Global dynamics of magnetic reconnection in VINETA II

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Abstract

Magnetic reconnection is a fundamental plasma process where a change in field line connectivity occurs in a current sheet at the boundary between regions of opposing magnetic fields. In this process, energy stored in the magnetic field is converted into kinetic and thermal energy, which provides a source of plasma heating and energetic particles. Magnetic reconnection plays a key role in many space and laboratory plasma phenomena, e.g. solar flares, Earth’s magnetopause dynamics and instabilities in tokamaks. A new linear device (VINETA II) has been designed for the study of the fundamental physical processes involved in magnetic reconnection. The plasma parameters are such that magnetic reconnection occurs in a collision-dominated regime. A plasma gun creates a localized current sheet, and magnetic reconnection is driven by modulating the plasma current and the magnetic field structure. The plasma current is shown to flow in response to a combination of an externally induced electric field and electrostatic fields in the plasma, and is highly affected by axial sheath boundary conditions. Further, the current is changed by an additional axial magnetic field (guide field), and the current sheet geometry was demonstrated to be set by a combination of magnetic mapping and cross-field plasma diffusion. With increasing distance from the plasma gun, magnetic mapping results in an increase of the current sheet length and a decrease of the width. The control parameter is the ratio of the guide field to the reconnection magnetic field strength $B_g/B_{xy}$. Cross-field plasma diffusion leads to a radial expansion of the current sheet at low guide fields. Plasma currents are also observed in the azimuthal plane and were found to originate from a combination of the field-aligned current component and the diamagnetic current generated by steep in-plane pressure gradients in combination with the guide field. The reconnection rate, defined via the inductive electric field, is shown to be directly linked to the time-derivative of the plasma current. The reconnection rate decreases with increasing $B_g/B_{xy}$, which is attributed to the plasma current dependency on axial boundary conditions and the plasma gun discharge. The above outlined results offer insights into the complex interaction between magnetic fields, electric fields, and the localized current flows during reconnection.
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Chapter 1

Introduction

Plasma dynamics are dominated by complex interactions between magnetic and electric fields and charged particles. Understanding the fundamental processes involved is key in explaining a multitude of observed plasma phenomena. A full description of the dynamics of a plasma is often impossible due to the complexity of the system. Instead, simplified models are utilized which can be extended or adapted as needed depending on a wide range of factors such as the plasma parameters and the relevant length- and time-scales. This is also the case for magnetic reconnection, which is a fundamental plasma process where a change in the connectivity of magnetic field lines leads to a topological rearrangement of a magnetic field and a new equilibrium configuration with lower magnetic energy [1–3]. One of the key aspects of magnetic reconnection is that through a coupling of global and local dynamics, energy stored in the magnetic field can be converted into kinetic energy and thereby provide a source of plasma heating and acceleration of fast particles [4–7].

Despite intense research on magnetic reconnection over the last decades, and a vast quantity of information obtained through space observations [8, 9], numerical simulations [10, 11], and experiments [12–19], the underlying plasma processes governing magnetic reconnection are still not well understood. Magnetic reconnection is influenced by a wide variety of factors and circumstances: For resistivity dominated plasmas, simple resistive magnetohydrodynamic (MHD) models can provide an adequate descrip-
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They are, however, insufficient to explain reconnection in collisionless plasmas and the observed fast timescale of the reconnection dynamics [1, 2]. Instead, more complex approaches are required, incorporating a multi-fluid description [23, 24], kinetic effects [25, 26], and instabilities [27, 28]. Reconnection is also often described using a simplified two-dimensional geometry, where it occurs in the so-called current sheet formed at the boundary between opposing magnetic fields. However, more recent findings suggest that a three-dimensional description, incorporating factors such as an additional component perpendicular to the reconnection plane (guide field) [29–31], is necessary to yield the complete picture of the reconnection processes [10, 32].

Much of the magnetic reconnection research done to date has been motivated by its relevance to a wide range of plasma dynamical phenomena observed in space and the laboratory [33–35]. One of the more spectacular phenomena observed is the solar flare [20,36], which is observed as a sudden brightening in the solar corona and chromosphere. It is associated with a rapid release of up to $6 \cdot 10^{25}$ Joules of energy (equivalent to approximately a sixth of the total energy output of the Sun per second), with the energy stored in the magnetic field of the corona being the only plausible source [38]. In this process, high energy charged particles are ejected into space which, as they reach earth, can have an impact on power grids, communications systems and satellites [39]. Small solar flares, called nanoflares [40, 41], are believed to be responsible for the high temperature of the Sun’s corona compared to the surface [42, 43].

![Figure 1.1: (a) Energy release of a solar flare imaged by AIA (Atmospheric Imaging Assembly) in EUV and RHESSI (Ramaty High Energy Solar Spectroscopic Imager) in X-rays. RHESSI X-ray fluxes are shown as red (blue) contours in the 4 – 10 (10 – 20) keV bands. Bright flare regions observed on the surface are superimposed in pink. Taken from ref. [5]. (b) Schematic of magnetic reconnection in Earth’s magnetopause. The reconnection occurs at the X-points in the shaded regions. Adapted from ref. [37].](image-url)
flare is depicted in figure 1.1a. It shows the hot flare loops, with regions with hard X-ray fluxes and bright flare regions observed on the surface superimposed.

At Earth's magnetopause, the interaction of the solar wind and the interplanetary magnetic field (IMF) with the dipole magnetic field of the Earth leads to magnetic reconnection [44–46]. A schematic is shown in figure 1.1b. The reconnection occurs at the magnetic X-points formed at the day- and nightside. At the dayside, reconnection allows for the particles of the solar wind to enter the magnetosphere. At the nightside, fast energetic particles accelerated by magnetic reconnection stream along the field lines to the poles where they excite atoms in the high altitude atmosphere and contribute to the aurora [47]. Figure 1.1b shows a simplified picture of magnetic reconnection with antiparallel field lines, but a guide field component perpendicular to the plane also plays a role [48,49]. Magnetic reconnection also occurs on the smaller length-scales of tokamak experiments where the interaction between the plasma and the magnetic field during sawtooth oscillations leads to plasma heating and particle acceleration which significantly degrade confinement [34,50–52].

The direct observation of reconnection in space and fusion plasmas is often hampered by technical limitations and harsh environments. Dedicated experiments enable a controlled environment and allows for a detailed study of the fundamental mechanisms involved in reconnection. A new linear device, VINETA II, has been designed for the study of magnetic reconnection and set up within the framework of this thesis. The focus of the present work is on the global description of reconnection in VINETA II. This entails characterizing the dynamics of the magnetic fields, the electric fields and the currents flowing in the plasma. The plasma currents and the current sheet geometry are of particular importance, and are influenced by a number of global factors and boundary conditions. Particular attention is devoted to the role played by the guide field during a magnetic reconnection event, the consequences of a plasma profile with steep gradients, and the impact of the axial boundaries and the size of the current source on the amplitude and geometry of the current sheet.

The thesis is organized as follows: Chapter 2 gives a brief overview of reconnection in collisional and collisionless plasma. Chapter 3 gives a description of the design criteria for a reconnection experiment and how this is realized in the VINETA II device, followed by an outline of the diagnostic tools used in Chapter 4. The experimental findings are
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compiled in chapter 5: It starts with a description of the localized high density plasma and the current sheet without a magnetic X-point topology. The effect of the X-point topology and how the current sheet geometry is modified by the ratio of the guide and in-plane field is then characterized. This is followed by an investigation of the time-evolution of the current sheet geometry and the plasma currents. The next part studies the plasma parameters during reconnection, followed by a discussion of the in-plane velocities and currents. Finally, the observed reconnection rates are discussed.
Chapter 2

Basics of magnetic reconnection

As already pointed out, the mechanisms behind reconnection are not well understood. This is particularly the case for collisionless plasmas where magnetic reconnection has been found to proceed much faster than predicted based on simple MHD models. This chapter gives a basic description of reconnection in collisional and collisionless plasmas. It starts with the motion of field lines in resistive MHD and a two-dimensional model where reconnection is facilitated through classical resistivity. The focus then shifts to collisionless reconnection, explaining the need for a generalized Ohm’s law and giving a description of a model of magnetic reconnection with two-fluid effects. The next part outlines how reconnection is modified by the addition of a third magnetic field component. Finally, various laboratory experiments and astrophysical plasmas are put in context of the different reconnection regimes.

2.1 Induction equation

As a starting point, a simple expression for the motion of field lines in a plasma can be obtained from the Maxwell-Faraday equation

$$\frac{\partial \mathbf{B}}{\partial t} = -\nabla \times \mathbf{E},$$

(2.1)
combined with Ohm’s law for non-zero resistivity

\[ E + v \times B = \eta j, \quad (2.2) \]

yielding

\[ \frac{\partial B}{\partial t} = -\nabla \times (\eta j - (v \times B)). \quad (2.3) \]

In terms of notation, \( B \) is the magnetic induction and is often referred to as the “magnetic field”, \( E \) is the electric field, \( j \) is the current density, \( \eta \) is the plasma resistivity and \( v \) is the velocity. Using Ampere’s law \( \nabla \times B = \mu_0 j \), an expression called the “induction equation” is obtained

\[ \frac{\partial B}{\partial t} = \nabla \times (v \times B) + D_m \nabla^2 B, \quad (2.4) \]

where \( D_m = \eta / \mu_0 \) is the magnetic diffusivity. The right hand side describes the motion of the field lines due to convection and diffusion, respectively [53]. For a characteristic length scale \( L_0 \), the spatial derivatives are approximated by \( 1 / L_0 \), thus giving \( \nabla \times (v \times B) \sim v_0 B / L_0 \) and \( D_m \nabla^2 B \sim D_m B / L_0^2 \). The ratio of the relative amplitude of the two terms is the Magnetic Reynolds number

\[ R_m = \frac{v_0 L_0}{D_m} = \frac{\mu_0 v_0 L_0}{\eta}, \quad (2.5) \]

which is a measure of the relative importance of magnetic convection and diffusion effects [54]. In many space plasmas, where the length scales and the conductivity are usually very large, \( R_m \) is also large and convection is dominant [33]. For such a case, the field lines can not move across the plasma since the motion sets up plasma currents opposing the change. This is called the “frozen-in-flux” condition and the field lines are tightly coupled to the plasma [55]. Hence, a motion of the plasma will bring the field lines with it.

On the other hand, if the diffusion term is dominant, the field lines are able to diffuse across the plasma. For \( R_m << 1 \), the convection term in the induction equation can be set to zero and the diffusion term determines the time scale for the magnetic field motion. This yields a diffusion time of

\[ \tau_D = \frac{\mu_0 L_0^2}{\eta}, \quad (2.6) \]
and a diffusion velocity

\[ v_d = \frac{L_0}{\tau_D} = \frac{\eta}{\mu_0 L_0}. \] (2.7)

Since the magnetic Reynolds number is usually large in typical astrophysical plasmas, the diffusion time is very long. For example, in the solar corona, a typical length scale of \( L_0 = 10^7 \) m, a coronal temperature of \( T = 10^6 \) K and a diffusivity of \( D_m = 1 \) m\(^2\)/s yields a diffusion time on the order of \( \tau_D = 10^{14} \) s, which is several orders of magnitude larger than observed [53].

### 2.2 Reconnection in resistive MHD

Magnetic reconnection in resistive MHD can be described as a combination of: (a) a global region where ideal MHD and frozen-in-flux are valid and the field lines move with the plasma, and (b) the local non-ideal region of the current sheet, formed at the boundary where the anti-parallel field lines cancel out, and where the resistivity is large enough for magnetic diffusion to be significant.

The size of the current sheet is determined by the length scale at which the magnetic field lines are able to move across the plasma and reconnect. In the resistive MHD model, the convective motion results in a thinning of the current sheet (and thus an increase in the current density) until it reaches a length scale where field line diffusion is significant [53].

#### 2.2.1 Sweet-Parker Model

The Sweet-Parker model describes steady state reconnection in resistive MHD [20]. The key element is a two-dimensional geometry where reconnection occurs in a long and thin diffusion region of length \( 2L \) and width \( 2l \), as shown in figure 2.1. It is further assumed that the plasma can be treated as an incompressible fluid, mass is conserved and the pressure is uniform. The plasma moves vertically into the diffusion region with the convective velocity \( v_{\text{in}} \), is accelerated along the layer and then expelled horizontally.
at an outflow velocity $v_{\text{out}}$. The equation of motion for a steady state plasma is given by

$$\rho (\mathbf{v} \cdot \nabla) \mathbf{v} = -\nabla p + j \times \mathbf{B}, \quad (2.8)$$

where $\rho$ is the mass density of the plasma. Mass conservation implies that the rate at which the plasma enters the diffusion region is equal to the rate at which it is expelled. This yields the expression

$$Lv_{\text{in}} = lv_{\text{out}}. \quad (2.9)$$

Since the diffusion region needs to be thin ($l \ll L$) for magnetic diffusion to be significant, the outflow velocity is much faster than the inflow. From the assumption of a uniform pressure, balancing the upstream magnetic pressure with the downstream dynamic pressure gives

$$\frac{B_{\text{in}}^2}{2\mu_0} \sim \frac{\rho v_{\text{out}}^2}{2}, \quad (2.10)$$

which yields an expression for the outflow velocity $v_{\text{out}} = B/\sqrt{\mu_0 \rho} = v_A$, which is the Alfvén speed. From equation (2.9) and the diffusion velocity given by equation (2.7), the inflow velocity is

$$v_{\text{in}} = \sqrt{\frac{\eta v_A}{\mu_0 L}} = \frac{v_A}{\sqrt{S}}, \quad (2.11)$$

where

$$S = \frac{\tau_D}{\tau_A} = \frac{\mu_0 L v_A}{\eta}, \quad (2.12)$$
is the Lundquist number and is given by the ratio of the resistive diffusion time scale $\tau_D$ to the Alfvén time scale $\tau_A$. This is the same as the magnetic Reynolds number with $L$ as the characteristic length scale and $v_A$ as the characteristic velocity. The Sweet-Parker reconnection rate is then given by

$$\frac{v_{in}}{v_A} = \frac{l}{L} = \frac{1}{\sqrt{S}}.$$  

(2.13)

An issue with the Sweet-Parker model is that the predicted reconnection rate $v_{in}/v_A$ is far too slow to explain the time scales of most observed events attributed to magnetic reconnection [1]. For example, the energy release observed in solar flares occurs on the order of minutes as opposed to days as predicted by the model [53]. Despite the problems that arise with the model, it is a useful starting point for understanding reconnection and the complex mechanisms involved. Furthermore, by relaxing some of the assumptions made and using a generalized Sweet-Parker model incorporating compressibility, downstream pressure, and anomalous resistivity, it shows good agreement with experimental results for reconnection in a highly collisional regime [20–22]. However, since most space plasmas and fusion devices have high Lundquist numbers and can be considered collisionless, a description beyond the Sweet-Parker model is usually required.

2.3 Beyond resistive MHD

2.3.1 Generalized Ohm’s law

A plasma can be thought to be collisionless when the mean free path of the particles is much greater than the characteristic length scale [1]

$$\lambda_{mfp} = \frac{v}{\nu_{coll}} \gg L$$  

(2.14)

where $v$ is the particle velocity and $\nu_{coll}$ the collision frequency. For a collisionless plasma, classical resistivity is not significant and equation (2.2) is no longer sufficient to describe the relationship between the current and the electric fields in the system. Instead, a generalized Ohm’s law, where additional terms are taken into account, is re-
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<table>
<thead>
<tr>
<th>Parameter</th>
<th>Dimension</th>
<th>Laboratory Experiments</th>
<th>Earth’s Magnetosphere</th>
<th>Solar Corona</th>
<th>Solar Interior</th>
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<td>Plasma density (n_0)</td>
<td>m(^{-3})</td>
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<td>(10^{16})</td>
<td>(10^{20})</td>
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<tr>
<td>Temperature (T_0)</td>
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<td>(10^4)</td>
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<tr>
<td>Length scale (L_0)</td>
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<td>(10^{-1})</td>
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<td>(10^8)</td>
<td>(10^7)</td>
</tr>
<tr>
<td>Magnetic field (B_0)</td>
<td>T</td>
<td>(10^{-1})</td>
<td>(10^{-8})</td>
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<td>(10^1)</td>
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<tr>
<td>Lundquist number (S)</td>
<td>(10^3)</td>
<td>(10^{13})</td>
<td>(10^{14})</td>
<td>(10^{13})</td>
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</tr>
</tbody>
</table>

\[ L_{\text{resistivity}} = \beta^{1/2} \delta_e \delta_i / \lambda_{\text{mpf}} \]
\[ L_{\text{Hall}} = c / \omega_{\text{pi}} \]
\[ L_{\text{pressure}} = \beta^{1/2} r_{L,i} \]
\[ L_{\text{inertia}} = c / \omega_{\text{pe}} \]

Table 2.1: Typical plasma parameters and characteristic length scales of the Generalized Ohm’s law. Laboratory experiments refer to TS-3, MRX and SSX (cf. section 3.1). Adapted from ref. [53].

The generalized Ohm’s law is equivalent to the equation of motion of the electrons and is given by [56]

\[ \mathbf{E} + \mathbf{v} \times \mathbf{B} = \eta \mathbf{j} + \frac{1}{ne} \mathbf{j} \times \mathbf{B} - \frac{1}{ne} \nabla \cdot \mathbf{p}_e + \frac{m_e}{ne^2} \frac{\partial \mathbf{j}}{\partial t}. \]  

(2.15)

The left-hand side is the electric field in the moving frame. The first term on the right-hand side is the electric field associated with Ohmic dissipation caused by electron-ion collisions. The second term is the Hall term and arises from treating the ions and electrons as two separate fluids. The next term comes from the force due to gradients in the electron pressure, where \( \mathbf{p}_e \) is the electron pressure tensor. The last term on the right hand side is the electron inertia effect. The relative importance of the terms depends on the plasma parameters as well as the length- and timescales involved. An estimate of which of the non-ideal terms are significant can be obtained by equating the individual terms with the convective electric field \( \mathbf{v} \times \mathbf{B} \) [53]. For example, for the Hall term \( \mathbf{j} \times \mathbf{B} / ne \), using \( j \sim B_0 / \mu_0 L_0 \) gives

\[ \frac{B_0^2}{ne \mu_0 L_0} \sim v_0 B_0, \]

(2.16)

and using the Alfvén speed, \( v_A \), as the characteristic velocity yields

\[ L_{\text{Hall}} \sim \sqrt{\mu_0 m_i n e} \sim c \sqrt{\frac{\epsilon_0 m_i}{ne^2}} \sim \delta_i, \]

(2.17)
where $\delta_i = c/\omega_{pi}$ is the ion-inertial length (or ion skin depth), $\omega_{pi}$ is the ion plasma frequency and $c = 1/\sqrt{\varepsilon_0 \mu_0}$ is the speed of light. Hence, the Hall term becomes important when the width of the current sheet is comparable to the ion-inertial length. Similarly, the electron inertia term $(m_e/ne^2)\partial j/\partial t$, with $\partial/\partial t \sim v_0/L_0$, yields

$$\frac{m_e v_0 B_0}{ne^2 \mu_0 L_0^2} \sim v_A B_0,$$

and

$$L_{\text{inertia}} \sim c \left( \frac{m_e e_0}{ne^2} \right)^\frac{1}{2} \sim \delta_e,$$

where $\delta_e = c/\omega_{pe}$ is the electron-inertial length and $\omega_{pe}$ is the electron plasma frequency. For the pressure term $\nabla \cdot p_e/ne$, using $\nabla \sim 1/L_0$ and $p \sim nk_B T$, one obtains

$$\frac{nk_B T}{ne L_0} \sim v_A B_0,$$

and

$$L_{\text{pressure}} \sim \beta^{1/2} r_{L,i},$$

where $\beta = nk_B T/(B_0^2/2 \mu_0)$ is the ratio of the plasma pressure and magnetic pressure and $r_{L,i}$ is the ion-gyro radius. Lastly, for $\eta j$

$$\frac{\eta B_0}{\mu_0 L_0} \sim v_A B_0,$$

and using the Spitzer resistivity

$$L_{\text{resistivity}} \sim \beta^{1/2} \frac{\delta_e \delta_i}{\lambda_{mpf}}.$$  

Table 2.1 lists various plasma parameters along with the characteristic length scales of the terms in the generalized Ohm’s law for four cases. The magnetosphere is an example of a collisionless environment, while the solar interior can be considered highly collisional.
2 Basics of magnetic reconnection

![Diagram of two-fluid model](image1.png)

**Figure 2.2:** (a) Geometry of two-fluid model. Adapted from ref. [57]. (b) Computer simulation of out-of-plane quadrupolar magnetic field. The coordinates are in units of $δ_i = c/\omega_{pi} = 1$. Adapted from ref. [24].

