core ion and electron temperature gradients. Analysis of this discharge is focused at 1890 ms (before the ECH is added) and 3010 ms (after the ECH is added). Turbulence linear growth rates are predicted by TGLF,\textsuperscript{77} and these results (see Fig. 6.8) show that the maximum ion growth rate is significantly reduced by the addition of ECH.

For this discharge, details of the neutral beam injection design precluded use of the PAAR measurement technique. As a result, vertical CER measurements
Figure 6.7: Traces showing the development of a QH-mode discharge with $B_\phi = -1.9$ T, $I_p = -1.1$ MA, and $\tau_{\text{NBI}} = 3.4$ Nm. Measurements in the core (blue, $\rho \approx 0.1$) and mid-radius (red, $\rho \approx 0.5$) show that the ion and electron temperature gradients change significantly once ECH is applied.

are the used to analyze the poloidal rotation instead. These measurements are compared to neoclassical predictions in Fig. 6.9. In the core of the plasma during the NBI phase, there is significant disagreement between the measurement and the neoclassical prediction, but once the ECH is added, the poloidal rotation is reduced significantly and shows agreement with the neoclassical prediction across nearly the entire profile.

Considering these measurements and the basic result from the TGLF calculations, that ITG turbulence is significantly reduced once ECH is added, there would appear to be good evidence for the generation of poloidal rotation via a Reynolds stress. However, more careful accounting of the Reynolds stress drive mechanism suggests that this effect cannot explain these results. Note that the disagreement seen in Fig. 6.9(a) is most significant at $\rho \approx 0.3$, but this is the region where TGLF calculations predict ITG growth rates to be small both before and after the ECH is added. Measurements of density fluctuations with the Doppler backscattering diagnostic$^7$ (DBS) also suggest that turbulence intensity is low in
this region at both times of the discharge. As shown in Fig. 6.10, the density fluctuation level for $0.3 < \rho < 0.45$ is constant and relatively small. Though the wavenumber probed here is more typical of the trapped electron mode (TEM), the observed reduction in fluctuations outside the core is correlated with the reduction of ITB turbulence predicted by TGLF (electron growth rates in this region at both times are near or below the $E \times B$ shearing rate). In this case then, though disagreement with neoclassical predictions can be correlated with a reduction in ITG turbulence, the location of the turbulence is not consistent with the turbulence being a direct drive of poloidal rotation via a Reynolds stress.

More quantitative comparisons of Reynolds stress drive and measured
Figure 6.9: Measurements and neoclassical predictions of poloidal rotation for (a) NBI and (b) NBI and ECH phases of a QH-mode discharge. Significant disagreement is seen in the core of the NBI phase, while in the NBI and ECH phase, agreement is nearly total.

Figure 6.10: Density fluctuation levels measured by DBS before and after the addition of ECH to a QH-mode discharge. Core fluctuations are constant and relatively weak while a significant reduction is seen for $0.5 < \rho < 0.7$. Probed wavenumbers are $k_\theta \approx 6 - 9.5 \text{ cm}^{-1}$. 
Figure 6.11: Reynolds stress per mass (a) before and (b) after the addition of ECH to a QH-mode discharge. This poloidal rotation drive term is decreased significantly after the addition of ECH, especially for $\rho < 0.5$.

Poloidal rotation are possible with simulations that can calculate the turbulence contribution to the stress tensor. The comparisons of these calculations with measurements is not routine, so a basic outline of the process is discussed here using results of the GTS simulation code. GTS is a gyrokinetic, delta-f, particle-in-cell code that can simulate electrostatic turbulence in full tokamak geometry. Here, the Reynolds stress drive is modeled as in Eq. (2.20). GTS can calculate the flux surface average of these fluctuating quantities for the main-ion and the results are shown in Fig. 6.11. These results corroborate the interpretation of the TGLF results by showing significant reduction in electrostatic-induced velocity fluctuations in the core after the ECH is added. However, the sign of the radial derivative of the poloidal rotation drive term changes at $\rho \approx 0.3$, indicating that the sign of the Reynolds stress driven poloidal rotation should change sign as well [see Eq. (2.20)]. This result is not seen in the measurements where poloidal rotation is consistently below the neoclassical prediction from $0.1 < \rho < 0.4$. Such comparisons are in the early stage because GTS does not include impurities in its simulation, but making rigorous comparisons such as these is important for future poloidal rotation research.