### 2.3.2 Two-fluid model

Two-fluid effects offer a possible explanation for the fast reconnection rate observed in collisionless plasmas. A schematic of the geometry is shown in figure 2.2a. The ions move through the broader ion diffusion region while the magnetic field frozen into the electron fluid continues inward and reconnects within the smaller electron diffusion layer. Since the length of the electron diffusion region is much smaller than the global plasma length scale, the bottleneck issue of Sweet-Parker reconnection is eliminated and reconnection can proceed much faster. The differential motion of the ion and electrons give rise to in-plane currents, which in turn gives an out-of-plane electric field through the Hall term in the generalized Ohm’s law and helps facilitate faster reconnection. The in-plane currents produce an out-of-plane quadrupole magnetic field which is taken as an indicator of the presence of two-fluid effects. This has been observed in both computer simulations [58, 59] and in experiments [60, 61]. A simulation of the out-of-plane quadrupolar magnetic field is shown in figure 2.2a, and the peaks of the magnetic field are located along the separatrix.

### 2.3.3 Current sheet instabilities

Recent work has shown that the stability of the current sheet, which is highly affected by the small-scale dynamics inside, plays an important role in magnetic reconnection.
2.4 Reconnection with a guide field

So far magnetic reconnection has been discussed in two-dimensional geometry with antiparallel field lines. However, the presence of a magnetic field component (guide field) directed perpendicular to the reconnection plane may significantly affect the reconnection dynamics. This is particularly relevant for reconnection in the magnetosphere [48, 65], where the perpendicular component has a magnitude comparable to the in-plane reconnection field, and for the problem of sawtooth crashes in tokamaks where the guide field is typically an order of magnitude larger than the reconnection field [52].

Figure 2.3: Field lines in circular current sheet at \( z = 0 \) mapped to \( z = \pm 2d \). The current sheet shape is shown in red, the blue lines represent the magnetic field lines of the combined field and the black lines the in-plane X-point topology.
2 Basics of magnetic reconnection

2.4.1 Magnetic mapping

With the addition of a guide field, there is no longer a magnetic null-point and the electric field has a component parallel to the magnetic field. The combined guide field and X-point topology results in a change in the shape of the current sheet through magnetic mapping due to the free motion of the particles along the field lines. This is illustrated in figure 2.3 where the field lines in a circular current sheet at $z = 0$ have been mapped to $z = \pm 2d$. Electrons moving along a magnetic field line are shifted in the $xy$-plane (azimuthal plane) and the original shape of the current density distribution is changed. The current sheet becomes squashed along one separatrix and elongated along the other, thereby changing its width and length [66].

2.4.2 Influence of guide field on reconnection in two-fluid regime

It has been found in both experiments [67,68] and computer simulations [29,69–71] that a guide field affects reconnection in a two-fluid regime by modifying the electron flow.

Figure 2.4: Top: Contours of the toroidal field measured in MRX for guide fields spanning $B_g \sim 0$ (left) to $B_g \sim B_{rec}$ (right). Bottom: Quadrupole field patterns produced using a Hall-MHD simulation. The contour scales are independently determined for each plot to maximize contrast. The contours have been plotted with the same spatial scale between simulation and experiment, setting $d_i = 5$ cm. Taken from ref. [67].
2.4 Reconnection with a guide field

Figure 2.5: Density contours measured in the CS-3D experiment for different directions of the guide field. The difference between neighboring lines is $\delta n_e = 0.2 \times 10^{10} \text{m}^{-3}$. Taken from ref. [68]

It has also been observed that the reconnection rate is significantly reduced when a guide field is included [67, 69]. However, the specific dependence on the guide field strength is not well understood. It has been proposed that the interaction of the Hall currents and the guide field result in a $\mathbf{j}_\perp \times \mathbf{B}_g$-force which opposes the reconnection flow [67, 68] and thereby reduced the reconnection rate. In the same studies, it has been suggested that this force is responsible for the tilt of the current sheet observed in both simulations and experiments. An example is shown in figure 2.5 for current sheet measurements at the CS-3D experiment. When a guide field is applied, there is a clear tilt of the current sheet, with the tilt direction depending on the direction of the guide field.
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2.5 Reconnection regimes

In addition to the conventional single X-line collisional and collisionless cases described previously, reconnection in space and laboratory devices can be subdivided in different regimes such as single/multiple X-line reconnection and hybrid branches [2,72]. The two key dimensionless parameters for the division of the regimes are the Lundquist number \( S \) and the effective plasma size \( \lambda \) given by

\[
\lambda = \frac{L}{\rho_s}.
\]  

(2.24)

Here \( L \) is the plasma size and \( \rho_s = c_s/\omega_{ci} \) the ion sound gyro radius, with \( c_s \) the ion sound speed, and \( \omega_{ci} \) the ion gyro frequency. The length of the current sheet is taken as \( L_{CS} = \varepsilon L \) where \( \varepsilon \) is typically chosen between \( 0 \leq \varepsilon \leq 0.5 \) [72]. For reconnection without a guide field, the transition between single X-line collisional and collisionless reconnection occurs when the current sheet length scale approaches the ion skin depth [72]. With a non-zero guide field, the transition instead occurs at the ion sound gyro radius [72, 73].

![Phase diagram showing various laboratory, heliophysical and astrophysical plasmas in which reconnection is believed to occur. Taken from ref. [72]](image)

Figure 2.6: Phase diagram showing various laboratory, heliophysical and astrophysical plasmas in which reconnection is believed to occur. Taken from ref. [72]
Figure 2.6 shows a phase diagram indicating the different reconnection regimes and where various laboratory, heliophysical and astrophysical plasmas belong. They are roughly divided into three groups. The first group includes high temperature fusion plasmas and the plasmas of Earth’s magnetosphere, and lie in the collisionless regime. The reconnection in these plasmas either occur with a single or multiple X-lines, depending on whether the plasma size is is larger than a critical plasma size $\lambda_c$. The second group are clustered along the black line given by $S = \varepsilon^2 \lambda^2$ (in the figure $\varepsilon = 0.5$) which separates the multiple X-line collisionless phase and the multiple X-line hybrid phase. It spans a huge range of plasmas from the solar corona to extragalactic jets. For small $\lambda$ and $S$ this line separates the single X-line collisional/collisionless phases. Due to their low Lundquist numbers and small plasma sizes, most reconnection experiments lie in these two regimes. The last group of plasmas lie in the multiple X-line collisional phase and e.g includes the solar chromosphere, molecular clouds and gamma ray bursts.
Chapter 3

Experimental Setup

The experimental setup of VINETA II is described in this chapter. It starts with an overview of some of the past and present experiments dedicated to the study of magnetic reconnection in section 3.1 and outlines some of the important design criteria and how reconnection is achieved in other devices. This is followed by a description of the new device VINETA II in section 3.2 and the reconnection drive methods in section 3.3. Section 3.4 addresses the sheath limited current and the use of a plasma gun as an additional source of electrons. Finally, the reconnection regimes accessible in VINETA II are discussed in section 3.5.

3.1 Overview of dedicated reconnection devices

Dedicated experiments provide a well-controlled plasma environment which allows for a detailed study of the fundamental mechanisms involved in reconnection. Various experimental approaches and configurations can be utilized, and the design is a compromise between technical limitations and the requirements for the specific subject to be investigated. The experimental configurations mostly consist of either toroidal or cylindrical geometry, which all offer their own advantages and drawbacks. For example, a linear device has a simpler geometry whose uniform guide field makes it ideal for the study of
guide field reconnection. However, such a device also has the issue of open end boundary conditions which limits the inductively driven axial current in the system as well as the achievable plasma temperature due to the lack of axial confinement. On the other hand, in a toroidal device, the plasma current can close toroidally and the axial confinement allows for higher temperatures. The drawback is that geometrical effects come into play and the boundary conditions are less well-defined than for a linear device.

Common to all reconnection experiments is a need for a reconnection drive. The distinction between driven and spontaneous reconnection can be defined as the former being due to some externally forced motion, e.g. through an $\mathbf{E} \times \mathbf{B}$-drift by an externally induced electric field, while the latter is due to some internal instability or loss of equilibrium and only has a weak dependence of the external coupling [1]. From the experimental point of view, the advantages of a controlled driving of the reconnection is the better reproducibility.

There are a number of past and present day devices devoted to the study of magnetic reconnection. One of the earliest experiments was carried out by Stenzel and Gekelmann on a linear device at UCLA [74]. Reconnection was achieved by a oscillating current through two parallel conductors. Plasma was generated using an oxide coated cathode, which also provided the necessary plasma current in response to the inductive field. The ions were unmagnetized in the experiment and the plasma was in the electron magnetohydrodynamics (EMHD) regime. Another linear device dedicated to the study of reconnection is the Three-Dimensional Current Sheet experiment (3D-CS) at the General Physics Institute Russian Academy of Sciences [75]. The experiment is dedicated to the study of current sheet formation and reconnection in three-dimensional geometry. An X-point topology is generated by parallel conductors and reconnection is driven

| Table 3.1: Relevant parameters of devices dedicated to the study of magnetic reconnection. n/a indicates that no information is available. |
|---------------------------------|---|---|---|---|---|
| **Geom.** | **Linear** | **Linear** | **Toroidal** | **Toroidal** | **Spheromak** | **Spheromak** |
| $L$ | 0.3 m | 0.2 m | 0.1 m | 0.3 – 0.5 m | 0.1 m | 0.12 m |
| $B$ | < 2 mT | < 300 mT | < 100 mT | 20 – 100 mT | 50 – 100 mT | 50 – 100 mT |
| $n_e$ | $10^{18}$ m$^{-3}$ | $10^{20}$ – $10^{21}$ m$^{-3}$ | $10^{18}$ m$^{-3}$ | $10^{19}$ – $10^{20}$ m$^{-3}$ | $10^{20}$ – $10^{21}$ m$^{-3}$ |
| $T_e$ | 10 eV | 10 eV | 20 eV | 5 – 15 eV | 10 eV | 10 – 30 eV |
| $\lambda_{mfp,e}$ | 2 m | n/a | 5 – 50 m | 0.01 – 0.2 m | 0.002 m | 0.1 m |
| $S$ | 1 – 10 | n/a | 1000 | 1000 | 700 | 1000 |
3.1 Overview of dedicated reconnection devices

by a plasma current generated by a pulsed voltage between two electrodes. The Versatile Toroidal Facility (VTF) at the Massachusetts Institute of Technology (MIT) [15], is used for investigating fast magnetic reconnection in a collisionless plasma. It utilizes a strong guide field and an X-point topology produced by toroidal coils. Reconnection is either driven by an additional toroidal coil whose inductive electric field induces a plasma current at the X-point, or by varying the current generating the X-point topology. The plasma is generated through electron cyclotron resonance heating (ECRH). Another toroidal device is the Magnetic Reconnection Experiment (MRX) at the Princeton Plasma Physics Laboratory (PPPL) [14]. The experiment is used for studies of reconnection in the two-fluid regime. The plasma is generated through induction by pulsing currents through toroidal field coils. A poloidal magnetic field is established by additional coils and the reconnection drive is achieved by increasing or decreasing this field. Another means for achieving reconnection is through the merging of spheromaks, where an X-point configuration is formed between the two plasma tori. This method is used for TS-3/4 at the University of Tokyo [13] and the Swarthmore Spheromak Experiment (SSX) at Swarthmore College [17]. Some of the relevant parameters for the experimental devices listed above are given in table 3.1.
3 Experimental Setup

3.2 VINETA II

The experiments were performed in the new linear device VINETA II, which was set up within the framework of this thesis. A schematic of VINETA II is shown in figure 3.1. It consists of a cylindrical stainless steel vacuum vessel (reconnection module), with a diameter of 1 m and a length of 1.6 m, in combination with two modules from the original VINETA experiment (Versatile Instrument for Studies on Nonlinearity, Electromagnetism, Turbulence and Applications) [76]. One side of the device is used for a turbomolecular pump, which provides a base pressure of \( p_b \approx 5 \cdot 10^{-5} \) Pa. The other side is used for a radio frequency (rf) plasma source. A large rectangular port serves as the main diagnostics access and is used for a positioning system that allows for probes to be scanned through the azimuthal plane of the device.

Magnetic guide field

The magnetic field coils shown in figure 3.1 generate a homogeneous axial guide field of \( B_g \leq 100 \) mT. The calculated guide field in the \( xz \)-plane is shown in figure 3.2a and the profile along the device axis is shown in figure 3.2b. The smaller coils are positioned to reduce the spatial ripple of the field to \( \tilde{B} < 3\% \) on axis in the transition region between
the reconnection module and the smaller modules. A gradient of the magnetic field with a maximum of \( \frac{\partial B_g}{\partial z} = 91.2 \text{ mT/m} \) is seen at the edge of the reconnection module on the source side. If needed, this magnetic field gradient can be eliminated by placing one of the smaller modules on the source side.

### Radio-frequency plasma source

Plasma is generated through rf-heating \( (f_{rf} = 13.56 \text{ MHz}) \) using either a helicon antenna [78] or a planar spiral antenna [77]. As neutral gas, Argon at a pressure of \( p \approx 0.15 \text{ Pa} \) is used. The helicon plasma source consists of a \( m = +1 \) helical antenna with a diameter of 10 cm. It can be operated either in a capacitive, inductive or helicon wave sustained discharge mode [78]. In the helicon mode, a plasma with an electron temperature of \( T_e \approx 3 \text{ eV} \) and a plasma density of \( n \leq 10^{19} \text{ m}^{-3} \) is generated. The chief rf-plasma source utilized for the measurements presented in this thesis was the spiral antenna. With a diameter of 20 cm, it allows for a wider plasma column with flat central temperature and density profiles. It can be operated either in the inductive mode (with some additional contribution due to capacitive coupling) or in the helicon mode [77]. Typical radial density profiles of the two modes are shown in figure 3.3. In the inductive discharge
3 Experimental Setup

Figure 3.3: Typical radial plasma density profile for the inductive mode (blue line) and helicon mode (red line) using the spiral antenna ($P_{rf} = 0.8$ kW and $1.8$ kW, respectively). The helicon wave sustained mode using the helicon antenna is plotted in black ($P_{rf} = 1.8$ kW). Figure taken from ref. [77].

mode a plasma density of $n \approx 10^{16}$ m$^{-3}$ and an electron temperature of up to $T_e \leq 6$ eV are achieved. This allows for the generation of a plasma with lower collisionality than with the helicon antenna. In the helicon mode, the antenna generates a plasma with a peak density of $n \approx 5 \cdot 10^{18}$ m$^{-3}$ and a temperature of $T_e \approx 2-3$ eV.

3.3 Reconnection drive

Reconnection in VINETA II is achieved by externally driving magnetic field lines towards an X-point, thereby inducing a current in the background plasma and altering the magnetic field topology. The advantage of this particular driving scheme is the high reproducibility. This allows for the profile of the relevant parameters to be obtained by scanning probes thought the azimuthal cross section of the device on a shot-by-shot basis, thereby eliminating the need for a probe array and reducing the disturbance of the plasma. Another advantage is that, through the time scale of the reconnection drive, it is possible to ensure that the field line motion is faster than the magnetic diffusion velocity. In terms of length scales, this means that the distance the magnetic field lines diffuses on the reconnection drive time scale $\tau_{drive}$, i.e., the resistive layer width $L_{res}$, is smaller than the collisionless length scales (cf. section 3.5). The resistive layer width is given by the expression for the magnetic diffusion length scale, (equation 2.6), with the diffusion time replaced by the reconnection drive time.
3.3 Reconnection drive

3.3.1 Reconnection Drive Setup

Two approaches are utilized for driving the reconnection in VINETA II. For the first method, called setup A, an oscillating current with a peak amplitude of $I \leq 1 \text{kA}$ through two parallel conductors generates an X-point topology and simultaneously drives the magnetic field lines (similar to the UCLA experiment). For the second approach, called setup B, a nearly constant current of $I \leq 2 \text{kA}$ through the conductor pair establishes a stationary X-point magnetic field topology. An oscillating current through two additional parallel conductors (similar to the VTF experiment) then modulates the magnetic field and generates an inductive electric field $E_{\text{ind}}$ which drives the reconnection.

The wire configuration used for setup A is shown in figure 3.4a, and consists of a pair of parallel flat conductors connected in series. They have a width of 10 cm and are positioned with a 30 cm radial separation at an angle of approximately $120^\circ$ with respect to the $x$-axis. Flat conductors were used in order to reduce the contribution of the skin effect to the resistivity. For setup B, shown in figure 3.4b, two cylindrical conductor pairs are utilized. The conductors used for the X-point topology are placed at the same positions as for setup A, and the drive conductors are positioned at a $90^\circ$ angle to these at radial positions of $r = \pm 0.4 \text{m}$. All conductors for setup B are placed within vacuum tight stainless steel tubes. These are grounded to the wall on one side and provide electrostatic shielding.

Figure 3.4: Reconnection module with the conductors producing the reconnection magnetic field for: (a) setup A and, (b) setup B. The rf-plasma source is placed on the right-hand side of the module. The copper colored conductors generate the X-point topology and example in-plane magnetic field lines are plotted at the axial position of the probe positioning system. For setup B, the drive coils are colored in black.
3 Experimental Setup

Figure 3.5: The calculated modulus of the in-plane magnetic field of the conductors for: (a) the conductors generating the the X-point topology \((I = 2 \text{kA})\), and (b) the drive conductors \((I = 1 \text{kA})\). The wall of the reconnection module is indicated by the blue circle. Example field lines are plotted in black.

The calculated modulus of the magnetic field as created by the conductors for setup B, is shown in figure 3.5. Figure 3.5a shows the X-point topology established for a \(I = 2 \text{kA}\) current, and figure 3.5b the absolute magnetic field of the reconnection drive conductors for a peak current of \(I = 1 \text{kA}\). The in-plane magnetic field is largest at the conductors and zero at the X-point. Due to mirror currents induced in the wall, the magnetic field approaches zero at the edge of the vessel. This, however, has not been considered in the simple calculation shown in figure 3.5. The in-plane magnetic field at the plasma edge (assuming a characteristic length scale of \(L = 10 \text{cm}\)) is approximately \(B_{xy} = 2 \text{mT}\). This is more than one order of magnitude larger than the peak magnetic field of the reconnection drive at the same position.

The placement of the conductors producing the X-point topology, along with the large dimensions of the reconnection module, makes it possible to establish a closed field line configuration in the plane, which reduces plasma transport to the walls along the reconnection field lines. For both setups, the drive of the reconnection is independent of the plasma generation, i.e., plasma parameters such as density and temperature can be varied without affecting the drive. For setup B the inductive field can also be set independently of the amplitude of the in-plane magnetic field.
3.3 Reconnection drive

Figure 3.6: (a) Schematic of a PFN circuit. (b) Calculated discharge voltage of the individual capacitors of a PFN. (c) Calculated (green) and measured (blue) current produced by the PFN.

3.3.2 Electronics

Pulse forming network

The conductor current which generates the X-point topology for setup B is supplied using a pulse forming network (PFN) [79, 80]. A schematic of a PFN is shown in figure 3.6a. The switching is done using a thyristor, and the system utilizes inductor-capacitor (LC) resonator circuits in a ladder network in order to create delays between the energy release of the capacitors. This allows for control of the output power in terms of duration, magnitude, and shape. Thereby, a rectangular output pulse can be achieved [80].