The same poloidal rotation dynamics are observed in another discharge that also has a NBI heated and NBI and ECH heated phase. This discharge is a
basic H-mode plasma (i.e. not QH-mode) with much lower auxiliary power input. Additional differences between this discharge and the one just presented are: total auxiliary power is kept constant when ECH is turned on, NBI torque is reduced significantly when ECH is turned on, the value of $T_e/T_i$ is less than one before ECH is turned on and greater than one after, and PAAR measurements are available at an intermediate time between the NBI heated and NBI and ECH heated phases. The development of this discharge is shown in Fig. 6.12. Though the NBI is a series of blips after the ECH is added, the fast-ion slowing down time is on the order of 100 ms and will serve to smooth the input power absorbed by the thermalized particles in the plasma.

**Figure 6.12**: Traces showing the development of a H-mode discharge with $B_\phi = -1.9$ T, $I_p = 1.1$ MA, $P_{\text{NBI+ECH}} = 4.2$ MW, and $\tau_{\text{NBI}} = 4.1 - 0.5$ Nm. Measurements in the core ($\rho \approx 0.1$) and mid-radius ($\rho \approx 0.5$) show that the ion and electron temperature profiles are changed significantly once the ECH is applied.

To determine the effect of the change in heating power, GKS linear stability calculations were performed. These show that the turbulence intensity is peaked at $\rho \approx 0.6$ for both phases of this discharge, and that the strongest mode in the primarily ECH heated phase is TEM while it is ITG during the primarily NBI
heated phase. The ratio of $\gamma_{\text{max}}$ to $\gamma_{E \times B}$ suggests that turbulence will be more significant for the ECH heated phase.

![Figure 6.13](image)

**Figure 6.13:** (a) Growth rates of the most unstable mode and (b) corresponding mode frequencies calculated by GKS for a H-mode discharge during a primarily NBI heated and primarily ECH heated phase. At both times, no instability is predicted for $\rho < 0.2$. The measurement of the $E \times B$ shearing rate, $\gamma_{E \times B}$, suggests the turbulence intensity is more significant during the primarily ECH heated phase.

The poloidal rotation measurements for the two phases of this discharge are shown in Fig. 6.14. As was seen for the QH-mode discharge just presented, a significant disagreement between the measurement and neoclassical prediction during the NBI phase was observed for $\rho < 0.4$, the region where turbulence intensity is predicted to be low. When the ECH is added and NBI power and torque reduced, the measurement and neoclassical prediction agree very well. These results are not consistent with a basic turbulence induced Reynolds stress driven poloidal rotation because the disagreement is seen where turbulence is not present to drive the poloidal rotation. As for the QH-mode discharge just presented, these measurements were made with the vertical CER system, but at a time between the transition from primarily NBI to primarily ECH heating, PAAR measurements are available. The comparison between these two measurements of poloidal rotation was shown in Fig. 3.23 where it can be seen that these two measurement methods
are in agreement.

![Graphs showing poloidal rotation measurements and neoclassical predictions for (a) NBI and (b) NBI and ECH phases of a H-mode discharge. Significant disagreement is seen in the core of the NBI phase, while in the NBI and ECH phase, agreement is good.](image)

**Figure 6.14**: Measurements and neoclassical predictions of poloidal rotation for (a) NBI and (b) NBI and ECH phases of a H-mode discharge. Significant disagreement is seen in the core of the NBI phase, while in the NBI and ECH phase, agreement is good.

### 6.4 Tests of theory for fast-ion friction

For both the QH- and H-mode discharges with a NBI and ECH phase that were discussed in Sec. 6.3, poloidal rotation measurements were observed to agree with neoclassical predictions after the addition of ECH. One hypothesis for this behavior is that friction with the fast-ions is driving the disagreement with neoclassical theory. For the H-mode discharge (Fig. 6.12), where the addition of ECH power is accompanied by a reduction in NBI torque, it is plausible that this results in significantly reduced fast-ion friction and then causes poloidal rotation to be neoclassical. For the QH-mode discharge (Fig. 6.7), the addition of ECH is not accompanied by a change in the NBI power or torque, but it is possible that the associated change in ion and electron temperature profiles altered the dynamics of the slowing down of the fast-ions and, therefore, had an effect on the fast-ion driven poloidal rotation.
However, it is unlikely that fast-ion friction is the cause of both observations of disagreement with neoclassical theory because, while the direction of the toroidal field is the same in these two discharges, the direction of the plasma current is different. As a result, due to the projection of the neutral beam injection direction onto the magnetic field line, it is expected that fast-ion friction would drive poloidal rotation in the electron-diamagnetic direction for the QH-mode discharge, and the ion-diamagnetic direction for H-mode discharge. The measurement of neoclassical poloidal rotation is more electron-diamagnetic for both discharges. Based on Eq. 2.23, it is expected that the QH-mode discharge would exhibit stronger fast-ion friction because of its higher proportional fast-ion density (compare \( n_e \) and \( P_{\text{NBI}} \) for the two discharges in Fig. 6.7 and 6.12). Also note that the more electron diamagnetic poloidal rotation (compared to the neoclassical prediction) observed in a large radius ITB plasma (Fig. 6.5) is consistent with the basic picture of the fast-ion friction.