The simplest type of PFN is the so-called Rayleigh line PFN, where all capacitances and inductances are of equal value [81, 82]. Other, more complicated designs, exist which can be utilized to reduce the flattop ripple and the rise/fall times of the output pulse by optimizing the component values and the number of elements used [83]. The calculated
3 Experimental Setup

Figure 3.7: (a) Schematic of reconnection drive circuit. (b) Frequency dependence of the maximum output current of the drive.

discharge voltage of the separate capacitors of a PFN is shown in figure 3.6b. The resulting calculated and measured current is shown in figure 3.6c and show good agreement. The discrepancy in amplitude and the longer rise/fall times of the measurement are due to the parasitic resistance and inductance of the circuit which are not taken into account in the simulation.

Reconnection drive electronics

The oscillating current used for the reconnection drive is produced via a resonant $LC$-circuit. A schematic is shown in figure 3.7a. The switching is done using insulated-gate bipolar transistors (IGBT). The inductance of the reconnection drive conductors ($L_c \approx 7 \mu H$) in combination with a variable capacitance ($C_{\text{var}} = 250 - 750 \mu F$) allows for a frequency range of $f = 60 - 100 \text{kHz}$. The frequency dependence of the maximum of the modulus of the current is shown in figure 3.7b. Due to the increasing impedance of the resonance circuit, the current amplitude decreases with increasing drive frequency. The choice of the frequency is based on a compromise between the time resolution of the available diagnostics and the data acquisition system (DAQ), and driving the field
3.3 Reconnection drive

lines faster than resistive magnetic diffusion.

### 3.3.3 Phases of the reconnection drive

The reconnection drive current and the current used for producing the X-point topology are shown in figure 3.8a. The PFN current is observed to be nearly constant during the reconnection drive pulse. Figure 3.8b shows the time-evolution of the drive current and the associated inductive electric field $E_{\text{ind}}$ at the X-point. An initially large electric field is induced as the drive is turned on and then follows the current with a $-\pi/2$ phase shift. A small bump in the electric field can be seen as the current changes direction at $t = 8.5\,\mu s$. This is related to the switching of a second IGBT and the addition of a second current component. Through the capacitance and inductance of the reconnection drive, the second current component can be tuned to counteract the effect of damping and produce a waveform where the two initial peaks have equal amplitude. The switching of the

![Figure 3.8](image-url)  

Figure 3.8: (a) Current producing the stationary X-point topology (blue line) and the current through the reconnection drive conductors (green line). (b) Drive current (green line) and the associated inductive electric field at the X-point (red line). I-IV indicates the different phases of the drive corresponding to either “pull” or “push” reconnection.
IGBTs also results in a high frequency ripple of the inductive electric field (“ringing”).

Depending whether the amplitude of the magnetic field is rising or falling, the magnetic field lines are either pushed away from, or pulled towards the conductors. For setup A, “push” reconnection occurs in the phases indicated by I and III in figure 3.8b, and “pull” reconnection during II and IV. For setup B, push reconnection occurs during the I and IV phases and pull reconnection during the II and III phases.

3.4 Plasma current

3.4.1 Boundary conditions

At the boundaries of the device wall, the sheath forms a potential barrier which electrostatically confines the more mobile electrons and adjusts itself such that the electron and ion fluxes to the wall are equal [84]. The resulting net current to the wall consists of both the electron and ion current density and is zero due to ambipolarity. The ions are accelerated in the pre-sheath and enter the sheath with a velocity greater than the ion sound speed $c_s$, which is the Bohm sheath criterion [84, 85]. The ion current density to

![Figure 3.9: Schematic one-dimensional profile of the plasma potential of a plasma bounded by two walls at $±d$. The blue line shows the plasma potential without an induced net current and the green line shows how the plasma potential is adjusted due to charge separation effects when applying an inductive electric field.](image)
the wall is hence given by [84]

\[ j_i = ne c_s. \]  \hspace{1cm} (3.1)

For electrons with a Maxwellian velocity distribution, the current density of the electrons to the wall is given by [84]

\[ j_e = -ne \sqrt{\frac{k_B T_e}{2\pi m_e}} \exp \left( \frac{-e\phi_p}{k_B T_e} \right). \]  \hspace{1cm} (3.2)

Figure 3.9 shows the one-dimensional plasma potential of a plasma bounded by two walls at \( \pm d \). When no net current is induced, the plasma potential is constant inside the main plasma region and drops off at the wall. If one attempts to drive a current in the system by an inductive electric field \( E_{\text{ind}} \), the charge separation will set up a gradient in the potential, resulting in an electrostatic field opposing the externally applied electric field. Thus, the plasma current that can be induced is limited by charge separation. Assuming that the heavier ions are unaffected by the electric field, the ion current to the walls remains unchanged. The lighter electrons, on the other hand, will be accelerated (to the right in figure 3.9). Thus, they will be impeded from flowing to the wall on the left and instead leave the plasma through the right-hand side. Assuming the number of ions and electrons leaving the plasma remains unchanged, the maximum possible current induced in the plasma is now given by twice the ion current. This sets the upper limit of the plasma current. In order to inductively drive a larger current in the plasma, additional electrons need to be injected (on the left-hand side in figure 3.9) in order to counterbalance the charge separation effects. An important consequence of the axial boundary conditions is that the resistivity \( \eta \) is no longer just given by the inductive electric field but has an electrostatic contribution,

\[ \eta = \frac{1}{j_x} \left( E_{\text{ind}} - \frac{\partial \phi}{\partial z} \right). \]  \hspace{1cm} (3.3)

This significantly affects the interpretation of the resistivity, and if not correctly taken into account may lead to an overestimate of effects such as anomalous resistivity.
3 Experimental Setup

![Diagram of plasma gun and grids]

Figure 3.10: Setup of the plasma gun and the grids. The gray parts are stainless steel and the blue parts are made of insulating Polytetrafluoroethylene (PTFE). A localized high density plasma is generated in the background rf-plasma. The non-negligible inductance of the ground connection is denoted by $L$.

3.4.2 Plasma gun

Due to the above described charge separation effects, an additional electron source is needed to provide the current in response to the inductive electric field. This is achieved using a plasma gun, which produces electrons through an arc discharge in the gas confined in the volume between its cathode and anode [18, 86]. A schematic drawing of the plasma gun is shown in figure 3.10. The particles exit the plasma gun through a ring anode with a diameter of 1.2 cm. The cathode and the anode are separated by a channel of steel washers, which provide a so-called wall-stabilized arc [87] and prevent the arc from forming along a random path within the plasma gun. Gas puffing, using a fast valve at the cathode, provides the high local pressure required for electrical breakdown without significantly increasing the overall background pressure in the device. The voltage required to ignite the arc, i.e., the breakdown voltage, is dependent on the gas pressure and the distance between the electrodes through Paschen’s Law [88]

$$U = \frac{apd}{\ln(pd) + b},$$

(3.4)

where $U$ is the breakdown voltage, $p$ is the pressure, $d$ is the distance between the electrodes and $a$ and $b$ are constants that depend on the composition of the gas. The Paschen curve for argon and a separation between the electrodes of $d = 2$ cm is shown in fig-
3.4 Plasma current

Figure 3.11: Paschen curve for argon and a separation between the electrodes of $d = 2\, \text{cm}$.

A minimum breakdown voltage of $U = 100\, \text{V}$ is required at a gas pressure of $p = 37\, \text{Pa}$. The current for the plasma gun is supplied by a PFN (cf. section 3.3.2).

The plasma gun is essentially a small plasma source which produces a localized high density plasma. The details of the plasma gun discharge mechanism, and in particular dependencies on external factors such as a magnetic field and a background plasma, are not well understood [89]. Triple probe [90] measurements indicate a high density plasma on the order of $n \approx 10^{18}\, \text{m}^{-3}$ and a peak temperature of a few electron volts at the axial position of the probe positioning system (at a distance of $z = 53\, \text{cm}$ from the plasma gun) [91].

The gun is positioned at the X-line and the plasma current closes through a grounded circular wire mesh grid (grid 2) as shown in figure 3.10. The grid has a diameter of $23\, \text{cm}$ and is placed at the rf-source side of the large vacuum chamber at a distance of approximately $z = 1\, \text{m}$ from the gun. Figure 3.12a shows the current flowing through grid 2 when biased to $U_{\text{bias}} = 50\, \text{V}$. Also plotted is the current supplied by the PFN. From figure 3.12a, it is clear the gun is capable of supplying a current in response to an electric field. There is also a current flowing without an external electric field, i.e., when grid 2 is not biased. This current is also plotted in figure 3.12a. The origin of this “background” current is not clear, but it is assumed to be associated with electrostatic fields in the plasma. In order to reduce the background current, an additional grid is placed close to the plasma gun (grid 1 as shown in figure 3.10). The resulting background current flowing to grid 1 and grid 2 is shown figure 3.12b. The background current to grid 2 is now reduced by approximately an order of magnitude ($I_{\text{grid 2}} \approx -10\, \text{A}$) and most of the
Figure 3.12: (a) Current flowing to grid 2 when the grid is biased to $U_{\text{bias}} = 50$ V (red line) and when the grid is unbiased (green line). Plotted for reference is the current supplied by the PFN. (b) Current flowing to ground when an additional grid is placed in front of the gun. The green line shows the current to grid 2 and the blue line the current to grid 1.

current instead flows to grid 1 ($I_{\text{grid 1}} \approx -200$ A).

### 3.4.3 Plasma current limitation

Figure 3.13 shows time traces of the plasma current $I_p$ and the inductive electric field measured at the X-point for setup B. The current of the PFN is not exactly constant, and adds a small contribution to the inductive electric field ($E_{\text{ind}} = 0.5$ V/m at $t = 0$). A small peak can be seen in $E_{\text{ind}}$ at the start of the plasma gun discharge at around $t = 5 \mu$s. As the reconnection drive is triggered, the inductive electric field initially acts to inhibit the background current, and a plasma current is only inductively driven once $E_{\text{ind}}$ becomes negative, i.e., points towards the gun. The consequence of the background current is that there is always a current sheet present, which is then modulated by the reconnection drive. Comparing the inductive field and the extracted current, it is clear that the latter does not follow the former exactly in time. The current only reaches its peak value of $I_p = -20$ A $1.1 \mu$s after $E_{\text{ind}}$ and the inductively driven part of the plasma current continues to flow $2.1 \mu$s after the inductive electric field has changed direction. The explanation for this delay is that the inductance of the external circuit of the plasma gun and grid setup inhibits the rise/fall time of the current: The reactance corresponding
3.4 Plasma current

Figure 3.13: Current from plasma gun (green) and inductive electric field at the X-point (red).

The phase shift can be estimated by

$$\theta = \arctan \left( \frac{X}{R} \right),$$  \hspace{1cm} (3.5)

where $\theta$ is the phase shift, $X$ is the reactance consisting of an inductive and capacitive component and $R$ is the resistance. The Spitzer resistivity, calculated for a density of $n \approx 10^{18}$ m$^{-3}$, an electron temperature of $T_e \approx 4$ eV, and a current channel with a length of 1 m and a cross-sectional area of $A = 10^{-4}$ m$^2$ yields a resistivity of $R = 1.4 \ \Omega$. The phase shift of $\theta = 21^\circ$ and equation (3.5), then results in a plausible reactance of $X = 0.5 \ \Omega$. Similar delays were seen in the VTF experiment, which where attributed to the self-inductance of the toroidal current channel [92].

If the current density is given by $j = E / (\eta_s + \eta_L)$, where $\eta_s$ is the Spitzer resistivity and $\eta_L$ is the resistivity due to collisions with neutrals (negligible in comparison to the Spitzer resistivity for the measurements presented in this thesis), a density and a temperature of $n = 10^{18}$ m$^{-3}$ and $T_e = 4$ eV, and an inductive electric field of $E_{ind} \approx 25$ V/m should result in a current density of $j_z \approx 200$ kA/m$^2$. In comparison, the measured current density has an amplitude on the order of $j_z \approx 10$ kA/m$^2$. Thus, the plasma current is clearly limited by some mechanism. The previous estimate is based only on the inductive electric field and does not account for any additional electrostatic field. Since the plasma gun is not simply a source of electrons but instead produces a high density plasma, it is possible that not enough additional electrons are provided to counteract the build-up
of an opposing electrostatic field. As described in section 3.4.1, without an electron source, the current is limited to twice the ion saturation current. For the temperature and density given above, this would yield a current on the order of \( j_z = 2 \cdot j_{i,\text{sat}} \), which is lower than the measured value, indicating that at least some fraction of additional electrons are supplied. The observed current density amplitude is likely due to a combination of the inductive and electrostatic field and is set by the ion saturation current together with an increase due to additional electrons provided by the plasma gun.

Since the ion saturation current is density dependent, one would expect the measured current to increase for higher densities. This can be achieved by operating the rf-antenna in helicon mode. Figure 3.14 shows the current flowing through grid 2 for a high \( n \approx 10^{18} \text{ m}^{-3} \) and a low \( n \approx 10^{16} \text{ m}^{-3} \) density background rf-discharge plasma. Also plotted are the currents without the reconnection drive. For the helicon discharge, there is a positive current flowing before the plasma gun discharge at \( t = 3.7 \mu\text{s} \). This is likely to be an ion saturation current due to the high densities achieved in a helicon discharge and the inductive electric field set up by the small time-variation of the current flowing through the conductors creating the X-point topology. Comparing the two cases in figure 3.14, it is clear that the current is increased significantly by an increase in the background density. However, the increase is mainly for the background current, and leads to a correspondingly smaller modulation degree of the plasma current due to the reconnection drive. Furthermore, the background current has a non-negligible time-dependence for the higher density measurement, peaking at around \( t = 60 \mu\text{s} \).
increase for the background current than the inductively driven current can be explained by the different time-dependence of the two: The amplitude of the plasma current does not only depend on the amplitude of the inductive electric field but also on the frequency, with a higher frequency yielding a smaller current due to a higher reactance.

The inductively extracted current is the most important in terms of the magnetic reconnection, and it is beneficial to keep the density of the rf-plasma low in order to reduce the influence of the background current. However, the rf-plasma improves the reproducibility of the plasma gun discharge and all the measurements presented in this thesis were carried out with a low density background plasma ($n \approx 10^{15-16} \text{ m}^{-3}$).

### 3.5 Reconnection regimes

As outlined in section 2.5, magnetic reconnection can be divided into different regimes based on the amplitude of the Lundquist number relative to the effective plasma size $\lambda$. Figure 3.15 shows a phase diagram similar to figure 2.6 for VINETA II based on esti-

<table>
<thead>
<tr>
<th></th>
<th>Case I</th>
<th>Case II</th>
<th>Case III</th>
</tr>
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<tbody>
<tr>
<td>Plasma density $n$</td>
<td>$10^{19}$ m$^{-3}$</td>
<td>$10^{15}$ m$^{-3}$</td>
<td>$10^{18}$ m$^{-3}$</td>
</tr>
<tr>
<td>Electron temperature $T_e$</td>
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<td>6 eV</td>
<td>4 eV</td>
</tr>
<tr>
<td>Magnetic guide field $B_{\text{guide}}$</td>
<td>100 mT</td>
<td>6 mT</td>
<td>6 - 100 mT</td>
</tr>
<tr>
<td>Characteristic length scale $L$</td>
<td>10 cm</td>
<td>10 cm</td>
<td>5 cm</td>
</tr>
<tr>
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<td>$&lt; 2$ mT</td>
<td>$&lt; 2$ mT</td>
</tr>
<tr>
<td>Drive time $\tau_{\text{drive}}$</td>
<td>10 $\mu$s</td>
<td>10 $\mu$s</td>
<td>10 $\mu$s</td>
</tr>
<tr>
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<td>30 cm</td>
<td>0.6 cm</td>
</tr>
<tr>
<td>Ion gyroradius $r_{\text{i}, i}$</td>
<td>0.3 cm</td>
<td>5 cm</td>
<td>0.3 - 5 cm</td>
</tr>
<tr>
<td>Electron gyroradius $r_{\text{L,e}}$</td>
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<td>0.1 cm</td>
<td>$7 \cdot 10^{-3}$ - 0.1 cm</td>
</tr>
<tr>
<td>Lundquist number $S$</td>
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<td>130</td>
<td>1.5</td>
</tr>
<tr>
<td>Diffusion time $\tau_D$</td>
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<td>6 $\mu$s</td>
</tr>
<tr>
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<td>4 cm</td>
<td>4.8 cm</td>
</tr>
<tr>
<td>Ion sound gyroradius $\rho_s$</td>
<td>1 cm</td>
<td>26 cm</td>
<td>1 - 22 cm</td>
</tr>
<tr>
<td>Electron magnetization: $r_{\text{L,e}}/L$</td>
<td>$1.4 \cdot 10^{-2}$</td>
<td>$5 \cdot 10^{-4}$</td>
<td>$1.3 \cdot 10^{-3}$ - $2.3 \cdot 10^{-2}$</td>
</tr>
<tr>
<td>$\omega_{ce}/\nu_{\text{coll}}$</td>
<td>3.7</td>
<td>214</td>
<td>5 - 84</td>
</tr>
<tr>
<td>Ion magnetization: $r_{\text{L,i}}/L$</td>
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<td>0.5</td>
<td>$5.8 \cdot 10^{-2}$ - $2 \cdot 10^0$</td>
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<tr>
<td>$\omega_{ci}/\nu_{\text{coll}}$</td>
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<td>0.1</td>
<td>$2.7 \cdot 10^{-3}$ - $4.6 \cdot 10^{-2}$</td>
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</table>
Figure 3.15: Phase diagram for VINETA II. The black arrows indicate the changes (increase/decrease) required for the specific quantities for the transition (green arrows) between the collisional and collisionless branch divided by the dashed line. The red box indicates the parameter region based on the rf-plasma parameters and the blue line the region based on the gun plasma parameters.

mates of the plasma and reconnection parameters for three different scenarios as shown in table 3.2: (I) the high density helicon discharge, (II) the low density and higher temperature inductively generated plasma, and (III) the gun plasma for a high and low guide field. The lower limit is given by the minimum guide field required to reliably sustain the rf-plasma and plasma gun discharge. The high and low density cases of the rf-generated plasma gives the two extremes of the operation parameters in the experiment. The red box in figure 3.15 indicates the parameter region accessible based on these parameters. The blue line represents the gun plasma. The collisional/collisionless regimes are divided by a dashed line which is given by \( S = \frac{\lambda^2}{4} \). Also indicated in the figure are the changes (increase/decrease) in quantities such as temperature, density and magnetic field required to transition between the collisional and collisionless regimes. Figure 3.15 and the large ion sound gyroradius at small guide fields implies that, though the plasma gun restricts the parameter space, the investigation of both single X-line collisional and collisionless reconnection should be achievable. However, the characteristic length scale \( L \) of the gun plasma, which is smaller than for the other two scenarios, is now on the order of the resistive layer width \( \Delta_{\text{res}} \), and the diffusion time \( \tau_D \) is shorter than the drive
time $\tau_{\text{drive}}$. This implies that reconnection in the device should be firmly in the collisional regime. This is further supported by the low Lundquist number ($S < 1$) and an electron mean free path which is smaller than the characteristic length scale. The parameters in table 3.2 also indicate that while the electrons are magnetized in all three operation scenarios, the ions are not, i.e., they do not meet the requirement that the characteristic length scale of the plasma is much larger than the gyroradius of the particle ($r_L/L << 1$), and that they on average complete at least one gyration before colliding ($\omega_{ci}/\nu_{\text{coll}} > 1$).
4.1 Plasma diagnostics

4.1.1 Langmuir probes

A Langmuir probe is a diagnostics tool for the local measurement of plasma parameters such as the plasma density \( n \), the plasma potential \( \phi_p \), and the electron temperature \( T_e \) [93]. In its simplest form, a Langmuir probe consists of an insulated wire with a small uncovered tip which serves as an electrode. Since only the exposed part of the probe is in electrical contact with the plasma, measurements can have high spatial resolution. The probe tip is biased with a variable voltage \( U_b \) and the resulting current \( I_p \), which is due to the flux of electrons and ions to the probe, is measured by the voltage drop over a shunt resistor. The current-voltage \((I_p - U_b)\) characteristic is then evaluated to determine the above mentioned plasma parameters. A typical probe characteristic is shown in figure 4.1. It can be divided into three regions: (I) the ion saturation region, (II) the transition region, (III) and the electron saturation region. The regions are roughly separated by the floating potential \( \phi_f \) and the plasma potential \( \phi_p \). At bias far below the floating potential \((U_b << \phi_f)\), the retarding potential of the probe repels all electrons and the current is only from the ion flux. Similarly, for a probe bias above the plasma
potential ($U_b \gg \phi_p$) ions are repelled and only electrons are collected by the probe, resulting in an electron saturation current. For probes with cylindrical geometry, which were used in this thesis, the current does not saturate in the two regions, as can be seen in figure 4.1. This is due to sheath expansion related to the increasing probe voltage [94]. In the transition region, the net current is given by both the electron and ion flux, with the contribution of the electrons increasing with increasing bias and determined by the electron energy distribution function (EEDF). Assuming a Maxwellian EEDF,

$$f(E) = n \left( \frac{m}{2\pi k_B T_e} \right)^{3/2} \exp \left( \frac{-E}{k_B T_e} \right),$$

(4.1)

the electron current in the transition region is exponential.