More quantitative investigations of poloidal rotation driven by fast-ion friction are possible on DIII-D due to the ability of co- and counter-current NBI to decorrelate \( P_{\text{NBI}} \) and \( \tau_{\text{NBI}} \). Discharges with high NBI torque have a population of fast-ions traveling in one direction while discharges with small NBI torque have a population of fast-ions with no net rotation. Therefore, all else being equal, a discharge with high NBI torque should exhibit more poloidal rotation driven by fast-ion friction than a discharge with low NBI torque.

To study the effect of fast-ion friction on poloidal rotation, two similar QH-mode discharges \( (B_\phi = 2.1 \, \text{T}, I_p = 1.1 \, \text{MA}, \, n_e = 2.5 \times 10^{19} \, \text{m}^{-3}) \) with different amounts of NBI torque were analyzed. The high torque discharge has 1.8 Nm of co-current torque provided by the neutral beams, and the low-torque discharge has \(-0.2 \, \text{Nm}\) of torque due to the addition of significant torque from the counter-current neutral beams. Based on results of Ref. 46, fast-ion friction should be significant in the high-torque discharge and insignificant in the low-torque discharge. The measurements and neoclassical predictions for these two cases are shown in Fig. 6.15. Unexpectedly, the agreement with the neoclassical predictions is very good in the high torque discharge, and significant disagreement is seen in
the core of the low torque discharge. These results show that fast-ion friction, on its own, cannot explain the observed discrepancies between measurements and neoclassical theory of poloidal rotation.

Figure 6.15: Measurements and neoclassical predictions of poloidal rotation for two similar QH-mode discharges with (a) a significant NBI torque of 1.8 Nm and (b) an insignificant NBI torque of 0.2 Nm. Agreement is worse in the low torque case.

6.5 Simultaneous modeling of Reynolds stress and fast-ion friction

It is possible that the measurements shown so far can be explained by a combination of turbulence induced Reynolds stress and fast-ion friction. If both effects are significant, interplay between them could yield the results shown, which are not consistent with either one of these models on their own. In order to estimate how these two effects could combine, the equation for parallel momentum conservation is used:

\[
\frac{\partial}{\partial t} \langle m n \mathbf{B} \cdot \mathbf{V} \rangle + \mu \langle B^2 \rangle \frac{V_\theta - V_{\theta}^{\text{neo}}}{B_\theta} = \langle \mathbf{B} \cdot \mathbf{F}_{\text{fast}} \rangle - \langle \mathbf{B} \cdot \nabla \Pi_{\text{turb}} \rangle. \tag{6.1}
\]
Here, $V_{\theta}^{\text{neo}}$ is the neoclassical poloidal rotation, $F_{\text{fast}}$ is the friction force of the fast-ions and $\Pi_{\text{turb}}$ is the turbulence induced stress tensor.

The value of $V_{\theta}$ (which is present both explicitly and implicitly in the first term of Eq. 6.1) can be solved for with a combination of simulations. TGLF is used to determine the turbulence induced stress tensor for carbon, NUBEAM$^{61}$ and ONETWO$^{80}$ are used to estimate the momentum that is deposited in the parallel direction by the fast-ions, NEO$^{34}$ is used to determine the neoclassical poloidal rotation, and XPTOR$^{81}$ is used to solve for the steady state poloidal rotation. Within XPTOR, poloidal rotation and the radial electric field are allowed to evolve while all other experimental parameters that are input are held constant.