There exists a vast number of different probe theories for evaluating the probe characteristic [95,96]. For the case of a cylindrical probe in an unmagnetized collisionless plasma in thermodynamic equilibrium, the original Langmuir theory has been developed [97]. For $U_b << \phi_f$ the flux of ions to the probe is determined by the Bohm-criterion [85], i.e., the ions are accelerated in the electric field of the presheath and enter the sheath with ion sound speed $c_s$. The ion saturation current then reads

$$I_{i,\text{sat}} = \exp(1/2) A c_e n e = \exp(1/2) A n e \sqrt{\frac{k_B T_e}{m_i}},$$

(4.2)
4.1 Plasma diagnostics

where $A$ is the (effective) surface area of the probe tip, $m_i$ is the ion mass and the factor $\exp(1/2)$ accounts for density depletion in the presheath. The electron saturation current is given by

$$I_{e,\text{sat}} = -Ane\bar{v}/4 = -Ane\sqrt{\frac{k_B T_e}{2\pi m_i}},$$  \hspace{1cm} (4.3)

where $\bar{v}$ is the mean electron velocity and $m_e$ the electron mass. Assuming a Maxwellian EEDF, the electron current in the transition region is given by

$$I_e = I_{e,\text{sat}} \exp\left(\frac{e(U_b - \phi_p)}{k_B T_e}\right).$$  \hspace{1cm} (4.4)

At the floating potential $\phi_f$, the electron and ion currents cancel out and no net current is drawn. It is negative with respect to the plasma potential, and equating (4.2) and (4.4) yields

$$\phi_f = \phi_p - \frac{k_B T_e}{2e} \ln \frac{2m_i}{\pi m_e} = \phi_p - a \frac{k_B T_e}{e},$$  \hspace{1cm} (4.5)

where $a \approx 5.2$ for Argon. The plasma potential is usually obtained from the minimum of the first derivative of the current, i.e., the turning point of the characteristic. The electron temperature can be found from an exponential fit to the electron current in the transition region. Having the electron temperature, using equation (4.2) yields the density

$$n = \frac{I_{i,\text{sat}}}{Ae \exp(1/2) \sqrt{\frac{m_i}{k_B T_e}}.}$$  \hspace{1cm} (4.6)

Depending on the plasma environment, the simple description given above is often not sufficient for evaluating the plasma parameters. Instead more complex probe theories or modified probe designs are required. For example, for a magnetized plasma, the charged particle motion is anisotropic since the particles can only move freely along the magnetic field. This significantly affects the particle flux to the probe surface [98] and thus the interpretation of probe characteristic. The plasma potential fluctuations generated in radio frequency (rf) produced plasmas can distort the averaged probe characteristic [99] which leads to a systematic overestimation of the electron temperature. This can be counteracted by using a probe design that compensates for these fluctuations [100, 101].

The necessity of sweeping the bias voltage and measuring the entire probe characteristic, limits the time-resolution of the Langmuir probe measurement. This means that it is not possible to use this method to obtain the plasma parameters during the plasma gun dis-
charge and the reconnection drive phase. Assuming a constant electron temperature, a single Langmuir probe can be used to estimate a change in density over time by measuring the fluctuations of the ion saturation current and using equation (4.2). Unfortunately the PFNs and the reconnection drive produce such a high electrical noise that it is impossible to use this method for measuring the density fluctuations during the reconnection drive phase. This noise is seen on all diagnostics, as well as on the ground line, as any unshielded cable will act as an antenna and pick it up. It consist of: (a) a fast component in the MHz range, which is produced by the switching of the PFNs and the IGBTs, and (b) a component proportional to the time-derivative of the current produced by the PFNs and the reconnection drive.

### 4.1.2 Emissive probe

An emissive probe allows for direct measurements of the plasma potential without having to determine the full probe characteristic [102]. The basic principle is to heat an electrode to emit electrons in order to neutralize the potential drop within the surrounding Debye sheet and cause the probe to float at the plasma potential. For a Maxwellian plasma and neglecting space charge effects, the plasma potential can be described by [103]

\[
\phi_p = \phi_f + \frac{k_B}{e} T_e \ln \left( \frac{j_{e,\text{sat}}}{j_{\text{em}} + j_{i,\text{sat}}} \right),
\]

(4.7)

where \( j_{e,\text{sat}} \) is the electron saturation current density, \( j_{i,\text{sat}} \) is the ion saturation current density and \( j_{\text{em}} \) is the emission current density. To keep the probe floating at the plasma potential requires an emission current of equal amplitude to the collected electron saturation current (neglecting the small ion saturation current). A sketch of the current-voltage characteristic for an emissive probe is shown in figure 4.2. It can be divided into collecting and emitting regimes, where the collecting regime is identical to an unheated Langmuir probe. The emission current is zero above the plasma potential, or more exactly at a bias of a few \( T_w/e \approx 0.2 \text{ V} \) (where \( T_w \) is the wire temperature) above the plasma potential due to the electrons in the high-energy tail of the EEDF overcoming the potential barrier [102]. Below the plasma potential, the emission current is constant and limited by the temperature of the wire.

This simplified picture of the emission neglects space charge effects which are always
Figure 4.2: Schematic diagram of the current-voltage characteristic for an emissive probe and its collecting and emitting components. The plasma potential is chosen to be zero. $U_b$ is the probe bias, $I_p$ is the current out of the probe, $I_e^*$ is the electron saturation current and $I_{e0}$ is the temperature limited electron emission current. Taken from ref. [102].

Present and can significantly alter the structure of the surrounding sheath [105]. As the temperature of the electrode is increased, the sheath potential decreases and the floating potential of the probe approaches the plasma potential until the emission reaches a point where the electric field at the surface is zero [106]. Further heating results in the formation of a virtual cathode, preventing additional electrons escaping into the plasma and hence limiting the emission. The consequence is that the floating potential of an emissive probe always underestimates the plasma potential by approximately $T_e/e$ [105]. Figure 4.3 shows the current-voltage characteristics of an emissive probe for different heating powers. As the heating increases, the floating potential approaches the plasma potential ($\phi_f = 10.6$ V) until it saturates at high emission.

One of the main advantages of the emissive probe over the standard Langmuir probe is that it can be used for measurements in plasmas with energetic electrons, e.g., beam plasmas or double plasma devices [107, 108]. There are a number of different measurement techniques for obtaining the plasma potential with an emissive probe. For this thesis, the “floating point with large emission method” was used [102]. This method assumes that the floating potential at high emission is very close to the plasma potential and is very useful for obtaining time-resolved measurements at timescales too fast for sweeping of the probe bias voltage. Measurements and numerical simulations have shown that the
floating potential in the limit of large emission can be described by [106]

\[
\phi_f = \phi_p - \alpha T_e,
\] (4.8)

where \(\alpha \approx 1-2\). This expression is similar to equation (4.5), though the origin of the constant \(\alpha\) differs from \(\alpha\). Combining the two equations and using measurements of the floating potential of a hot and cold probe, the temperature can be determined through

\[
T_e = e^{\frac{\phi_{f,h} - \phi_{f,c}}{\alpha - \alpha}},
\] (4.9)

with \(\phi_{f,h}\) the floating potential of the hot probe and \(\phi_{f,c}\) the floating potential of the cold probe. This assumes the absence of fast electrons, which would significantly alter the floating potential of the cold probe.

The heating of the probe can be achieved in several ways, e.g., through Joule heating [109–111], self-emission [112, 113], and laser heating [114, 115]. Figure 4.4 shows the emissive probe setup used for the measurements in this thesis. The probe consists of a small tungsten wire loop which is heated by a DC current. The current is supplied by a battery together with a small electric circuit that allows for the current to be varied. The probe is electrostatically shielded up to the probe tip in order to reduce the antenna pick-up.
4.2 Current and magnetic fluctuation diagnostics

In order to study magnetic reconnection, it is essential to obtain a detailed description of the electromagnetic fields and currents that, along with the evolution of plasma parameters such as temperature and density, characterize the system. The total plasma current is measured using Rogowski coils that encompass the whole plasma cross section. The time-dependent magnetic fields are measured using $\dot{B}$-probes, and yield a number quantities of interest. From the magnetic field $\mathbf{B}$, the vector potential $\mathbf{A}$ can be obtained via $\mathbf{B} = \nabla \times \mathbf{A}$. The axial component $A_z$ allows for evaluating the in-plane magnetic field lines which are represented by the contours of constant vector potential. The magnetic flux of the conductors can be subdivided into the common flux of the field lines encircling both conductors, and the private flux encircling the individual conductors. The magnetic flux through a surface $S$ is given by $\Phi = \int_S \mathbf{B} \, d\mathbf{S}$. Since the magnetic field is almost independent of the axial coordinate $z$, this can be simplified to the flux per axial length $\Phi/z = \int (-B_x dy + B_y dx)$. The magnetic field is zero at the wall due to the mirror currents and hence the common flux is given by $\Phi_c/z = \int_0^R B_y dx$ and the private flux by $\Phi_p/z = \int_0^{d/2} B_x dy$, where $R$ is the radius of the vacuum chamber and $d$ is the separation of the conductors. The common magnetic flux is equivalent to the $z$-component of the vector potential at the X-point $A_z(0,0)$. Since the induced axial electric field can be derived from $\mathbf{E}_z = -\partial \mathbf{A}_z/\partial t$, $\mathbf{E}_z(0,0)$ gives the rate of change of the flux. The measurement of the in-plane magnetic field can also be used to determine the current density through $\nabla \times \mathbf{B} = \mu_0 \mathbf{J}$.
Figure 4.5: (a) Schematic of a Rogowski coil. (b) Calibration signals of Rogowski coil. The measured voltage output of the Rogowski coil is given by the green line. The red dotted line is the current obtained by integrating the voltage signal and scaling it to the known current (blue line).

4.2.1 Rogowski coil

A Rogowski coil is a simple diagnostic for measuring alternating currents [116, 117]. A schematic of a Rogowski coil is shown in figure 4.5a. It consists of a solenoidal coil with its ends brought together. To avoid a contribution from magnetic flux through the Rogowski coil, one of the signal wires returns through the solenoidal coil. The Rogowski coil has the useful trait compared to other current transformer that the output of the coil does not saturate for large currents [93]. The voltage output of the coil is proportional to the rate of change of current through the coil and is given by

$$U_{Rog} = \frac{\mu_0 NA}{l} \frac{dI}{dt} \quad (4.10)$$

where $N$ is the number of windings, $A = \pi a^2$ is the area of one winding with radius $a$, and $l = 2\pi R$ is the circumference of the Rogowski coil with radius $R$. It is assumed that the windings are evenly spaced and small compared to the radius of the coil. The current is determined by integrating the measured voltage. Calibration of the Rogowski coil, and obtaining the integration constant, is done by measuring the voltage output of a known current. A calibration example is shown in figure 4.5b, where a calibration constant of $2.3 \cdot 10^7$ is obtained.
4.2 Current and magnetic fluctuation diagnostics

Figure 4.6: (a) Schematic of a $\dot{B}$-probe. The measured output voltage is a combination of the induced voltage in the loop and the voltage due to the coupling to electrostatic potential fluctuations of the plasma. (b) Setup using a Helmholtz coil arrangement (red) for the absolute calibration of the $\dot{B}$-probe.

4.2.2 Magnetic fluctuation probes

Induction coils, or $\dot{B}$-probes, are widely used for measurements of magnetic fluctuations [93, 116] and can be inserted directly into the plasma to give highly localized measurements [118]. A schematic of a $\dot{B}$-probe is shown in figure 4.6a. The principle of the probe is based on Faraday’s law

$$\oint_C \mathbf{E} \cdot d\mathbf{l} = - \int_A \dot{\mathbf{B}} \cdot dA,$$  \hspace{1cm} (4.11)

and for $N$ number of windings with contour $C$ enclosing an area $A$, the induced voltage is

$$U_{\text{ind}} = -NA\dot{B}(t).$$  \hspace{1cm} (4.12)

For sinusoidal fluctuations, the induced voltage is given by

$$U_{\text{ind}} = -NAB\omega \cos(\omega t)$$  \hspace{1cm} (4.13)

Hence, the sensitivity of the probe scales linearly with the area, the number of windings and the frequency. While it seems beneficial to simply increase $N$ and $A$ in order to improve the sensitivity, a number of different factors have to be taken into account when designing a $\dot{B}$-probe. First, the spatial resolution of the probe and its disturbance of
the plasma is dependent on the size of the probe, and the area is preferably kept as small as possible. Increasing the number of windings is a feasible way to increase the sensitivity. However, the maximum measurable frequency of the coil is of the order of \( f = \frac{R_0}{L} \) [119], where \( L \) is the inductance of the coil and \( R_0 \) the resistance terminating its output. Since the inductance of the coil is proportional to its radius, \( r \), and to the square of the number of windings \( L \propto rN^2 \), the upper limit of the frequency response decreases quadratically with the number of windings. When designing the probe, one also has to take into account the resonance frequency of the coil \( f_{\text{res}} = \frac{1}{\sqrt{LC}} \).

A crucial consideration in determining the required sensitivity is the amount of noise picked up by the probe. A well-known problem with \( \dot{B} \)-probes is their capacitive coupling to electrostatic potential fluctuations [119]. Reduction of this capacitive pick-up is especially necessary in rf-generated plasmas where large potential fluctuations can be expected. A highly effective capacitive pick-up reduction is achieved by a differential measurement of the voltage of the probe [120]. Since rotating the induction coil by 180° will result in the inductive signal changing sign while the electrostatic pickup remains unchanged, the electrostatic pick-up is thus eliminated.

**Design and data evaluation**

A sample probe designed for measuring the reconnection magnetic field is shown in figure 4.7. The probe consists of three perpendicular induction coils which are used for measuring the magnetic field in the \( x \)-, \( y \)- and \( z \)-direction. The coils are made using Kapton-coated wire with a diameter 0.1 mm and the probes are coated with epoxy in order to provide further insulation from the plasma. In order to minimize noise pick-up, the signal wires of the probe are twisted and electrostatically shielded. The coils measuring
4.2 Current and magnetic fluctuation diagnostics

Figure 4.8: (a) Frequency dependence of the sensitivity of the probe (blue line). The frequency dependence of the integrated signal is also shown (green line). (b) Phase shift dependence on frequency before (blue line) and after integration (green line)

The $x$- and $y$-components of the magnetic field, each consist of $N = 20$ windings with an area of approximately $A = 2 \text{ cm}^2$ ($1 \times 2 \text{ cm}$). For the $z$-direction, $N = 20$ windings with an area of $A = 1 \text{ cm}^2$ ($1 \times 1 \text{ cm}$) is used. Hence, the probes have a spatial resolution of $1 \text{ cm}$ in the $xy$-plane. It is assumed that the reconnecting magnetic field is mainly in the azimuthal plane and only varies slightly along the $z$-axis. In order to improve the sensitivity, the probes were made larger in the $z$-direction, thereby increasing the area of the $x$- and $y$-direction coils.

The sensitivity calibration of the $\hat{B}$-probe is done using the nearly uniform magnetic field of a Helmholtz coil arrangement which is given by $B_{\text{H}} = (4/5)(3/2)(\mu_0 N I / R)$. The schematic of the Helmholtz coils setup is shown in figure 4.6b. The coil pair has a separation and radius of $3.5 \text{ cm}$ and consists of four windings each. A voltage with a frequency range of $f = 10 \text{ kHz} - 10 \text{ MHz}$ drives a current through the coils. The amplitude of the current is obtained either with a current monitor or from the voltage drop over a $50 \Omega$ resistor connected in series with the coils. The measured frequency dependence of the sensitivity of a $\hat{B}$-probe is shown in figure 4.8a. For frequencies below $f = 1 \text{ MHz}$ the sensitivity increases linearly and has a maximum at the resonance frequency of $f_{\text{res}} = 4.5 \text{ MHz}$. For measuring incoherent signals of the fluctuating magnetic field, it is necessary to first integrate the signal to obtain a calibration constant. The integrated signal is also shown in figure 4.8a. It is constant up to $f = 1 \text{ MHz}$, which is well above the frequency of the reconnection drive. For measurements of fluctuating magnetic fields with
a frequency below $f = 1 \text{ MHz}$, the magnetic field is simply determined by multiplying this constant with the integrated voltage obtained in the measurement. Figure 4.8b shows the phase shift dependence on the frequency. A nearly constant phase shift of $\pi/2$ with respect to the current can be seen for frequencies well below the resonance frequency. For the integrated signal, there is no phase shift below the resonance frequency.

A test of the capacitive pick-up of the probe was done using a parallel-plate capacitor. Electric fields of up to 5 V/m are generated inside the capacitor and the frequency response is shown in figure 4.9. As a reference, the voltage response to a magnetic field of $B=0.1 \text{ mT}$ is shown. The pick-up is smaller than the noise level of the data acquisition system for frequencies below the resonance frequency, and orders of magnitude smaller than the induced voltage.

**$\dot{B}$-probe measurement of magnetic reconnection**

The spatiotemporal evolution of the reconnection magnetic field is obtained by scanning a $\dot{B}$-probe on a shot-by-shot basis through the entire plasma cross section using the positioning system. Figure 4.10a shows the calculated modulus of the magnetic field in the azimuthal plane for a coil current of $I=1 \text{ kA}$. The magnitude of the measured
4.2 Current and magnetic fluctuation diagnostics

Figure 4.10: (a) Modulus of the magnetic field in the azimuthal plane calculated for a coil current of $I=1\,\text{kA}$, with contours of constant vector potential representing magnetic field lines. The black rectangle indicates the measurement region. (b) Measurement with a $\dot{B}$-probe in vacuum. (c) One-dimensional cuts of the measured (dotted blue line) and calculated (green line) magnetic field.

The magnetic field in vacuum is shown in figure 4.10b. Radial cuts through the X-point of the measured and the calculated magnetic field are plotted in figure 4.10c. There is a good agreement close to the X-point and the discrepancy further out is due to mirror currents in the vessel wall, which has not been taken into account in the calculated value for the magnetic field.

### 4.2.3 Loop voltage probe

In order to obtain the amplitude of the vector potential (and thereby the inductive electric field) from the magnetic field measurements, it is necessary to start the spatial integration where the magnetic field is zero, i.e., at the wall. Since the measurement region of the positioning system lies far from the wall, an offset of the vector potential is introduced when evaluating the measurement. To correct for this offset, a loop voltage probe is used to measure the inductive electric field at the starting point of the integration. The probe consist of a radially inserted square loop with one side with length $\Delta z$ parallel to the $z$-axis. As a result, the open-loop voltage $U$, measured outside the vacuum chamber wall, stems only from the induced voltage of the section in the $z$-direction. The electric field is then given by $E_z = -\dot{A}_z = U/\Delta z$ and the offset of the vector potential can be found...
by time-integration of the electric field. Figure 4.11 shows the vector potential obtained from the loop probe and the offset-corrected vector potential at the X-point.