The fast-ion contribution is estimated by recasting the torque profile calculated in ONETWO as a total toroidal force on all the ions. This rate of momentum input is projected onto the parallel direction and distributed to the main-ion and impurity according to the formulas in Ref. 82. The poloidal damping term, $\mu$, is taken to be the result of the neoclassical viscous stress and is determined by NCLASS.$^{32}$

This new, “comprehensive” analysis technique was applied to the H-mode discharge analyzed in Sec. 6.3 (Fig. 6.12), and the results are shown in Fig. 6.16. With respect to the purely neoclassical predictions seen in Fig. 6.14, the comprehensive analysis comes much closer to the measurement during the NBI phase of this discharge, but the agreement is reduced in the NBI and ECH phase. Note that the prediction is well outside the measurement error for both times at many locations.

The relative contributions of the Reynolds stress and the fast-ion friction to the poloidal rotation in this comprehensive analysis are shown in Fig. 6.17. As expected, the fast-ion friction is largest during the NBI phase and negligible during the NBI and ECH phase when $\tau_{\text{NBI}}$ has been reduced by an order of magnitude. During both phases, the slowing down of the fast-ions is not taken into account, but this choice is justified by the fast that toroidal rotation has equilibrated at the two analysis times. The turbulence induced Reynolds stress calculation is similar during both phases. Though this model is still rather simplistic, its results
Figure 6.16: Measurements and a comprehensive simulation of poloidal rotation for the (a) NBI and (b) NBI and ECH phases of a H-mode plasma. The simulation includes turbulence induced Reynolds stress and fast-ion friction. Relative to the strictly neoclassical predictions in Fig. 6.14, the calculation is much closer to the measurement during the NBI phase.

are important because they show that both the fast-ion friction and turbulence induced Reynolds stress can drive significant amounts of poloidal rotation.

Figure 6.17: Volume integrated turbulence induced Reynolds stress and fast-ion friction contributions to poloidal momentum [see Eq. (6.1)] for the (a) NBI and (b) NBI and ECH phases of a H-mode discharge. Both effects can be comparable in size, as seen in (a) when NBI torque is significant.
6.6 Conclusions

In this chapter, the comparison of poloidal rotation measurements and theory has been significantly expanded beyond the basic neoclassical paradigm. As was outlined in Secs. 2.3 and 2.4, plausible explanations for disagreement with neoclassical predictions are turbulence induced Reynolds stress and friction with fast-ions. Some indication of turbulence reduction being correlated with the accuracy of neoclassical theory was found for ITB plasmas. However, large changes in turbulence intensity and propagation direction did not manifest in the accuracy of neoclassical theory in the region of the plasma that would be expected from a turbulence induced Reynolds stress model. An example of fast-ion friction being a poor predictor of neoclassical accuracy was also presented. Simple modeling suggests that both these effects can be significant drivers of poloidal rotation.

The examples presented here were specifically chosen to highlight the unexpected trends in disagreement between measurements and theory of poloidal rotation. All together, they demonstrate that neoclassical theory does not contain an accurate description of poloidal rotation and that a turbulence induced Reynolds stress and fast-ion friction can potentially drive significant poloidal rotation. More advanced poloidal rotation models, including these two effects, at a minimum, are necessary to make accurate predictions. The implications of these findings are discussed more in Chapter 7.

6.7 Acknowledgements

Chapter 6 contains material from Physics of Plasmas vol. 21. Chrystal, Colin; Burrell, Keith H.; Grierson, Brian A.; Staebler, Gary M.; Solomon, Wayne M.; Wang, Weixing X.; Rhodes, Terry L.; Schmitz, Lothar; Kinsey, Jon E.; Lao, Lang. L; deGrassie, John S.; Mordijck, Saskia; Meneghini, Orso, American Institute of Physics, 2014. The dissertation author was the primary investigator of this paper.
Chapter 7

Conclusions

The work in this dissertation can be separated into two complementary components. The diagnostic portion is covered in Chapters 3 and 4, and the experimental results portion is covered in Chapters 5 and 6. The experimental investigations were enabled by the diagnostic work that developed a new method for measuring poloidal rotation with CER. Ultimately, the work of this dissertation shows the complexity of poloidal rotation in tokamaks is beyond current prediction capabilities. A complete theory of poloidal rotation will need to include neoclassical effects, turbulence induced Reynolds stress, and fast-ion friction.