The loop probe is also useful for direct measurements of the inductive electric field at the X-point, and thus the rate of change of the magnetic flux. This was done using a loop probe setup consisting of three loop probes with lengths of $\Delta z = 13.7$ cm placed next to each other at three axial positions. This allows for determining if the electric field has an axial gradient. The probe setup is placed approximately at the center of the large vacuum chamber such that the middle of the three loop probes corresponds to the inductive field at the axial position of the positioning system.
Guide field reconnection in the VINETA II device

5.1 Localized plasma source

Before introducing an X-point topology and driving reconnection, it is useful to first characterize the plasma and background current produced by the plasma gun. While the issues associated with Langmuir probe measurement, as outlined in section 4.1, makes it challenging to reliably determine the plasma parameters, estimates can be obtained using the magnetic field diagnostics along with emissive probe measurements of the plasma potential.

5.1.1 Magnetic field and current sheet geometry of background current

Figure 5.1a shows the azimuthal profile of the axial background current density $j_z$ as produced by the plasma gun, and figure 5.1b the associated in-plane magnetic field. The axial current density can be approximated by a two-dimensional Gaussian distribution
whose full width at half maximum (FWHM) is mainly set by the size of the opening of the plasma gun and the diffusion of the plasma across the guide field ($B_g = 18$ mT in the current case). The in-plane current density $j_{xy}$ is shown in figure 5.1c and the associated axial magnetic field is plotted in figure 5.1d. The current circulates in the azimuthal plane and can be assumed to be the diamagnetic current created by steep radial plasma pressure gradients. The in-plane current is discussed in more detail in section 5.6.2.

### 5.1.2 Plasma potential and temperature

The temporal evolution of the plasma potential at the position of the current density peak, obtained using an emissive probe, is plotted in figure 5.2a. Also plotted is the floating potential measured with the same probe without any heating. The plasma current is plotted for reference. The plasma potential initially has positive value of $\phi_p = 21.5$ V
5.1 Localized plasma source

Figure 5.2: (a) Time-evolution of the plasma potential (red line) and the floating potential (blue line) at the position of the current density peak. The plasma current is plotted for reference (green line). (b) Azimuthal plasma potential profile averaged over the time interval $\Delta t_2$ as indicated in (a). The contour level at twenty percent of the peak current density is plotted in black. (c) One-dimensional cuts through the center point of the plasma potential (red line) and the floating potential (blue line) averaged over the time interval $\Delta t_1$ (dotted lines) and $\Delta t_2$ (solid lines). A one-dimensional cut of the current density profile is plotted for reference (green line).

in the rf-plasma and then rapidly drops to negative values as the plasma gun discharge starts. After an initial peak, the plasma potential settles on a fairly constant value of $\phi_p \approx -15$ V. The negative value of the plasma potential is due to a surplus of electrons which results in negative space-charge. This could either come from a greater loss of ions than electrons from the plasma due to confinement of the electrons, or as a result of the electrons injected by the plasma gun (or a combination of both).

Figure 5.2b shows the azimuthal profile of the plasma potential averaged over the time interval $\Delta t_2$ as indicated by the shaded gray region in figure 5.2a. The contour level of the axial current density at 20% of the peak value is plotted for reference (cf. figure 5.1a). The plasma potential has a peak at the same position as the axial current...
density. It is negative at this position and is positive at the edge of the measurement region. This is further highlighted in figure 5.2c, where one-dimensional cuts through the peak of the plasma potential are plotted for two different time-intervals. For reference, the current density profile averaged over $\Delta t_2$ is also plotted in the diagram, along with one-dimensional cuts of the floating potential averaged over the two time-intervals $\Delta t_1$ and $\Delta t_2$. The plasma potential of the rf-produced plasma has a flat profile with $\phi_p \approx 20$ V and has a peak value of $\phi_p = -14$ V during the plasma gun discharge. The floating potential of the rf-discharge has a flat profile with a small dip at the center. Similar to the plasma potential, at $\Delta t_2$ the floating potential becomes negative, and has a peak value of $\phi_f = -30$ V.

The profile of the plasma potential results in an in-plane electrostatic field. Figure 5.3a shows the electrostatic field in the azimuthal plane as obtained from the gradient of the plasma potential, and figure 5.3b a one-dimensional cut through the center. The electrostatic field points radially inwards and has a hollow profile with a local minimum at the center and peaks along the edge of the current sheet.

An estimate of the temperature can be obtained from equation (4.8) via the measurements of the plasma and the floating potential. The so-obtained temporal evolution of the temperature at the position of the current density peak is plotted in figure 5.4a. It has an initial value of $T_e \approx 6$ eV in the rf-generated plasma and then, after a minimum of
5.1 Localized plasma source

Figure 5.4: (a) Time-evolution of the electron temperature $T_e$ at the position of the current density peak (black line). The plasma current is plotted for reference (green line). (b) Temperature profile in the azimuthal plane averaged over the time-interval $\Delta t_2$. The contour level at twenty percent of the peak current density is also plotted (black line). (c) One-dimensional cut through the center of the temperature (solid black line) averaged over the time interval $\Delta t_1$ (dotted black line) and $\Delta t_2$ (solid line). A one-dimensional cut of the current density profile is plotted for reference (green line).

$T_e = 2.4 \text{ eV}$ as the gun first discharges, settles at around $T_e \approx 3.6 \text{ eV}$. The time-averaged azimuthal temperature profile of the gun plasma is plotted in figure 5.4b. The radial profiles of the rf- and gun-plasma are plotted in figure 5.4c. The temperature of the rf-plasma has a maximum of $T_e \approx 6 \text{ eV}$ at the radial position of the plasma gun and then decreases to approximately 5 eV at the edge of the measurement region. The amplitude is consistent with measurements done with Langmuir probes in similar rf-plasma discharges [77]. The temperature of the gun plasma has a maximum at the same position as the current density peak and decreases towards the edge before again rising slightly. It should be kept in mind that the azimuthal profiles of the plasma potential and floating potential are not obtained simultaneously and the temperature does not account for changes in the plasma parameters between the measurement runs. This method for measuring the temperature also assumes that fast electrons have no influence.
5 Guide field reconnection in the VINETA II device

Figure 5.5: Azimuthal profiles of: (a) the in-plane $j \times B$-force obtained from measurements, (b) the $\nabla p$-force derived from estimates of the density profile and the measured temperature profile, and (c) the residual of the two forces. The arrows indicate the direction of the force.

5.1.3 MHD force balance and plasma density

From the assumption of ideal MHD force equilibrium $j \times B = \nabla p$, the $j \times B$-force yields a measure of the pressure gradient force. The plasma pressure profile can be obtained using the temperature profile in figure 5.4b and by assuming a density profile that follows the current density profile in figure 5.1a. The plasma density amplitude is scaled such that the resulting $\nabla p$-force balances the $j \times B$-force as well as possible. The $j \times B$-force is shown in figure 5.5a, the estimate of the $\nabla p$-force in figure 5.5b, and their difference in figure 5.5c. There is a reasonably good match between the two forces. They both have a minimum at the center, a ring shaped region where the force is largest and fall off towards zero outside the current sheet. The $j \times B$-force, however, has a flatter distribution and hence the residual is biggest outside the current sheet. Inside the current sheet the residual has a mean value of 2.4 Pa/m. For the $\nabla p$-force in figure 5.5b, a peak density amplitude of $n_e = 7 \cdot 10^{17} \text{ m}^{-3}$ is obtained. This is a reasonable value since triple probe measurements indicate that the peak density is on the order of $n_e \approx 10^{18} \text{ m}^{-3}$ [91].
5.2 Guide field dependence of current sheet geometry

5.2.1 Magnetic mapping

In order to better understand the current sheet geometry during reconnection, it is useful to first determine how the in-plane X-point magnetic field topology affects the current sheet (without any reconnection drive). As described in section 2.4.1, magnetic mapping will lead to an elongation of the current sheet along the separatrix in one direction and a squashing in the other direction. The degree to which the current sheet is deformed depends on the relative amplitude of the guide field and the in-plane magnetic field $B_g/B_{xy}$. The pitch angle of the field lines is shown in figure 5.6a for $B_g = 18$ mT and a magnetic field topology calculated for a 1.8 kA current through the conductors. The angle is zero at the X-point and largest near the conductors. The separatrix is plotted as black arrows for reference. Figure 5.6b shows a circular current sheet with a diameter of 1.2 cm (based on the aperture of the gun anode), mapped axially over a distance of $z = 53$ cm for two different guide field strengths ($B_g = 18$ mT and 90 mT). The area enclosed by the initial circle is conserved, but the shape is deformed into an ellipse. The orientation of the elongation depends on the direction of the magnetic field lines of the in-plane field. For a current flowing in the opposite direction through the conductors, the sheet would be rotated by $90^\circ$. For $B_g = 90$ mT, the current sheet width and length, taken as the minor and major axis of the ellipse (1 cm and 1.4 cm, respectively), are similar to the gun aperture. However, for the lower guide field case, the in-plane field has a signifi-

![Figure 5.6: (a) Pitch angle of the magnetic field line for a guide field of $B_g = 18$ mT. (b) Circular current sheet (black) mapped axially over a distance of $z = 53$ cm for $B_g = 18$ mT (blue) and $B_g = 90$ mT (red). (c) Same as for (b) but with a current sheet shifted with respect to the X-point.](image-url)
significant effect on the sheet, with pronounced elongation along the separatrix. The resulting ellipse has a minor axis of 0.5 cm and a major axis of 3.1 cm. If $B_g/B_{xy}$ is sufficiently small (or the axial distance is sufficiently large), the current sheet can elongate around the conductors, taking on a figure-eight shape. Figure 5.6c shows mapping examples where the initial circular current sheet is shifted with respect to the X-point and centered on $(x,y) = (0.5, 0.5)$ cm. The magnetic mapping moves the current sheet towards the separatrix. The center of the resulting ellipses are shifted to $(x,y) \approx (0.5, 0.4)$ cm for $B_g = 90$ mT and to $(x,y) \approx (0.9, 0.2)$ cm for $B_g = 18$ mT.

### 5.2.2 Current sheet geometry

In order to investigate how the current sheet evolves along the $z$-axis, a probe setup consisting of three $\dot{B}$-probes placed at three axial position was used. The probes have a spacing of $\Delta z = 14$ cm, with the middle probe placed at a distance of $z = 53$ cm from the gun. Figure 5.7 shows the current density $j_z$ at the three consecutive planes for different values of the guide field $B_g$. The measurements are done using setup B, but without the inductive field of the reconnection drive. The contour level corresponding to 50% of the peak value is also plotted. The field lines at the positions of an elliptical fit to this contour level are mapped from one plane to the other in the same way as the examples given in figure 5.6. The results of the magnetic mapping are also plotted in figure 5.7. The current sheet can be seen to be elongated along the separatrix in one direction and compressed in the other. This elongation is more pronounced the greater the distance between the measuring plane and the gun is, as well as the lower the guide field is. The gun is not placed exactly at the X-point, but rather at $(x,y) = (0.2, 0.5)$ cm. Consequently the elongation is not symmetrical along the separatrix. Furthermore, the X-point is shifted by the localized current density. This is especially the case at high guide field measurements, and in some cases the X-point is even shifted to the edge of the current sheet. The peak value of the current density also decreases with the distance from the plasma gun and with decreasing guide field. The dependence of the current sheet shape on the guide field and the distance from the starting plane is a strong indicator that the mapping of the field lines has a significant impact on the current sheet geometry. However, the change in the peak amplitude of the current density implies that magnetic mapping is not the sole contributor to the spatial evolution of the current sheet. For
5.2 Guide field dependence of current sheet geometry

$B_g = 6 \text{ mT}$, the current sheet is elongated enough that it follows the S-like shape of the separatrix and the approximation of the current sheet as an ellipse is no longer viable. Note that, due to the relatively small value of the current density, the current sheet is not well-resolved. The noise-level is obtained by taking the standard deviation of the

![Figure 5.7: Current density $j_z$ at three axial position and for different guide field strengths. The green line indicates the contour at half maximum of the current density peak, and the blue line is the result of magnetic mapping from one plane to the other. The field lines, given by contours of constant vector potential, are plotted in black.](image-url)
Figure 5.8: Fit parameters for two-dimensional Gaussian fits to the current density profiles in figure 5.7 (square markers) and to the profiles of measurements without an X-point topology at \( z = 53 \) cm (black triangular markers). (a) Peak current density \( \hat{j}_z \). (b) Area \( A \) of current sheet at FWHM. (c) Difference in area \( \Delta A \) between the planes at \( z = 53 \) cm and 39 cm (\( \Delta A_1 \)) and the planes at \( z = 67 \) cm and 53 cm (\( \Delta A_2 \)). The difference in area between \( z = 53 \) cm and 39 cm based on the diffusion coefficient is plotted in black.

Current density in a 5×5 cm area in the lower left corner of the \( B_g = 90 \) mT measurement where the current density is expected to be zero. This gives a standard deviation of \( 2\sigma \approx 440 \) kA/m\(^2\). The half maximum current density is \( j_z = 480 \) kA/m\(^2\) at \( z = 67 \) cm and \( B_g = 6 \) mT, which is well in the \( 2\sigma \) range.

Another point of interest in figure 5.7 is that a positive current density can be seen at the edge of the current sheet. This reverse current, i.e., a current directed oppositely to the main current in the sheet, is mainly located along the separatrix in the region where the current sheet has been compressed. However, at \( B_g = 90 \) mT it is observed all around the edge of the sheet. Similar to the main current, the amplitude of the reverse current decreases with the guide field strength and with distance from the gun. However, the reverse current is here at most 10% of the peak amplitude of the current sheet peak value. The origin of the reverse current is discussed in section 5.4.2.

To have a more in-depth look on how the sheet geometry evolves axially and with guide field strength, fit parameters such as the area, angle and minor and major axis at FWHM are obtained from two-dimensional Gaussian fits to the current density profiles. Because of the difficulty with the Gaussian fit at \( B_g = 6 \) mT (due to the shape and amplitude of the current sheet), this case has been omitted in the subsequent figures. Figure 5.8a shows the peak current density of the Gaussian fit \( \hat{j}_z \) at the different planes for the different...
guide field strengths. For comparison, the peak values of the current density at \( z = 53 \) cm for measurements without the X-point topology are also plotted. The amplitude decreases with distance from the gun and increases with guide field strength, with a change of approximately a factor of six between the lowest and highest guide field case (a factor of ten if the \( B_g = 6 \) mT measurement is included). Overall, the amplitude of the current density is larger without an X-point topology. This is likely due to the inductive electric field of the small time-variation of the X-point topology opposing the current.

The value of the areas \( A \) of the cross-section of the current sheets at FWHM are plotted in figure 5.8b. The current sheet area is largest for the lowest guide field, decreases with higher guide field, and eventually starts to plateau. The area expands with axial distance from the gun. This, together with the decrease in amplitude and the fact that this is seen for both with and without an X-point topology, indicates cross-field plasma diffusion: Figure 5.8c shows the difference in area \( \Delta A \) between the three planes and how the area expands with distance from the gun. The difference in area as obtained from cross-field plasma diffusion is also shown. The diffusion coefficient is calculated from estimates of the temperature and density [84], and is highest for \( B_g = 12 \) mT with a value of approximately \( D = 5 \) m\(^2\)/s and decreases to a value of \( D = 0.1 \) m\(^2\)/s for \( B_g = 90 \) mT (with X-point). Correspondingly, the change in the measured area is largest for the low guide field and drops with increasing guide field. The agreement between the diffusion calculation and the measurement suggests that diffusion is the main contributing factor in the expansion of the current sheet area. In summary, the current sheet geometry is set by a combination of magnetic mapping and the plasma diffusion. The discrepancy in figure 5.7 between the current sheet geometry of the measurement and the expectation from magnetic mapping is explained by the plasma diffusion across the magnetic field.

The major and minor axis of the best fit to the measured data are shown in figure 5.9a and 5.9b, respectively. The minor and major axis of the fit to the shape due to the calculated magnetic mapping is also shown. The measured major axis increases with decreasing guide magnetic field and with distance from the gun, as expected from an elongation due magnetic mapping, albeit slightly less than for the calculated values. The minor axis length plotted in figure 5.9b displays a decrease with increasing guide field which is not consistent with magnetic mapping. For a current sheet mapped from the gun, the minor axis length would be expected to be almost unchanged for very high guide field and highly compressed for low guide field. While the calculated minor axis length based
on the mapping of the current sheets in figure 5.7 does exhibit the trend of a small width at low guide fields and a large width at high guide fields, there is a maximum around $B_g = 24 \text{ mT}$. The result of the magnetic mapping is dependent on the starting area, which changes due to cross-field plasma diffusion for different guide field amplitudes and axial starting position. This is the explanation for the observed trend of the minor axis width obtained from the magnetic mapping in figure 5.9b. The discrepancy between the calculated and the observed value, and in particular the decrease of the minor axis with increasing guide field, can be explained by the expansion of the current sheet from cross-field diffusion dominating over the compression due the mapping. As the guide field increases and the cross-section of the current sheet becomes more circular, the minor and major axis approach the same value. This is made evident in figure 5.9c, where the ratio of the minor and major axis is plotted.
5.2 Guide field dependence of current sheet geometry

Another fit parameter is the angle of the current sheet. By magnetic mapping, the angle of a current sheet centered on or close to the X-point should be the angle of the separatrix. The angle of the current sheet for the different guide field amplitudes is shown in figure 5.9d. The observed values are close to the separatrix angle at low guide fields, and around zero at high guide fields where the cross-section of the current sheet becomes circular. The peak value of the angle around $B_g = 24$ mT is attributed to the influence of the reverse current density on the Gaussian fit.

5.2.3 Comparison of high and low in-plane field

To further examine how the current sheet geometry is affected by the guide field and cross-field plasma diffusion, a comparative measurement was done where $B_g/B_{xy}$ was unchanged, but where the amplitude of the guide field and the in-plane field was halved.

![Figure 5.10](image)

Figure 5.10: Comparison of two measurements with the same magnetic field ratio $B_g/B_{xy}$ but with different in-plane and guide magnetic field amplitudes. Azimuthal current density profile for: (a) $B_g = 18$ mT, and (b) $B_g = 9$ mT. (c) One-dimensional cuts through the current density maximum as indicated in (a) and (b). The dashed lines show one-dimensional Gaussian fits to the curves.
Table 5.1: Fit parameters for high and low magnetic field

<table>
<thead>
<tr>
<th></th>
<th>High B</th>
<th>Low B</th>
</tr>
</thead>
<tbody>
<tr>
<td>In-plane field $B_{xy}$</td>
<td>2mT</td>
<td>1mT</td>
</tr>
<tr>
<td>Guide field $B_g$</td>
<td>18mT</td>
<td>9mT</td>
</tr>
<tr>
<td>Peak current density $J_z$</td>
<td>-8.0 kA</td>
<td>-5.6 kA</td>
</tr>
<tr>
<td>Angle $\theta$</td>
<td>20.4°</td>
<td>13.8°</td>
</tr>
<tr>
<td>Area $A$</td>
<td>9 cm²</td>
<td>16 cm²</td>
</tr>
<tr>
<td>Major axis $L$</td>
<td>5.0 cm</td>
<td>6.5 cm</td>
</tr>
<tr>
<td>Minor axis $l$</td>
<td>2.3 cm</td>
<td>3.1 cm</td>
</tr>
<tr>
<td>Diffusion coefficient $D$</td>
<td>1.5 m²/s</td>
<td>4.5 m²/s</td>
</tr>
</tbody>
</table>

Figure 5.10a shows the measured current density profile at $z = 53$ cm for $B_g = 18$ mT and $B_{xy} = 2$ mT, and figure 5.10b for $B_g = 9$ mT and $B_{xy} = 1$ mT. One-dimensional cuts through the current density peak are shown in figure 5.10c together with Gaussian fits. The fit parameters for two-dimensional Gaussian fits to the current density profiles are given in table 5.1. From figure 5.10 and table 5.1 it can be seen that the peak current density is smaller and the current sheet has a larger area when the magnetic field is smaller. Correspondingly, the diffusion coefficient is also larger. This is further evidence that the current sheet geometry is not solely determined by the magnetic mapping due to the magnetic field ratio, but also through cross-field plasma diffusion.
5.3 Driven magnetic reconnection

The above discussed studies of the current sheet geometry where done without the reconnection drive. The main difference between the two driving methods is that for setup A the coils simultaneously establish the X-point topology, drive the magnetic flux and induce the electric field. In setup B the X-point topology is always present and the inductive electric field is independently created.

5.3.1 Magnetic flux

The common magnetic flux per axial length $\Phi / z$ is defined by the axial component of the vector potential at the X-point. $\Phi / z$ in vacuum and in a plasma is shown in figure 5.11a for setup A and in figure 5.11b for setup B. The axial plasma current is shown for reference. Indicated are the phases where the field lines are either pushed away from or pulled towards the parallel conductors setting up the X-point topology. For both setups there is an initial excess of flux before the reconnection drive starts due to the background current of the plasma gun. For setup A, the excess of flux disappears as the inductive electric field reduces the current to zero in the initial push reconnection phase (figure 5.11a). Since the total plasma current never changes direction, its contribution to the flux is always positive. Consequently, the magnitude of the flux with a plasma

![Figure 5.11](image-url)

Figure 5.11: Common magnetic flux per axial length in vacuum (solid blue line) and with a plasma (dotted red line) for: (a) Setup A, and (b) Setup B. The plasma current is plotted for reference (solid green line).
does not get smaller than the flux in vacuum. This can be seen in the initial push phase. In the first pull phase, there is an increase in the flux compared with the vacuum case, indicating a build-up of the common flux as expected (corresponding to reduction of the private flux). As the flux changes sign, and the second push phase sets in, there is a reduction of the common flux due to the build-up of the private flux. Finally, in the last pull phase, there is still a reduction and the common flux approaches the vacuum flux toward the end of the phase. This is because a current is still flowing due to the background current and the phase shift between the plasma current and the inductive electric field (cf. section 3.4.3), which results in a positive contribution to the common flux.

The magnetic flux per axial length for setup B, shown in figure 5.11b, is positive and is modulated by the reconnection drive (modulation $< 10\%$). The inductive electric field is never large enough to completely inhibit the plasma current and there is always a positive contribution of the background current to the flux. Consequently, the common flux is never reduced to the magnitude of the vacuum flux. In the first pull phase, there is again a build-up of the common flux as the current first returns to the background value and then increases due to the current induced by the reconnection drive. This continues into the second pull phase, where the difference between the vacuum and plasma measurement is the highest. The excess flux is reduced in the last push phase. As for setup A, the background and inductively extracted currents still contribute to the flux in this phase, and the vacuum magnitude of the common flux is never recovered.

For both setups, the modification of the vacuum flux is relatively small, with a maximum change in flux of 13\% for setup A and 2\% for setup B. This is due to the limitation of the current extracted from the gun. For setup B, the change is particularly small because the stationary magnetic flux is much higher than the modulation of the flux by the reconnection drive.

### 5.3.2 Electric field

The azimuthal profile and one-dimensional cuts through the X-point of the maximum of the modulus of the inductive electric field measured in vacuum is shown in figures 5.12(a) and (b) for setup A and 5.12(c) and (d) for setup B. For both setups, the
5.3 Driven magnetic reconnection

inductive electric field is largest close to the drive conductors, has a saddle point at the X-point and decreases towards zero at the wall of the vacuum chamber. The inductive electric field is much larger at the X-point for setup A compared to setup B \( E_{\text{ind}} = 180 \text{ V/m} \) and \( 27 \text{ V/m} \), respectively) due to the smaller separation between the conductors for setup A. This consequently leads to a larger inductively driven plasma current. The position of the saddle point does not match perfectly with the X-point for setup B but is located at \((x, y) = (-1.0, 0) \text{ cm}\). This is from a slight misalignment of the two conductor pairs.

5.3.3 Current sheet evolution for setup A

Figure 5.13 shows the axial current density and the modulus of the in-plane magnetic field for driven reconnection in setup A and with a guide field of \( B_g = 15 \text{ mT} \). The reconnection drive current, the associated inductive electric field at the X-point and the plasma current are displayed in figure 5.13a. It is evident that the electric field is large enough to initially inhibit the background current completely. The total current provided
by the plasma gun has a peak of $I_p = -37$ A. The axial current density $j_a$ and the modulus of the magnetic field at four time instants (corresponding to different phases of the reconnection drive as indicated by roman numerals I-IV in figure 5.13a) are shown in figure 5.13(b-e) and (f-i), respectively. The current density plots show that the current sheet is elongated along the separatrix. This is particularly clear in figure 5.13e, where

![Figure 5.13: Current sheet evolution in setup A. (a) Current trace of the reconnection drive (blue line), the associated electric field at the X-point (red line), and the plasma current (green line). Indicated are four time points I-IV for which the axial current density (b-e) and the modulus of the in-plane magnetic field (f-i) are plotted. Overlaid are contours of constant vector potential representing magnetic field lines. These are plotted for the same contour level values for all four time-instances. For reference, the contour level of the current density at twenty percent of the peak value is plotted (green)
5.3 Driven magnetic reconnection

Figure 5.14: (a) The axial current density and (b) the modulus of the in-plane magnetic field for a measurement with setup A and a guide field of $B_g = 30$ mT. For reference, the contour level of the current density at twenty percent of the peak value is plotted (green). Adapted from ref. [121].

it is oriented from lower left to top right. For figure 5.13(c-d), where the direction of the in-plane magnetic field is in the opposite direction, the current is elongated from the top left to the lower right. This is, however, less pronounced due to the lower current density and magnetic field at these time points. At I, when the magnetic field is large, the plasma current is small and has little influence on the magnetic field topology. As the magnetic field decreases and the plasma current increases, it starts to significantly contribute to the in-plane magnetic field at the center, resulting in an elongation of the in-plane neutral point into a neutral sheet (II). As the magnetic field generated by the conductors approaches zero, the in-plane magnetic field is dominated by the plasma current and has an O-shaped topology (III). This has also been reported in independent studies [122]. As the drive current reverses direction, the neutral sheet of the in-plane magnetic field is again formed but now in the push reconnection phase (IV). For a higher current density at the X-point, e.g., due to a larger background current or a stronger guide field, the contribution of the plasma current to the in-plane magnetic field at the center is dominant even at the phase when the magnetic field of the conductors is largest [121]. Such a case is shown in figure 5.14, where the neutral sheet of the in-plane magnetic field has broken apart to form three minima.

5.3.4 Current sheet evolution for setup B

Reconnection measurements using setup B are shown in figure 5.15 in the same representation as in figure 5.13. The current through the parallel conductors used to create the
X-point topology was 1.8 kA and the guide field was $B_g = 18$ mT. The inductive electric field is much lower than for setup A with a maximum of $E_{\text{ind}} = 27$ V/m at the X-point. This is reflected in a lower extracted current with a peak value of $I_p = -20$ A. The motion of a single field line is highlighted in figure 5.15(f-i). At time instant I the plasma gun discharge and reconnection drive has not yet started and there is no current flowing (figure 5.15b). Hence, the field topology is solely defined by the two parallel conductors. The field line remains stationary until the gun starts discharging (still without reconnection drive), after which the background current initially pushes the field lines away from
the conductors that create the X-point topology, until it reaches the position at time instant II. As the drive sets in, the field line is pushed further outwards until the drive current peaks and the inductive electric field changes direction (time instant III). At this point in time, the current mainly consists of the inductively inhibited background current and the reverse current (figure 5.15c). The inductive electric field then starts to extract an additional current as the field line starts to be pulled back. At time instant IV, the plasma current has reached its maximum. In figure 5.15e, the current density of the sheet has been increased due to the inductive electric field, while the reverse current has completely disappeared. Shortly afterwards, as the drive current reaches its maximum and the inductive electric field changes direction again, the field line is once more pushed away. From figure 5.15(f-i) it is clear that the neutral sheet of the in-plane magnetic field remains throughout the entire reconnection drive cycle and does not evolve into an O-shaped topology or even three minima as is the case for setup A.

5.3.5 Comparison of setup A and setup B

Both experimental setups have various advantages and drawbacks. For setup A, the coils producing the X-point topology also drive the magnetic flux and produce the axial inductive electric field. As a result, there is a significant change in the in-plane X-point topology, and the plasma current dominates the field topology for most of the phases of the drive, even resulting in an O-shaped topology. Setup B has the benefit that the inductive electric field, and thereby the plasma current, can be set largely independently of the amplitude of the magnetic field. It also has the advantage that the in-plane X-point topology remains throughout the drive cycle. As a consequence, the plasma current does not dominate the in-plane field topology in the same way. This, combined with the lower plasma current amplitude, makes it possible to avoid having the neutral sheet of the in-plane magnetic field breaking apart to form three minima. The lower inductive electric field does, however, mean that the modulation of the background current is smaller. Hence, the difference in the common flux between vacuum reconnection and reconnection with a plasma becomes smaller. Due to above given advantages, setup B was used for the results presented in the rest of this dissertation.
5.4 Influence of reconnection on the plasma current

Since the electric fields and the current flowing in the system are integral to the process of reconnection, it is useful to look in more detail into how they interact and the various factors influencing the current.

5.4.1 Current sheet geometry during reconnection

In terms of the current sheet geometry, it is clear that though the reconnection drive increases the current density, the general shape remains unaffected, i.e., the elongation and tilt of the sheet along the separatrix can be seen at all time phases. This differs from what is expected from the simplified picture of reconnection in two dimensions where the current sheet geometry is set by the motion of the field lines with the plasma and has a size dependent on the length scale where field line diffusion is significant. For example, Sweet-Parker like reconnection requires a long and thin current sheet to establish a resistivity that makes magnetic diffusion possible. Such a relationship between the current sheet and the reconnecting field lines is not observed. Sweet-Parker reconnection also requires a self-consistently formed current sheet, which is not the case in the measurements where the size is instead defined by the plasma gun aperture and the boundary conditions. If the current sheet is thereby externally set to be smaller than the relevant magnetic diffusion length scale, there would be no need for the current sheet to be narrowed further.

In experiments with guide field reconnection in the two-fluid regime, a tilt of the current sheet has also been observed. The tilt is explained by a $j \times B$ force set up by the interaction of the in-plane current and the guide field. In VINETA II the explanation for the tilt instead lies in the magnetic mapping. Hence, the shape of the current sheet is mainly set by a combination of the boundary conditions of the plasma gun, magnetic mapping and cross-field plasma diffusion.
5.4 Influence of reconnection on the plasma current

5.4.2 Current response to electric field

For setup B, the background current contributes more than a third of the total current, and it is obvious that it does not simply flow in response to the externally induced electric field. Instead, the current is a result of a total electric field consisting of: (a) the electrostatic field associated with the background current, (b) the externally imposed inductive electric field, and (c) an electrostatic field component opposing the inductive electric field and associated with the sheath limitation of the plasma current. While the inductive electric field has no reversal of direction within the azimuthal plane, the reverse current density observed at some time phases indicates that this is not true for the total electric field. Figure 5.16 shows the plasma current obtained by integrating the current density $j_z$ over the azimuthal plane for the measurement in figure 5.15. Also plotted are the contributions of the positive and negative current density to the total current. The positive component originates from the reverse current density that can be seen in figure 5.15(c-d). The positive component is largest when the inductive electric field is directed such to reduce the main current density in the sheet, and tends to zero at the peak of the total current. A reverse current was also observed at the edge of the current sheet in experiments conducted on CS-3D [123, 124]. The authors attribute the electric field reversal to an additional electric field $E' = v \times B$ due to the motion of outflow plasma jets with velocity $v$ in a strong transverse magnetic field. The origin of the electric field reversal is different in VINETA II. Here, it is from the inductive electric field being directed

![Figure 5.16: Plasma current obtained from the current density integrated over the area in the azimuthal plane (green line). Also plotted are the positive (red line) and negative (blue line) components of the total current.](image-url)
oppositely to the electrostatic field associated with the background current during some phases of the reconnection drive. The location of the reverse current density along the separatrix is consistent with magnetic mapping and a current flowing in the opposite direction of the main current in the sheet. Considering only the total current, it would be tempting to make the wrong conclusion that the current sheet disappears when the total current goes to zero. The reverse current may influence the electric field if it is large enough to significantly alter the temporal evolution of the current.

The influence of the plasma current on the vacuum inductive electric field can be studied from the difference $\Delta E$ between the inductive electric field with and without plasma. Figure 5.17a shows the axial inductive electric field at the X-point with a plasma and in vacuum, and figure 5.17b shows $\Delta E$ (for $B_g = 18$ mT). The electric field difference is clearly proportional to the time-derivative of the current, i.e., to the time derivative of the magnetic field generated by the plasma current. It has a maximum of $\Delta E = 2.7$ V/m at $t = 6.1 \mu$s. Similar results were obtained in VTF, where the reconnection electric field was found to peak when the change in current density $dj/dt$ at the X-point was largest [125].

The inductive electric field with a plasma can be described as a superposition of the vacuum field of the drive conductors and a field set up by the temporal evolution of
the plasma current. The profile of $\Delta E$ in the azimuthal plane at $t = 6.1 \mu s$ is shown in figure 5.17c. The difference is highest at the position where the current density has a peak and decreases with distance from the current sheet. The magnitude of the peak value of $\Delta E$ depends on the temporal evolution of the total current, but is nearly independent of the in-plane redistribution of the current density due to magnetic mapping and cross-field plasma diffusion.

### 5.4.3 Guide field dependence of plasma current

It has already been demonstrated that the peak current density shows a strong dependence on the guide field and the axial position. The change in the plasma current with the guide field has also investigated. The guide field can affect the current in a number of different ways: (a) by influencing the resistivity through changes in temperature and density, (b) by affecting the plasma gun discharge, and (c) by altering the boundary conditions. Lacking a complete picture of the evolution of the plasma parameters during the reconnection drive phase, it is not possible to determine the significant factors that dominate the guide field dependence of the plasma current.

Figure 5.18 shows Rogowski coil measurements of the plasma current for high and low

![Figure 5.18](image-url)
guide fields at two different axial positions: behind the first grid at a distance of $z = 20$ cm from the gun, and before the second grid at a distance of $z = 90$ cm. The current drawn by grid 2 to ground is also plotted in figure 5.18. Comparing the three currents for $B_g = 6$ mT, the peak value of the current $\hat{I}_{p}$ decreases with distance from the gun ($\hat{I}_{p} = 20.3$ A at $z = 20$ cm and $\hat{I}_{p} = 15.6$ A at $z = 90$ cm). While field line mapping and cross-field plasma diffusion do not directly influence the amplitude of the current, the current sheet is elongated past the area of the Rogowski coil for low guide fields (at $z = 90$ cm), which systematically reduces the measured current. Since the size of the Rogowski coil is comparable to the grids, part of the current flow goes through the device wall instead of grid 2. This leads to a smaller peak current of $\hat{I}_{p} = 11.9$ A through grid 2. The peak also occurs later for the current to grid 2 than for the plasma current ($t = 11.5 \, \mu s$ and $t = 10.1 \, \mu s$, respectively). In contrast, for $B_g = 90$ mT, the three current traces are very similar and only exhibit a small difference in amplitude and phase ($\hat{I}_{p} = 18.3$ A at $t = 11.3 \, \mu s$ for $z = 20$ cm and $90$ cm, and $\hat{I}_{grid2} = 17.4$ A at $t = 11.7 \, \mu s$ for the current to grid 2). The increase in the guide field has also led to a decrease in the amplitude of the current at $z = 20$ cm compared to the measurements for $B_g = 6$ mT, while the current at $z = 90$ cm and the current flowing to grid 2 has increased. Furthermore, the peaks of $I_p$ at $z = 20$ cm and $z = 90$ cm now occur later, while for the current to the grid it occurs earlier. The similarity in amplitude and phase of the three currents indicate that the current closure mainly occurs through the grid. The dependence of the peak current amplitude on distance from the gun and guide field strength is discussed further below.
5.4 Influence of reconnection on the plasma current

Figure 5.20: Plasma current for different guide field strengths. The blue markers show the current at a distance of \( z = 20 \) cm from the gun and the green markers the current flowing from grid 2 to ground. (a) Background current with an X-point topology. (b) Reduction in current due to the X-point topology as obtained by subtracting the background current from the current in (a).

Figure 5.19a shows the peak amplitude of the current at \( z = 20 \) cm and to grid 2 for different guide field amplitudes. The amplitude of the current to the grid is smallest at low guide fields, which is consistent with more current flowing through the walls instead of the grid. At the \( z = 20 \) cm position, the current first increases as the guide field is increased, with a peak at around \( B_g = 18 \) mT, and then decreases as the guide field is increased further. For the high guide field measurements, the amplitudes are about the same at the two axial positions, indicating that most of the current flows through the grid. For \( B_g > 30 \) mT, the current amplitude only shows a weak dependence on the increase of the guide field. If the current change at low guide fields is mainly due to the current sheet expanding past the Rogowski coil and grid 2, then above the guide field strength where the entirety of the current flows through the grid, the current amplitude should remain unchanged. An estimate of the guide field magnitude required for the current sheet to reach the edge of the grid can be obtained by mapping the circular area of the plasma gun aperture to the grid position and adding the increase in radius by cross-field diffusion. The latter is estimated by the increase in area of the background current sheet between the measurement planes of the \( \dot{B} \)-probes. The contour level at 10% of the peak current density would be expected to reach the edge of the grid around \( B_g = 24 \) mT.

To further understand the dependence between the current and the guide field, three separate measurements were done: (1) with background current only, (2) with background current and X-point topology, and (3) with background current and reconnection drive.
Figure 5.21: Plasma current for different guide field strengths. The blue markers show the current at a distance of $z = 20$ cm from the gun and the green markers the current flowing from grid 2 to ground. (a) Plasma current with the reconnection drive (without the X-point topology field). (b) Current due to the reconnection drive as obtained by subtracting the background current from the current in (a).

This makes it possible to determine the individual contributions to the current shown in figure 5.19a. Figure 5.19b shows that the amplitude of the background current decreases with increasing guide field strength. A small difference between the current to the grid and the current measured at $z = 20$ cm is observed, but diminishes with increasing guide field until they have the same amplitude at $B_g > 15$ mT. This can be attributed to a decrease in the cross-field plasma diffusion as the guide field is increased. However, from the measurements of the expansion of the current sheet, the sheet edge is estimated to reach the grid edge only when the guide field is around $B_g = 6$ mT. It is difficult to determine the reason for the observed guide field dependence of the background current since it is highly dependent on the axial electrostatic field and the discharge parameters of the plasma, neither of which are well-defined and both of which are affected by the guide field strength.

Figure 5.20a shows the background current in the X-point topology. The current displays a similar guide field dependence as in figure 5.19a, with the current to grid 2 being smallest at low guide fields and the plasma current measured at $z = 20$ cm having a peak around $B_g = 18$ mT. Figure 5.20b shows how the X-point topology affects the amplitude of the background current, by subtracting the background current in figure 5.19b from the total current in figure 5.20a. The difference is negative for low guide fields and hence the reduction in the current is then largest. For high guide fields, the reduction is small,
which indicates that the X-point topology then has a small influence only. Figure 5.20 also shows that the reduction becomes larger than the increase in the background current for $B_g < 18 \text{ mT}$. Magnetic mapping and cross-field diffusion of the current sheet past the edge of the Rogowski coil at $z = 20 \text{ cm}$ can be ruled out and the reduction in the measured current is at least not fully due to the diagnostics.

The current generated by the reconnection drive without the X-point topology is shown in figure 5.21a. The current decreases with increasing guide field. It is also higher than the currents in figure 5.19a, indicating that the X-point topology reduces the current. This is attributed to a combination of the field line mapping and a reduction in the current by the inductive field that is created by the small variations of the current producing the X-point topology. The contribution of the reconnection drive to the current is shown in figure 5.21b. The current to the grid is largely unaffected by changes in the guide field. It is smaller than the current at $z = 20 \text{ cm}$ for all guide field strengths. At $z = 20 \text{ cm}$, the current has a small increase in amplitude for $B_g < 18 \text{ mT}$.