The new measurement method developed for this dissertation, PAAR, determines poloidal rotation by measuring the poloidal asymmetry in toroidal angular rotation. The asymmetry is measured with CER view chords on the high- and low-field side of the tokamak midplane. The high-field side view chords were added specifically to enable the work of this dissertation. Significant hardware modifications were needed to make this upgrade to the CER system. As part of that upgrade, a new method for correcting the astigmatism of a Czerny-Turner spectrometer was developed. Since the PAAR method relies entirely on tangential CER view chords, poloidal rotation measurements made by this method do not need to account for the finite excited state lifetime of the ions that have undergone a charge exchange reaction, and no atomic physics calculations are used to determine the effect of the energy dependent charge exchange cross section. In addition to this inherent advantage, the PAAR method has higher precision and better
spatial resolution than the measurement of poloidal rotation with vertical CER views. Consistency between these two measurement techniques is seen, but the PAAR method enables more detailed studies of poloidal rotation. In addition, a new method for creating a spatial calibration for the CER diagnostic has been developed to ensure the accuracy of the PAAR method.

The measurements made with the PAAR technique have been used to obtain new and noteworthy results. Measurements during the formation of an ITB showed poloidal spin-up associated with the transport barrier forming. The resulting high poloidal rotation makes a significant contribution to the increased $E \times B$ shearing rate associated with the transport barrier. These measurements also showed that neoclassical predictions of poloidal rotation are inaccurate in the core of ITB plasmas. The lack of predicted turbulence in these regions suggested that a turbulence induced Reynolds stress could not explain these results.

Further tests of turbulence induced Reynolds stress showed that when ion and electron temperature profiles were significantly modified by the addition of ECH in NBI heated plasmas, the associated changes in predicted turbulence could not explain the changes in poloidal rotation measurements. Although turbulence intensity and disagreement with neoclassical predictions were correlated, the location of the change in turbulence was not the same as the location where the accuracy of the neoclassical predictions changed.

It was observed that a basic model of fast-ion friction could correctly predict the sign of the discrepancies observed for most of the cases analyzed where neoclassical predictions were inaccurate. However, more stringent tests of the fast-ion friction showed that similar plasmas with different fast-ion rotation speeds had nearly equal poloidal rotations, and the neoclassical prediction of poloidal rotation was most accurate for the plasma with large fast-ion rotation. Since this plasma is expected to have the most fast-ion friction, this analysis showed that fast-ion friction cannot be the sole cause of the inaccuracy of neoclassical predictions.

A basic analysis that combined the turbulence induced Reynolds stress and fast-ion friction effects suggested that both of these effects can drive significant poloidal rotation. An accurate theory of poloidal rotation will need to include
realistic models for at least these two effects. The fast-ion drive mechanism is of particular importance because of the uncertainty in the fast-ion dynamics of future burning plasmas. Significant fast-ion friction also enables for complex interplay between fast-ion transport mechanisms, e.g. Alfvén eigenmodes, and $E \times B$ shear through the poloidal rotation.

The results of this dissertation have shown that, when using the most accurate poloidal rotation measurements that do not rely on atomic physics calculations, poloidal rotation in the core of tokamaks is clearly not accurately predicted by neoclassical theory. The observed discrepancies are not wholly consistent with the turbulence induced Reynolds stress or the fast-ion friction model. Results indicate that both of these effects could drive significant poloidal rotation, however, the models that are currently available do not yield quantitative agreement with experiment. Creating accurate poloidal rotation predictions is therefore more complicated than the current paradigm. New poloidal rotation theory and models are needed to understand this important aspect of tokamak rotation.

7.1 Future Work

The development of more advanced poloidal rotation models is a challenging but important avenue of future work. Including accurate turbulence induced Reynolds stress and fast-ion friction effects is difficult because the modeling of turbulence and fast-ion dynamics are themselves active areas of current research. Nevertheless, such comprehensive models could easily determine if basic trends in the inaccuracy of neoclassical theory align with expectations. These models could also be used to predict poloidal rotation for future tokamaks and thereby calculate whether or not poloidal rotation will be an important component of $E \times B$ shear. Though the prediction of toroidal rotation is another difficult problem that bears on this question, if it could be determined that poloidal rotation can create $E \times B$ shear on the same level as the pressure gradient term, it could be concluded that poloidal rotation dynamics cannot be ignored for future tokamaks. The pressure
gradient term is relatively easy to determine if it is assumed that tokamaks meet their performance goals.

The disagreement observed within the core of ITB plasmas was the most significant disagreement observed within this dissertation. A defining characteristic of the core of ITB plasmas is reversed magnetic shear, so these results in this work suggest that further investigation of poloidal rotation during reversed magnetic shear may yield interesting observations. Though future tokamaks are unlikely to operate with plasmas that have an ITB, these results could have implications for the stability of the plasma during the initial current ramp. At this time, toroidal rotation will likely be low, so the poloidal rotation could be very important for plasma stability.