From figures 5.19 - 5.21 it is concluded that the guide field dependence of the current is mainly set by a combination of the background current and the X-point topology, while the amplitude of the inductively drawn current remains largely unchanged. As described previously, it is difficult to pinpoint exactly what governs the change in current. It is most likely due to a combination of the geometry of the current sheet and the consequent change in the current path, as well as a change in the plasma parameters.
5 Guide field reconnection in the VINETA II device

5.5 Plasma parameters during reconnection

After the investigation of the time-evolution of the plasma current and the current sheet geometry, the next step is to characterize how the plasma parameters change due to the X-point topology and the reconnection. Noise pick-up due to the reconnection drive, as outlined in section 4.1, is an issue and in an attempt to reduce it, the second switching of the IGBT was omitted for the measurements described in this section. This has the additional effect of reducing the inductive electric field, and thereby the plasma current.

Figure 5.22: (a) Axial current density profile in the azimuthal plane. (b) In-plane magnetic field. (c) In-plane current density. The arrows indicate the direction of the current flow. (d) Time-dependent magnetic field in the $z$-direction. For reference, the contour level at 20% of the peak value of the axial current density and the separatrix of the in-plane field are also plotted (green and black, respectively).
5.5 Plasma parameters during reconnection

5.5.1 Magnetic field and current density

Figure 5.22a shows the axial current density at the time point when the plasma current has a peak. The associated magnetic field is shown in figure 5.22b. The measurement was carried out with similar parameters as for the measurements shown in figure 5.15 ($B_g = 18$ mT). Since the inductive electric field is lower, the current density is lower as well. The in-plane current is shown in figure 5.22c and the associated out-of-plane magnetic field is plotted in figure 5.22d. Similar to the case with only the background current, the current circulates in the azimuthal plane and follows the geometry of the current sheet. The current does not exhibit the flow pattern expected from Hall-currents and consequently the magnetic field does not have a quadrupolar structure. The magnetic field instead has a peak value at the same position as the axial current density. The origin of the in-plane current is discussed in more detail in section 5.6.2.

5.5.2 Plasma potential

The time-evolution of the plasma potential and the floating potential at the current density peak is plotted in figure 5.23a. Both the floating potential and the plasma potential are modulated by the reconnection drive, but neither of them are in phase with the current trace. The plasma potential peaks slightly before and the floating potential after the current. The plasma potential measurement is also more sensitive to the high frequency noise produced by the switching of the IGBT than the floating potential measurement. This is most likely a filtering effect due to the change in impedance of the probe circuit between hot and cold operation.

Figure 5.23(b-c) shows the plasma potential measurement in the azimuthal plane for a time instance before the reconnection drive and for the time instance when plasma current has a peak ($t_1$ and $t_2$ as indicated in figure 5.23a). As for the plasma potential measurement discussed in section 5.1.2, the plasma potential is negative in the region of the current sheet and positive at the edge of the measurement region. With the X-point topology, the plasma potential profile is elongated along the separatrix. At $t_2$ the position of the peak has shifted outside the current sheet to $(x, y) = (5.5, 1.5)$ cm. The main region of interest lies within the current sheet where the bulk of the plasma is located.
Additionally, the plasma potential measurement is more sensitive to noise outside the current sheet due to the lower plasma density.

The gradient of the plasma potential yields the in-plane electrostatic field shown in figure 5.23(d-e) for $t_1$ and $t_2$. At $t_1$, the electrostatic field points radially inwards, has a minimum near the X-point, and is largest outside the current sheet where the plasma potential changes sign. At $t_2$, the electrostatic field has increased due to a steepening
5.5 Plasma parameters during reconnection

Figure 5.24: (a) Time-evolution of the electron temperature at the position of the current density peak (black line). The plasma current is plotted for reference (green line). (b-c) Temperature profile in the azimuthal plane for the time points $t_1$ and $t_2$ as indicated in (a). The contour level at 20% of the peak current density is plotted in green and the separatrix in black.

of the plasma potential. Since the plasma potential no longer has a minimum inside the current sheet, the electrostatic field points towards a position outside the sheet located at $(x, y) = (5.5, 1.5)$ cm. It has a mean value of 170 V/m inside the current sheet, which is much higher than the axial inductive electric field of the reconnection drive. Similar results were reported in [126] where the in-plane electrostatic field was found to be an order of magnitude higher than the axial inductive electric field.

5.5.3 Temperature

As before, the electron temperature is obtained from the difference between the plasma potential and the floating potential. The temporal evolution of the electron temperature at the location of the current density peak is plotted in figure 5.24a. At $t_1$ the temperature
has reached a fairly constant value of $T_e = 4$ eV, and is later on modulated by the reconnection drive. The temperature profile in the azimuthal plane is plotted in figure 5.24b for $t_1$ and in figure 5.24c for $t_2$. The temperature has a local maximum at the center position and is elongated similar to the current sheet. It decreases towards the edge of the current sheet, before again rising at the edge of the measurements region. At $t_2$ the temperature has globally decreased and reaches unphysical negative values at the edge of the current sheet.

The thermal equilibration time ($\tau_m = 0.4 \mu s$) and the slowing down time for the electrons ($\tau = 0.6 \mu s$), is an order of magnitude smaller than the electron transit time $\tau_{tr} = d ne / j_z$, with $d$ being the distance between the plasma gun and grid 2. The energy dissipation of electrons is thus expected to play an important role and ohmic heating $P = E \cdot j$ should contribute to the temperature. The maximum of the axial component $P_z = E_{ind} \cdot j_z$ is shown in figure 5.25. The ohmic heating is highest where the current density has a peak, with a value of $P_z = 176$ kW/m$^3$, and falls off to zero outside the current sheet. This is consistent with the temperature profiles in figure 5.24. Note that, though the in-plane currents are small, the in-plane electric field is larger than the axial inductive electric field and the in-plane component contributes to the total dissipation ($\langle P_{xy} \rangle \approx 50$ kW/m$^3$ inside the current sheet). However, since the in-plane electric field is less reliable than the axial inductive electric field, this has been omitted here. For comparison, the spiral antenna operating at a power of 350 W in a cylindrical volume of radius 10 cm and a length of 1 m yields a power density of $P = 11$ kW/m$^3$. Though $P_z$ is probably overestimated since it does not incorporate any axial electrostatic field contribution, calculating it from the current density and an estimate of the resistivity ($P_z = \eta_s j_z^2$) still yields a significant peak
value of $P_z = 40 \text{ kW/m}^3$. Thus, Ohmic heating would be expected to have a significant contribution.

From Ohmic heating, the electron temperature is expected to increase with the current density. Triple probe measurements of the temperature indicate that this is the case [91]. However, it is not consistent with the temperature obtained from the difference in the plasma potential and floating potential, which instead increases as the plasma current decreases (figure 5.24a). The observed temporal evolution of the temperature could be an artifact of the diagnostic technique. As previously discussed, all diagnostics pick up the noise of the reconnection drive. From figure 5.23a, it is evident that the plasma potential measurement is sensitive to the high frequency component of the noise, and it is possible that the modulation of the temperature originates in the lower frequency component of the reconnection drive noise. The evaluation of the temperature also assumes that the plasma parameters do not change between the consecutive measurement of the azimuthal profiles of the plasma potential and the floating potential. Furthermore, if fast electrons are present, the floating potential measurements would be highly affected. On the other hand, if figure 5.24a is a reliable representation of the temporal evolution of the temperature, it would require cooling of the electrons. This is possible only through anomalous processes like turbulent transport (e.g. via electron temperature gradient (ETG) turbulence [127]).

### 5.5.4 MHD force balance and density

The $\mathbf{j} \times \mathbf{B}$-force during the reconnection phase can be divided into two components corresponding to $\mathbf{j}_z \times \mathbf{B}_{xy}$ and $\mathbf{j}_{xy} \times \mathbf{B}_g$. The former is shown in figure 5.26a for the time point of the current density peak. The $\mathbf{j}_z \times \mathbf{B}_{xy}$-force accelerates the plasma fluid towards the X-point from the top right and bottom left regions, and away from the X-point in the bottom right and top left part. Since the current sheet is not centered at the X-point, the force has a larger amplitude in the top right quadrant. Figure 5.26b shows the $\mathbf{j}_{xy} \times \mathbf{B}_g$ component of the force. It points radially inwards, has maxima where the in-plane current is largest, and a minimum at the center of the current sheet where the in-plane current is small. The total $\mathbf{j} \times \mathbf{B}$-force is shown in figure 5.26c. Though the in-plane current is smaller than the axial current, the guide magnetic field is much larger than the in-plane magnetic field. Consequently the $\mathbf{j}_{xy} \times \mathbf{B}_g$ component has the biggest
5 Guide field reconnection in the VINETA II device

Figure 5.26: (a) $j \times B_{xy}$ component of the $j \times B$-force. (b) $j_{x(y)} \times B_y$ component of the $j \times B$-force. (c) total $j \times B$-force. (d) Estimate of the $\nabla p$-force. The arrows indicate the direction of the forces. For reference, the contour level at 20\% of the peak current density is plotted in green and the separatrix in black.

contribution to the total $j \times B$-force. The flow direction is predominantly radial and, compared to figure 5.26b, the inflow is increased in the top right and lower left quadrant, and decreased in the top left and lower right quadrant.

The force due to the pressure gradient $\nabla p$ is obtained using the same approach as in section 5.1.2. The pressure $p = n k_B T_e$ is determined assuming ideal MHD force balance and a density that follows the current density profile. Using the temperature profile from figure 5.24 and estimating the density from ideal MHD force balance, one obtains a $\nabla p$-force profile as shown in figure 5.26d. It points radially outwards and has maxima where the gradient of the current sheet is largest.

The two forces in figure 5.26c and figure 5.26d have similar structures, but with a shift of the location of the maxima. Figure 5.27 shows the residual of the two forces. Inside the current sheet, the residual has a mean value of 4.4 Pa·m. The forces do not match as well
as for the gun plasma without the X-point topology and reconnection (cf. figure 5.5). The discrepancy can be explained by a combination of the ideal MHD-force balance description not being fully valid during the reconnection drive phase, and uncertainties in the assumptions made in the estimate of the $\nabla p$-force.

Figure 5.27: Residual of the $j \times B$-force and the $\nabla p$-force. The color indicates magnitude and the arrows the direction. The contour level at twenty percent of the peak current density and the separatrix area plotted for reference, (green and black, respectively).
5.6 In-plane particle drifts and currents

5.6.1 Particle drifts

The fields and gradients in the plasma give rise to particle motions. From the magnetic and electric fields and the plasma parameters discussed in previous sections, it is possible to obtain estimates for some of the particle drift velocities. The in-plane $E \times B$-drift is plotted in figure 5.28. It can be divided into two components corresponding to the drift caused by the in-plane magnetic field and the axial inductive electric field (figure 5.28a), and a drift due to the in-plane electrostatic field and the guide field (figure 5.28b). For the reconnection dynamics, the drift due to the in-plane magnetic field is the more important one since it is responsible for transporting particles in and out of the diffusion region. The drift velocity is largest near the conductors where the in-plane field is largest and becomes small at the X-point. This drift is determined by the guide field, and differs significantly from what would be expected in two-dimensional reconnection. Without the guide field, the amplitude of the velocity is given by $E_{\text{ind}}/B_{xy}$, and would instead increase towards the X-point, where it diverges due to the vanishing magnetic field. The $E_{xy} \times B_g$-drift shown in figure 5.28b is highest outside the current sheet where the in-plane electrostatic field is the largest. The velocity is two orders of magnitude larger than the drift in figure 5.28a. Even if it is taken into account that the electrostatic field measurement is reliable only inside the current sheet and ignoring the drifts outside, this component of the $E \times B$-drift still clearly dominates.
5.6 In-plane particle drifts and currents

Another important drift is the electron diamagnetic drift, which is given by $v_{\text{dia,e}} = -\nabla p \times B / (e n B^2)$ and shown in figure 5.29a. Due to the steep density gradients, the diamagnetic drift is large compared to the $E_{\text{ind}} \times B_{xy}$-drift and is in the same direction as the total $E \times B$-drift. A further drift of interest is the polarization drift $v_{\text{pol}} = (m/qB^2) dE/dt$. Though it is largest outside the sheet, as expected from the magnitude of the electrostatic field, it is still significant inside the sheet where it has a mean value of $v_{\text{pol}} = 43$ km/s. Outside the sheet, it is mainly pointing radially outwards, while inside the sheet no preferential direction can be discerned.

The relative amplitude of the different drifts shows how the combination of a strong guide field and steep gradients in the plasma dominate the particle motion. A strong guide field means that the $E_{\text{ind}} \times B_{xy}$-drift is small and thus the transport of plasma towards the X-point by $E \times B$-drift becomes small. At the same time, a high guide field results in strong $E_{xy} \times B_g$- and diamagnetic drifts.

5.6.2 In-plane currents

As previously shown, there is an in-plane current flowing in the plasma. This current and associated magnetic field are shown in figure 5.30 for $B_g = 18$ mT and $B_g = 6$ mT. As demonstrated in section 5.5.1, for a guide field of $B_g = 18$ mT, the current flow circulates within the current sheet, follows the current sheet geometry and has local maxima where the azimuthal profile of the axial current density has the steepest gradients. The time
Figure 5.30: (a) In-plane current density $j_{xy} = \nabla \times B / \mu_0$ for $B_g = 18 \text{ mT}$ and $6 \text{ mT}$. The direction of the current flow is indicated by the arrows. (b) Time-varying magnetic field in $z$-direction dependent out-of-plane magnetic field created by the current has a peak at the position of the axial current density peak. For the lower guide field case, the current is significantly altered. It flows inward towards the X-point and is highest in the top left and lower right quadrants. The current results in an out-of-plane magnetic field with a quadrupolar structure (figure 5.30b). Although an out-of-plane quadrupole magnetic field can be a signature that two-fluid effects play a role [128], the measured field differs from what is expected in the Hall model where the differential particle flow set up by the motion of the field lines results in a field with maxima located along the separatrix. Furthermore, experiments and simulations of magnetic reconnection including two-fluid effects show that the quadrupole field is severely reduced or even has completely vanished already when the magnitude of the guide field is comparable to the reconnection magnetic field [12, 129]. The in-plane current and associated magnetic field must then have some other origin, which is discussed below.
In-plane particle drifts and currents

\[ j = n e E \times B / B \times B \text{ ind} \]
\[ j = n e E \times B / B \times B \text{ xy} \]

Figure 5.31: In-plane current set up by the \( E \times B \)-drift for \( B_g = 18 \text{ mT} \) and \( 6 \text{ mT} \) (assuming unmagnetized ions). The direction of the current flow is indicated by arrows. (a) Current due to \( E_{xy} \times B_g \)-drift. (b) Current due to \( E_{\text{ind}} \times B_{xy} \)-drift.

The drifts described in section 5.6.1 can be used to determine the origin of the observed in-plane current. Though the \( E \times B \)-drift usually does not in contribute to currents since it affects ions and electrons equally, if only the electrons are magnetized their motion may well result in a net current. This current is depicted in figure 5.31a for the \( E_{\text{ind}} \times B_{xy} \)-drift and in figure 5.31b for the \( E_{xy} \times B_g \)-drift. The former is associated with a current due to the field line motion and has similarities to what is expected from the Hall effect. The amplitude of this current is at least two orders of magnitude smaller than the measured current and approximately an order of magnitude smaller than the noise level and is, hence, not expected to contribute to the out-of-plane quadrupolar magnetic field. The \( E_{xy} \times B_g \)-drift sets up a circulating current flow. The current pattern is distorted and not centered at the peak value of the axial current density. The time-dependence of the
in-plane electric field generates an ion polarization drift, and the associated current is plotted in figure 5.32. The amplitude is larger than both the observed current and the other current contributions listed above by at least a factor of three. Though it appears to be mainly directed radially outwards, there is no clear flow pattern. Since such a current is not seen in the measurement, it is believed that the polarization drift velocity is not a significant contributor.

Due to the steep electron pressure gradients, there is a significant electron diamagnetic drift within the current sheet. The resulting current is shown in figure 5.33 for the two guide fields. The current circulates within the current sheet and, as expected, has local

![Figure 5.32: In-plane ion polarization drift current for $B_g = 18$ mT and 6 mT. Color indicates magnitude and arrows direction.](image)

![Figure 5.33: In-plane component of electron diamagnetic current for $B_g = 18$ mT and 6 mT. The direction of the current flow is indicated by arrows.](image)
maxima where the gradients are steepest. For $B_g = 18 \text{ mT}$, the diamagnetic current is very similar to the observed current in figure 5.30a. This shows that the diamagnetic current is a strong contributor to the total in-plane current. For $B_g = 6 \text{ mT}$ the difference is more pronounced. While the observed current is mainly directed towards the X-point from the left and the right, the diamagnetic current flows in both directions (figure 5.33), which means another current contribution must play a significant role.

As already pointed out, the field-aligned current has an in-plane component which becomes significant at low values of $B_g/B_{xy}$. The resulting current density is plotted in figure 5.34. It exhibits a flow along the separatrix and is directed towards the X-point from the left and the right, and directed away from it going towards the top and the bottom. Though all of the above discussed currents probably contribute to some degree to the in-plane current seen in the measurement, the most significant contributions come from the field-aligned current and the diamagnetic current (possibly with some contribution of the $E_{xy} \times B_g$-drift). The combination of these two components is shown in figure 5.35. An increase in the amplitude of the current is observed in the top left and bottom right quadrant where their components flow in the same direction. Conversely, the current is reduced in the top right and the bottom left quadrant where the components have opposite directions. This is particularly evident for the measurement with a lower guide field. Comparing the current pattern for $B_g = 6 \text{ mT}$ in figure 5.35 with the measured pattern in figure 5.30, a quite good qualitative and quantitative agreement is found. It is concluded that the in-plane currents mainly originate from the steep temperature
and density gradients in the current sheet and the relative amplitude of the guide and
in-plane magnetic field. It should be noted that the diamagnetic current and the in-plane
component of the field aligned current are present even without the reconnection drive.
If it can be assumed that the density scales with the current density, the inductive electric
field mainly affects the in-plane current by increasing the amplitude of the axial current
density and thereby steepening the pressure gradients.
5.7 Generalized Ohm’s law

A key question regarding magnetic reconnection is which of the terms in the generalized Ohm’s law are significant. The relative importance of the terms can be assessed based on their characteristic length scales. These are compiled in table 5.2 for a magnetic field of $B_g = 18$ mT, a density of $n_e = 10^{18}$ m$^{-3}$ and a temperature of $T_e = 4$ eV.

$L_{\text{Hall}}$ is on the order of one meter and, compared to the other values which are on the order of a few millimeters, indicates that the Hall term should dominate. All other terms are similar in magnitude and much smaller than the dimension of the current sheet which is on the order of few centimeters. While this indicates that they are not important, given by the high collisionality of the plasma, the resistivity term is expected to be significant.

The amplitude of each term in the generalized Ohm’s law (2.15) is obtained using the plasma parameters discussed in the previous sections. The amplitudes of the terms, averaged over a region given by the contour level at half of the peak current density, are given in table 5.3. The electron inertia term $(m_e/ne)\partial j_{\parallel}/\partial t$ is orders of magnitude smaller than all the other terms and can be neglected. The Hall $j \times B/ne$ and the pressure term $\nabla p_e$ are both significantly larger than the resistivity term $\eta s j$. However, the Hall term and the pressure term are of similar magnitude but different sign and cancel out when combined. This is due to the fact that the Hall term is dominated by the in-plane current which is mainly given by the electron diamagnetic current driven by the electron pressure gradient.