Another observation was that the addition of ECH caused the magnitude of poloidal rotation to decrease significantly, at which point the neoclassical prediction was accurate. Though the addition of ECH has many effects on transport in all channels (that are poorly understood), if this were a property of electron heating in general, it is expected that a burning plasma would have neoclassical poloidal rotation. In this scenario, the prediction of poloidal rotation for future tokamaks is greatly simplified. Another possibility is that the electron heating serves to slow poloidal rotation. In this case, the poloidal rotation is unlikely to make a significant contribution to $E \times B$ shear in future tokamaks. For this reason, it is important to study the origin of this observation in ECH heated discharges.

The poloidal viscous damping within a plasma is fundamental to poloidal rotation dynamics, but no direct measurements have ever been made of this parameter. If the damping rate were comparable to the CER measurement rate, it would be possible to see transient poloidal rotation dissipate, and use the dissipation rate to determine the viscous damping. An experiment designed by Burrell to attempt this measurement uses a high temperature, low density QH-mode plasma to decrease the collisionality and prompt neutral beam torque to provide transient poloidal rotation. Measurements of the viscous damping with this method have not yet been made, but are a promising avenue of future work that can help determine the origin of the inaccuracy of neoclassical poloidal rotation.
The inaccuracy of neoclassical theory could have implications for neoclassical calculations of the bootstrap current. If some of the inaccuracy of the neoclassical theory of poloidal rotation can be attributed to the calculation of parallel force balance, the bootstrap current calculations could be affected. Experiments designed to test poloidal rotation and bootstrap current simultaneously could shed light on this potential linkage. The accuracy of the bootstrap current calculation is of the utmost importance for the design of future reactors and the development of steady-state tokamak scenarios.

The new CER views that were added to enable PAAR analysis can also measure impurity density asymmetry within a flux surface. These measurements are already being made because they are necessary for making PAAR measurements in highly rotating plasmas. However, using these measurements to investigate the theory of parallel impurity transport has not yet been done. These results would complement work done by Reinke\textsuperscript{9,64} by using CER and investigating low-Z impurities. DIII-D is an ideal environment for testing these theories because the centrifugal force and the electrostatic effect of trapped fast ions (both effects are observed to be significant on DIII-D) can be varied with the use of co-current and counter-current NBI.
**Glossary**

**ASDEX Upgrade** A medium sized tokamak in Garching, Germany, that is run by the Max-Planck-Institut für Plasmaphysik.

**Banana** See Trapped.

**Blip** A short neutral beam injection pulse, typically 10 ms in duration.

**Bootstrap current** Toroidal plasma current that is self-generated in the plasma due to a combination of the pressure gradient and the trapped particle orbits.

**CCD** Charge coupled device, solid-state based image detectors that can have high sensitivity and high signal-to-noise.

**CER** Charge exchange recombination spectroscopy, a diagnostic technique for determining ion rotation, temperature, and density. Many other acronyms exist, e.g. CXRS, CHERS, etc.

**CMM** Coordinate measuring machine, a mobile device that can measure distances on complex shapes on the multiple-meter scale to sub-millimeter accuracy.

**Co-current neutral beam** A neutral beam whose injection direction is aligned with the usual direction of $I_p$ ($> 0$).

**Collisionality** The collision frequency normalized to the frequency of trapped particle orbits, denoted by $\nu^*$, see Eq. (2.12).

**Counter-current neutral beam** A neutral beam whose injection direction is anti-aligned with the usual direction of $I_p$ ($< 0$).
Cross section effect A broad term that refers to any effect the energy dependence of the charge exchange cross section has on CER measurements.

Czerny-Turner spectrometer A spectrometer that uses two mirrors and a diffraction grating.\textsuperscript{47}

DBS Doppler backscattering, a diagnostic that uses backscattered microwaves to measure turbulence flows and fluctuations.

DIII-D A medium sized tokamak San Diego, CA, run by General Atomics where the research for this dissertation was performed.

Discharge A pulse of the tokamak that generates a plasma.

Disruption An event that very quickly terminates the plasma current and ends the plasma discharge.

ECH Electron cyclotron heating, a method for adding heat to the plasma by launching microwaves that resonate with the electrons in the plasma.

EFIT A computer program that determines the magnetic equilibrium reconstruction.\textsuperscript{2}

Electron diamagnetic The direction electrons drift due to the pressure gradient. For the standard $B_\varphi$ direction ($< 0$), the electron diamagnetic drift direction is up on the outboard midplane.

ELM Edge localized mode, a fast relaxation of the pressure gradient in the edge of H-mode plasmas that transiently releases particles and heat.

ETG Electron temperature gradient, a type of plasma turbulence driven unstable by the electron temperature gradient.

Fast-ions A class of ions that travel much faster than $V_{th}$ and have a non-Maxwellian distribution function. In DIII-D, fast-ions are the result of NBI.
Flux surface A toroidal surface in the plasma that is traced by magnetic field lines and has a constant value of $\psi$.

FORCEBAL A front-end code for NCLASS.\textsuperscript{33}

Full, half, third energy component The different energy neutrals injected by the neutral beams due to the presence of $D^+_2$ and $D^+_3$ in addition to $D^+$ in the beam’s ion source.

GKS A gyrokinetic instability simulation code.\textsuperscript{75}

GTS A gyrokinetic particle-in-cell tokamak plasma simulation code.\textsuperscript{79}

Gyro-orbit cross section effect A combination of the cross section effect and an ion’s gyro-motion that can greatly complicate measurements of vertical velocity in tokamaks.

Gyrokinetic A regime of plasma study that averages over the gyro motion, frequently used to study turbulence.

H-mode High confinement mode, an operating regime characterized by a large gradient in pressure at the edge of the plasma that is due to a transport barrier.

High-field side The side of the plasma with major radius lower than the major radius of the magnetic axis (where the toroidal magnetic field is stronger).

Impurity Ions in the plasma that are not the main-ion.

Integrating sphere A very high reflectivity enclosure that emits Lambertian light from its aperture.

Ion diamagnetic The direction ions drift due to the pressure gradient. For the standard $B_\phi$ direction ($<0$), the ion diamagnetic drift direction is down on the outboard midplane.
**ITB** Internal transport barrier, a region of very large gradients that is in the core of the plasma.

**ITER** A tokamak being built in France that will be the largest tokamak to date and whose basic goal is to create a burning plasma.

**ITG** Ion temperature gradient, a type of plasma turbulence driven unstable by the ion temperature gradient.

**JET** A large tokamak in Culham, United Kingdom, run by Culham Centre for Fusion Energy and used by all European fusion laboratories.

**L-mode** Low confinement mode, in contrast to H-mode, this operating regime does not have a transport barrier at the edge of the plasma.

**Low-field side** The side of the plasma with major radius higher than the major radius of the magnetic axis (where the toroidal magnetic field is weaker).

**Magnetic pumping** A dissipative mechanism due to a combination of collisions and a time varying magnetic field.

**Main-ion** The ion species that makes the dominant contribution to the positive charge in the plasma. Deuterium is the main-ion for most DIII-D discharges.

**MAST** A medium sized spherical tokamak in Culham, United Kingdom, run by the Culham Centre for Fusion Energy.

**MSE** Motional Stark effect, a neutral beam based diagnostic that measures the pitch of the magnetic field line.

**NBI** Neutral beam injection, a method for adding heat and momentum to the plasma by injecting accelerated ions that have been neutralized.

**NCLASS** A neoclassical plasma simulation code that uses Hirshman and Sigmar’s moment approach.\(^{30}\)
NEO A neoclassical plasma simulation code that solves the drift kinetic equation in orders of $\rho^*$.³⁴

Neoclassical A theory of tokamak transport that accounts for spatial variation in the magnetic field.

Neutral beam A device that is the source of NBI.

Neutron browning The transmission reducing effect of neutron exposure on fiber optics.

NSTX A medium sized spherical tokamak in New Jersey run by Princeton Plasma Physics Laboratory.

NUBEAM A plasma simulation code that predicts neutral beam fast-ion properties.⁶¹

PAAR Poloidal asymmetry in angular rotation, a method for determining poloidal rotation with tangential velocity measurements at multiple points on a flux surface.

Parallel The direction of the magnetic field.

Perpendicular The direction that is perpendicular to the magnetic field.

Poloidal The direction that travels the short way around the torus, denoted by $\theta$.

QH-mode Quiescent high confinement mode, an operating regime similar to H-mode that does not exhibit ELMs.

Reversed magnetic shear $B_\theta$ that increases with radius due to a hollow current profile.