The parallel component of the electric field at the X-point is important in terms of the axial current and the reconnection magnetic field, and it is useful to divide the generalized Ohm’s law into a perpendicular and a parallel component. Since the in-plane magnetic field is zero at the X-point, the $v \times B$ term and the Hall term both vanishes and the

<table>
<thead>
<tr>
<th>$L_{\text{resistivity}}$</th>
<th>$L_{\text{Hall}}$</th>
<th>$L_{\text{pressure}}$</th>
<th>$L_{\text{inertia}}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$1.8 \cdot 10^{-3}$ m</td>
<td>$1.4 \cdot 10^{9}$ m</td>
<td>$7.1 \cdot 10^{-3}$ m</td>
<td>$5.3 \cdot 10^{-3}$ m</td>
</tr>
</tbody>
</table>

Table 5.2: Characteristic length scales of the different terms in the generalized Ohm’s law for a magnetic field of $B_g = 18$ mT, a density of $n_e = 10^{18}$ m$^{-3}$, and a temperature of $T_e = 4$ eV.
5 Guide field reconnection in the VINETA II device

\[
\langle \eta_s j \rangle \quad \langle j \times B/(ne) \rangle \quad \langle \nabla p_e/(ne) \rangle \quad \langle (m_e/ne^2) \partial j/\partial t \rangle
\]

<p>| | | | |</p>
<table>
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<th></th>
<th></th>
<th></th>
<th></th>
</tr>
</thead>
<tbody>
<tr>
<td>1.7 ± 0.12</td>
<td>252 ± 31</td>
<td>226 ± 31</td>
<td>0.07 ± 0.02</td>
</tr>
</tbody>
</table>

Table 5.3: Magnitudes of the terms in the generalized Ohm’s law averaged over a region given by the contour level at half of the peak current density.

parallel Ohm’s law reduces to

\[
E_\parallel = -\left( \frac{\partial A_\parallel}{\partial t} + \nabla \phi \right) = (\eta_s + \eta_{\text{anom}})J_\parallel - \frac{1}{ne} \nabla p_\parallel + \frac{m_e}{ne^2} \frac{\partial j_\parallel}{\partial t}, \tag{5.1}
\]

where the electrostatic field and the anomalous resistivity have been included. The time-evolution of the terms in the parallel Ohm’s law at the X-point is shown in figure 5.36. The electron inertia term is negligible and has been omitted. The electric field created by the axial pressure term goes in the same direction as the inductive field. It is evident that these terms are not sufficient to balance the inductive field, hence, there must be an additional contribution. This term has been denoted \( E^\ast \) and is also plotted in figure 5.36b. This contribution could be associated with anomalous resistivity in the current sheet, which would then be an order of magnitude larger than the Spitzer resistivity. Another candidate is an electrostatic field opposing the inductive electric field. Such a field would arise if the plasma gun does not provide enough additional electrons to counteract the current limitation due to the axial boundary conditions. This electrostatic field would be required to have a peak amplitude of 35 V/m in order to balance the

![Figure 5.36](image-url)

Figure 5.36: Time-evolution of the terms in the parallel generalized Ohm’s law at the X-point. The inductive electric field is given by the blue line, the pressure gradient by the red line, and the resistivity term by the green line. The residual of these three terms is given by the black line.
contributions of the inductive electric field and the axial pressure gradient term.

## 5.8 Reconnection rate

One of the key parameters of magnetic reconnection is the rate at which the reconnection proceeds. In the Sweet-Parker model, the reconnection rate is defined as the ratio between the inflow velocity and the Alfvén-speed at the edge of the current sheet. Taking the $E_{\text{ind}} \times B_{xy}$-drift velocity as the inflow velocity and the Alfvén-speed calculated from the plasma parameters, one obtains a Sweet-Parker rate of $v_{\text{in}}/v_A = 5 \cdot 10^{-3}$. A more general approach is to define the reconnection rate via the magnetic flux transferred across the separatrix from the inflow region to the outflow region, which is given by the axial inductive electric field at the X-point. Without a plasma, there is nothing to inhibit the motion of the field lines, and consequently this case gives the upper limit of the reconnection rate. In order to compare the reconnection rate for different scenarios (varying guide field, density etc.), it is useful to normalize it to this vacuum value.

The reconnection rate is taken at the point in time when the plasma has the greatest influence on the reconnection, i.e., when the electric field difference $\Delta E$ at the X-point is largest (cf. figure 5.17 in section 5.4.2). Since $\Delta E$ is proportional to the time-derivative of the current, the reconnection rate depends on the axial current flow. The normalized

![Figure 5.37](image.png)  

**Figure 5.37:** (a) Reconnection rate $E_{\text{ind},p}/E_{\text{ind},v}$ at different guide field strengths, measured at three axial positions. (b) Reconnection rate at different ratios of the guide and in-plane magnetic field $B_g/B_{xy}$. The green line represents the standard case, the red line measurements with half the in-plane field and the blue line measurements with a high background plasma density.
5 Guide field reconnection in the VINETA II device

<table>
<thead>
<tr>
<th></th>
<th>Case I</th>
<th>Case II</th>
<th>Case III</th>
</tr>
</thead>
<tbody>
<tr>
<td>$B_g$</td>
<td>6 – 78 mT</td>
<td>6 – 60 mT</td>
<td>12 – 60 mT</td>
</tr>
<tr>
<td>$B_{xy}$</td>
<td>2 mT</td>
<td>1 mT</td>
<td>2 mT</td>
</tr>
<tr>
<td>rf-discharge</td>
<td>inductive</td>
<td>inductive</td>
<td>helicon</td>
</tr>
</tbody>
</table>

Table 5.4: Relevant parameters for three cases for which the reconnection rate is determined.

reconnection rate $E_{\text{ind},p}/E_{\text{ind},v}$ is shown in figure 5.37a for three axial positions ($z = 39$ cm, 53 cm and 67 cm). It is averaged over five reconnection shots and the error bars are given by the standard deviation. The rate is highest for the lowest guide field, where it almost reaches the vacuum value, and then quickly decreases and approaches a plateau level for $B_g \geq 20$ mT, with only a slight decrease at larger guide fields. This is consistent with the current measurements in section 5.4.3 and the reduction in the measured current due to the X-point topology, which is largest for small guide fields and almost zero for $B_g \geq 18$ mT. The reconnection rate seem to decrease slightly with axial distance from the plasma gun for $B_g \leq 20$ mT with a maximum difference in the rate of 0.06 between $z = 39$ cm and $z = 67$ cm. This indicates that there is a gradient in the inductive electric field along the $z$-axis, which in turn implies that there is a reduction in the axial current along the device. This requires that part of the current is closing radially to the wall. It is, however, unclear why this would be the case since magnetic mapping and cross-field transport should not be sufficient for the edge of the current sheet to reach the wall even for lowest guide field case.

To further investigate the dependence of the reconnection rate on the ratio of the guide field and the in-plane field, measurements were done with an in-plane field reduced by half. This increases the maximum ratio of the guide and the in-plane magnetic field by a factor of two. The reconnection rate at $z = 53$ cm is plotted against the magnetic field ratio $B_g/B_{xy}$ in figure 5.37b for three different cases: (I) the standard case, which corresponds to similar plasma parameters and in-plane magnetic field as for the measurements presented before, (II) the case with a lower in-plane magnetic field, and (III) the case with a high background plasma density. The relevant parameters for the three cases are compiled in table 5.4. $B_{xy} = 2$ mT is used for case I, which corresponds to position far enough outside the current sheet that the magnitude is largely unaffected by the magnetic field generated by the plasma current. The reconnection rate for case II is lower than the standard case and does not exhibit the same steep rise at lower guide fields. Sim-
5.8 Reconnection rate

Similar to case I, it exhibits a small reduction as $B_g/B_{xy}$ increases. Thus, the reconnection rate does not simply saturate at high guide fields. Combining these two measurements, the reconnection rate is approximately $E_{\text{ind,p}}/E_{\text{ind,v}} = 0.95$ at $B_g/B_{xy} = 3$, quickly decreases to approximately $E_{\text{ind,p}}/E_{\text{ind,v}} = 0.87$ at $B_g/B_{xy} = 9$ and then decreases slightly further for larger field ratios, eventually reaching a rate of approximately $E_{\text{ind,p}}/E_{\text{ind,v}} = 0.85$ at $B_g/B_{xy} = 60$. Also plotted in figure 5.37b is the reconnection rate for a measurement with a high density helicon discharge (case III). While the density does not remain constant as the guide field is changed, it is still an order of magnitude higher than for the inductive discharge mode for all guide fields. The reconnection rate is systematically lower than for the standard case. The decrease in the rate corresponds to an increase in $\Delta E$ and $dI_p/dt$. As discussed in section 3.4.3, increasing the background density results in a higher background current with a significant time-dependence during the reconnection drive phase.

A decrease in the reconnection rate with increasing $B_g/B_{xy}$ is also observed in experiments and simulations with Hall reconnection. There is no consensus on the physical interpretation of this trend, but one explanation is that the reconnection flow is counteracted by an additional $j \times B$-force due to the interaction of the guide field and the Hall current, thus reducing the reconnection rate [12,129]. As previously discussed, two-fluid effects are here found to be negligible and should not have a significant influence on the reconnection rate. Since the observed reconnection rate depends on $dI_p/dt$, it is affected by similar factors as the plasma current. As discussed in section 3.4 and section 5.7 the current flowing in response to the inductive electric field is limited by electrostatic fields due to the axial sheath boundary conditions. This explains why the inductive electric field with a plasma, and thus the reconnection rate, at most exhibits a reduction of twenty percent from the vacuum case. As with the plasma current, it is difficult to determine the exact cause of the observed guide field dependence of the reconnection rate. However, based on the importance of the sheath boundary conditions on the amplitude of the plasma current, the reconnection rate can be expected to be sensitive to any changes in the boundary conditions. The amplitude of the plasma current, and thereby the reconnection rate, is also influenced by the guide field dependence of the plasma gun discharge. In addition, the plasma resistivity is changed by varying plasma parameters. Because of the lack of detailed information of how exactly both are influenced by the guide field, it is not possible to determine their relative importance for the reconnection rate.
Summary and conclusions

VINETA II is a device dedicated to the study of magnetic reconnection, with the purpose of gaining insight into the fundamental physical processes involved. Magnetic reconnection is governed by the interaction of the large-scale global dynamics and the small-scale dynamics within the current sheet. The global dynamics of magnetic reconnection was investigated within the framework of this thesis. It includes the characterization of the magnetic field topology, the electric fields, and the current flowing in the plasma. Guide field magnetic reconnection occurs in many space plasmas and experiments, and particular emphasis was put on how the guide field component influences the reconnection dynamics. Special attention was also devoted to the consequences of a plasma profile with steep gradients and the influence of the axial boundary conditions.

The study of reconnection in an experimental device necessitates a means for driving the reconnection. This was achieved either by an oscillating current through two parallel conductors (setup A) similar to the UCLA reconnection experiment [130], or by modifying a stationary X-point topology by an oscillating current through an additional conductor pair (setup B) similar to the VTF [15] experiment. The driving methods utilized in VINETA II offer high reproducibility and allow for the spatial distribution of the relevant parameters to be obtained from scanning probes on a shot-by-shot basis.

Setup A has the disadvantage that the amplitude of the in-plane magnetic field changes
6 Summary and conclusions

significantly during the drive cycle. For some phases of the drive, the plasma current has the dominating contribution to the field topology and the neutral sheet of the in-plane magnetic field forms three minima or even results in an O-shaped topology. Similar results were reported in ref. [122]. Setup B has the advantage that the X-point topology remains throughout the entire drive cycle and allows for the inductive field to be set independently of the amplitude of the magnetic field. This driving method was used for the majority of the results presented.

A linear device offers a simple geometry with a uniform guide field and well-defined axial and radial boundary conditions. However, the open end boundary conditions limit the inductively driven axial current, as well as the achievable plasma temperature due to the lack of axial confinement. The sheath limitation of the current is a well known issue and was also reported in ref. [126]. It necessitates an additional source of electrons in order to provide the current in response to the inductive electric field. In VINETA II, this is achieved using a plasma gun. Though originally taken as a source of additional electrons, the plasma gun was found to also generate a high density plasma, which establishes a collision-dominated operation regime. Due to a collision frequency much larger than the gyrofrequency of the ions, only the electrons are magnetized. The ions were also unmagnetized in the UCLA experiments [122], where the small amplitude of the magnetic field led to an ion gyroradius larger than the characteristic length scale of the plasma.

The main findings are:

i) Axial electric fields

Electric fields play a crucial role in the reconnection dynamics. Through the $E \times B$-drift, plasma is transported into the diffusion region and the electric field at the X-point determines the rate of flux transfer across the separatrix. The electric fields can be described by the generalized Ohm’s law. For magnetic reconnection in the collisional regime, the resistivity term dominates, while in the collisionless regime other non-ideal MHD terms such as the Hall term, the pressure gradient term, and the electron inertia term have to be considered. For the collision dominated plasma in VINETA II, the electron inertia term is negligible, and though both the Hall term and the pressure term are significant, the sum of the two is small. This is because the Hall term is linked to the diamagnetic current
via the pressure gradient. The axial plasma current flows in response to a combination of electrostatic fields and the externally induced electric field of the drive. This becomes clear if one considers the background current, which is present without the drive, and from the low amplitude of the measured current density. The current density is higher than the upper sheath limit, but its magnitude at the X-point is much lower than expected from $j_z = E_{\text{ind}}/\eta_s$. Even if this simple expression is extended to the parallel generalized Ohm’s law at the X-point, which takes the pressure gradient term and the inertial term into account, the low magnitude of the current density cannot be explained. It is concluded that the gun does not provide a sufficient amount of additional electrons and that there is consequently an electrostatic field opposing the inductive electric field. A significant axial electrostatic field was also reported in ref. [126]. If the role of the electrostatic field is ignored, other effects such as anomalous resistivity are overestimated. This is particularly important for linear devices but is also valid for other configurations.

While the axial inductive electric field does not reverse its direction within the azimuthal plane, the total electric field does. This is evident from the reverse currents observed at the edge of the current sheet. A reverse current was also reported in ref. [123], where the electric field reversal was attributed to outflow plasma jets with velocity $v$ in a strong transverse magnetic field $B$ giving rise to an additional electric field $E' = v \times B$. In VINETA II, the reverse current is instead attributed to a combination of the electrostatic field associated with the background current and the inductive electric field during the phases when the two fields are oppositely directed. The existence of a reverse current makes it possible to have a current sheet even if the total current is zero.

ii) Current sheet geometry

In simplified two-dimensional reconnection models, such as the Sweet-Parker model, current flows perpendicular to the reconnection plane tend to be located in the region of the X-point where the magnetic field is zero. For magnetic reconnection with a significant guide field component, the reconnection electric field has a component parallel to the magnetic field, and the current location is rather determined by global features [131]. If two-fluid effects are important, a $j \times B$-force due to the interaction of the in-plane current and the guide field leads to a tilt of the current sheet [12, 68]. A tilted current sheet is also observed in the present experiment and the current sheet position and geometry turns out to be determined by the current source, magnetic mapping and cross-field
6 Summary and conclusions

plasma diffusion.

The circular cross-section of the current sheet established by the background current is set by the aperture of the plasma gun and the cross-field diffusion of the plasma. Consequently there is an increase in area and decrease in the current density amplitude with decreasing guide field strength and increasing axial distance from the plasma gun. This current sheet is modified by the addition of the X-point topology of the in-plane magnetic field and the axial inductive electric field. Due to the field-aligned current, the magnetic X-point topology deforms the current sheet through magnetic mapping. It is elongated along the separatrix in one direction and compressed in the other, resulting in a tilt. The extent of the deformation depends on the ratio of the guide field to the in-plane magnetic field and the distance from the plasma gun. For a large guide field ($B_g/B_{xy} = 45$) the distortion is small and the cross-section of the current sheet retains its circular shape. For a smaller guide field ($B_g/B_{xy} = 3$) the current sheet is significantly elongated and follows the S-shape of the separatrix.

The inductive electric field of the reconnection drive increases the current density, but the shape remains the same and the deformation of the current sheet along the separatrix persist throughout the entire reconnection cycle. It is important to note that the current sheet is not only the result of the self-consistent electric and magnetic fields, but strongly depends on global boundary conditions and inhomogeneities.

iii) In-plane particle drifts and currents

In addition to axial currents, in-plane currents are also observed. For most measurements ($B_g/B_{xy} > 3$) these result in an out-of-plane magnetic field which peaks at the same position as the axial current density. However, for a ratio of the guide and in-plane magnetic field of $B_g/B_{xy} = 3$, the currents result in a quadrupolar magnetic field structure. The presence of a quadrupolar magnetic field is often taken as a signature of reconnection with two-fluid effects [60], but this is not the case here: The measured magnetic field structure is different from what is expected from Hall currents, i.e., the peak values are not located along the separatrix. Furthermore, from simulations [69] and experiments [67], the Hall-quadrupolar field is expected to be strongly reduced or even vanished when the amplitude of the guide magnetic field is on the same order as the reconnection magnetic field. From the relevant in-plane particle drifts, it is concluded that
the in-plane current instead originates from the combination of the electron diamagnetic current and the component of the field-aligned current in the azimuthal plane. The latter is the cause for the quadrupolar structure of the magnetic field. These in-plane currents are also present even without reconnection. The main effect of the inductive field and the reconnection drive on the in-plane currents is to increase the amplitude of the axial current density and thereby steepening the gradients of the plasma pressure and the plasma potential.

iv) Guide field dependence of the axial plasma current and the reconnection rate

The peak amplitude of the axial current current changes with the guide field magnitude. There is a decrease in the background current with increasing guide field. The total current is also decreased by the X-point topology, with the largest reduction at small guide fields and only a small reduction for $B_g \geq 18 \text{ mT}$. This is attributed in part to an elongation of the current sheet past the Rogowski coil and grid 2 due to magnetic mapping and cross-field plasma diffusion, and in part to a reduction in the current because of changing axial boundary conditions and plasma parameters. There is only a relatively weak guide field dependence of the contribution of the inductive electric field to the total current, with a small increase in magnitude at small guide fields.

The reconnection rate is defined by the axial inductive electric field $E_{\text{ind}}$ at the X-point, i.e., the rate at which the magnetic flux is transferred across the separatrix from the inflow region to the outflow region. In order to compare how the reconnection rate changes with the guide field, it is normalized to the vacuum inductive electric field. The difference in the electric field $\Delta E$ between the vacuum reconnection case and the case with a plasma is largest when the time-derivative of the axial plasma current has a peak. This means that there is a $90^\circ$ phase shift between the plasma current and $\Delta E$. Similar results where reported in the VTF experiment, where the inductive electric field was observed to increase by a change in $dj/dt$ [125]. The dependence on the reactance of the global electrical circuit is seen from the phase-shift of $\theta = 21^\circ$ between the peak of $E_{\text{ind}}$ and the peak of the plasma current. While this is of obvious importance in linear geometry, where the largest contribution to the reactance comes from the inductance of the external circuit, it is also relevant for toroidal devices. For VTF, for example, the self-inductance of the toroidal current channel results in a systematic delay between the reconnection drive electric field and the reconnection rate at the X-line [132].
Since the inductively driven current is limited by the axial sheath boundaries, the extent to which the vacuum electric field can be altered is also limited. The normalized reconnection rate is consequently always higher than $E_{\text{ind},p}/E_{\text{ind},v} = 0.85$. It is found to be highest when the magnetic field ratio is small and almost reaches the vacuum reconnection rate for $B_{\parallel}/B_{xy} = 3$ with a rate of $E_{\text{ind},p}/E_{\text{ind},v} = 0.95$. As the guide field is increased, the reconnection rate first decreases to approximately $E_{\text{ind},p}/E_{\text{ind},v} = 0.87$ at $B_{\parallel}/B_{xy} = 9$ and then only decreases slightly further for larger field ratios, reaching approximately $E_{\text{ind},p}/E_{\text{ind},v} = 0.85$ at $B_{\parallel}/B_{xy} = 60$. A decrease in reconnection rate with increasing $B_{\parallel}/B_{xy}$ is also observed in experiments and simulations with Hall reconnection. There is no consensus on the physical interpretation of this observation, but one explanation is that a $\mathbf{j} \times \mathbf{B}$-force of the in-plane current and the guide field opposes the reconnection flow and thereby reduces the reconnection rate [12]. Two-fluid effects are negligible in the presented work. Instead, the dependence of the reconnection rate on the guide field strength is attributed to a combination of the sheath limitation of the current, changes in the boundary conditions (e.g. through the change in the current sheet geometry by cross-field plasma diffusion and magnetic mapping), and the effect of the guide field on the actual plasma gun discharge and the plasma resistivity through changes in the plasma parameters.
Bibliography


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