Reynolds stress The component of the total stress tensor due to fluctuating velocities.
**Safety factor**  The ratio of the number of times the magnetic field traverses the tokamak poloidally to the number of times it traverses the tokamak toroidally.

**Shot**  A pulse of the tokamak systems that may or may not generate a plasma.

**TCV**  A medium sized tokamak in Lausanne, Switzerland, run by École polytechnique fédérale de Lausanne.

**TEM**  Trapped electron mode, a type of plasma turbulence driven unstable by the electron temperature and density gradients.

**TFTR**  A large tokamak that was in New Jersey and run by Princeton Plasma Physics Laboratory.

**TGLF**  A gyrokinetic instability simulation code.\(^{77}\)

**Toroidal**  The direction that travels the long way around the torus, denoted by \(\varphi\).

**Transport barrier**  A region of the plasma where transport is reduced due to \(E \times B\) shear.

**Trapped**  A property of particles or particle orbits that prevents them from traveling to the high field side of the plasma due to the magnetic mirror force.

**Vent**  A time period when the tokamak is opened to the outside environment and maintenance is typically performed.

**View chord**  The line of sight from which light is collected for an optical diagnostic, e.g. CER.

**XPTOR**  A transport equation solver.\(^{81}\)
Symbols

\( a \) The plasma minor radius

\( \alpha \) The magnitude of the cross section effect

\( \mathbf{B} \) The magnetic field

\( B_0 \) Magnetic field at the magnetic axis

\( \langle A \rangle \) Flux surface average except for in Sec. 2.1 where \( \langle A \rangle \) is the normalized average over the distribution function

\( C \) Collision operator

\( D \) Dispersion

\( \mathbf{E} \) The electric field

\( E_{\text{beam}} \) Accelerator voltage of a neutral beam

\( \epsilon \) Inverse aspect ratio, \( a/R \)

\( E_r \) Radial electric field

\( \omega_{E \times B} \) \( \mathbf{E} \times \mathbf{B} \) shearing rate, see Eq. (1.4)

\( f_a \) Distribution function for species \( a \)

\( \gamma \) The angle between \( \hat{V}_b \) and \( \hat{\phi} \)

\( \gamma_{\text{max}} \) The growth rate of the most unstable mode
\( I_p \) The total toroidal current in the plasma

\( I(\psi) \ RB_\varphi \)

\( J \) Current density

\( K \) Poloidal velocity flux function used in PAAR analysis (see Sec. 3.3.3)

\( k \) Poloidal velocity flux function, see Eq. (2.11)

\( \lambda_0 \) Rest wavelength of an atomic emission line

\( m \) Mass

\( \mu \) Neoclassical viscous damping coefficient

\( n \) Density

\( n_b \) Density of neutral beam neutral particles

\( n_p \) Density measured by CER at a position on the outboard midplane

\( \nu^* \) Collisionality

\( \omega \) Rigid rotor angular rotation, see Eq. (2.11)

\( \omega_c \) Gyrofrequency

\( \omega_{\text{max}} \) The frequency of the most unstable mode

\( \Omega^* \) Rigid rotor angular rotation used in PAAR analysis (see Sec. 3.3.3)

\( P \) Pressure, i.e. \( nT \)

\( P_{\text{aux}} \) Auxiliary power input to the plasma where aux can be NBI, ECH or both

\( \Phi \) The electrostatic potential in the plasma

\( \varphi \) The toroidal angle coordinate, positive in the counter clockwise direction when the tokamak is viewed from above
Total stress tensor

Viscosity tensor

Poloidal flux function, $\nabla \psi \times \nabla \varphi = B_{\theta}$

Electric charge

Major radius

Radial coordinate from magnetic axis

Major radius of magnetic axis

A radial coordinate representing the toroidal flux contained within a flux surface normalized to the total toroidal flux contained within the plasma boundary

Gyroradius

Gyro radius normalized to the minor radius, $a$

Major radius of $n_p$

Normalized vector representing the direction photons travel from the neutral beam to the lens for a CER view chord

Temperature

Excited state lifetime of a charge exchange electron

The torque applied by NBI on the plasma

The poloidal angle coordinate, positive is down on the outboard midplane of the tokamak

Mean velocity

Velocity

Angle between $\hat{s}$ and $\hat{V}_b$
$\dot{V}_b$ Neutral beam injection direction

$V_{\text{LOS}}$ Line of sight velocity

$V_{th}$ Thermal speed, $\sqrt{2T/m}$

$Z$ Atomic number

$z$ Displacement from tokamak midplane
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