Two-Dimensional Heat Transport in Tokamak Plasmas

PROEFSCHRIFT

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Contents

1 Introduction 1
  1.1 Why fusion? ........................................... 1
  1.2 The new challenge in fusion research: burn control .......... 3
    1.2.1 Understanding transport with the aim to control it .... 3
    1.2.2 Two-dimensional transport .......................... 3
  1.3 List of publications .................................. 6
    1.3.1 Journal publications ............................... 7
References ............................................... 9

2 1D and 2D transport in relation to burn control 11
  2.1 Introduction ......................................... 11
  2.2 Transport in a tokamak .............................. 11
  2.3 The standard case: 1D transport ...................... 13
    2.3.1 Classical, neoclassical and turbulent transport ...... 13
    2.3.2 Anomalous (turbulent) transport ................... 16
    2.3.3 Scaling laws ........................................ 16
    2.3.4 Profile consistency ................................. 17
  2.4 Magnetic island transport effects and stability .......... 17
    2.4.1 Magnetic island topology .......................... 17
    2.4.2 Effects of island geometry on transport ............ 18
    2.4.3 Magnetic island stability ........................... 19
  2.5 Conclusion ........................................... 23
References ............................................... 24

3 The TEXTOR tokamak, dynamic ergodic divertor, heating system and diagnostics 27
### References

#### 4 The Electron Cyclotron Emission diagnostic

4.1 Introduction .......................... 41
4.2 Electron Cyclotron Emission .................. 42
4.3 Electron Cyclotron Wave propagation ............. 44
4.4 Heterodyne ECE detection ................. 46
4.5 The 2D ECE-Imaging system ............... 48
4.5.1 ECEI optics ........................ 48
4.5.2 ECEI electronics ...................... 50
4.6 ECE noise-limits to resolution and signal enhancement ...... 52
4.6.1 Fourier transform filtering ............. 53
4.6.2 Singular value decomposition ............ 54

References .................................. 57

#### 5 Pressure profile consistency in TEXTOR

5.1 Introduction .......................... 59
5.2 Profile models .......................... 61
5.2.1 Critical Temperature gradient transport .......... 61
5.2.2 Profile consistency ..................... 62
5.2.3 Coincidence between profile consistency and temperature profile stiffness .......... 65
5.3 Comparison between models and experiments ......... 68
5.3.1 Experimental tests of model validity .......... 68
5.3.2 Experimental setup ...................... 68
5.4 Profile consistency and profile stiffness features from TEXTOR plasmas ........................................ 70
5.4.1 Central heated plasmas .................................. 70
5.4.2 Off-axis heating ........................................... 74
5.5 Summary and conclusions ................................ 77
References ..................................................... 79

6 Heat pulse propagation studies around magnetic islands induced by the Dynamic Ergodic Divertor in TEXTOR .......................................................... 83
6.1 Introduction .................................................. 84
6.2 Experimental setup ........................................ 85
6.3 Derivation of $\chi_e$ inside and outside the island ........ 88
6.4 Heat pulse propagation in and around a large $m/n = 2/1$ island 90
   6.4.1 Experimental results $m/n = 2/1$ island .............. 90
   6.4.2 Fourier analysis $m/n = 2/1$ island .................. 91
   6.4.3 Heat transport coefficients inside the island $m/n = 2/1$ island 95
6.5 Heat pulse propagation in and around an $m/n = 3/1$ island .... 96
   6.5.1 Experimental results $m/n = 3/1$ island .............. 96
   6.5.2 Fourier analysis $m/n = 3/1$ island .................. 98
   6.5.3 Heat transport coefficients inside and outside the $m/n = 3/1$ island 102
6.6 Discussion .................................................. 102
   6.6.1 Heat diffusivity in island .............................. 102
   6.6.2 Critical island width ................................. 104
   6.6.3 Conclusion ............................................. 105
References ..................................................... 106

7 Determining the sawtooth mode structure with 2D ECE-Imaging ........ 111
7.1 Introduction ................................................ 111
7.2 Sawtooth models .......................................... 113
   7.2.1 Sawtooth models bases on resistive internal kink instability 114
   7.2.2 Sawtooth model based on Quasi-interchange instability .... 117
7.3 Experimental setup ........................................ 118
   7.3.1 Noise reduction with Singular Value Decomposition .... 119
   7.3.2 Data normalization .................................... 120
7.4 Experimental results ....................................... 124
   7.4.1 General observations ................................. 124
<table>
<thead>
<tr>
<th>Section</th>
<th>Page</th>
</tr>
</thead>
<tbody>
<tr>
<td>7.4.2 Two-dimensional observations</td>
<td>125</td>
</tr>
<tr>
<td>7.5 Summary</td>
<td>130</td>
</tr>
<tr>
<td>References</td>
<td>132</td>
</tr>
<tr>
<td>8 Evaluation and outlook</td>
<td>137</td>
</tr>
<tr>
<td>8.1 Profile-consistency</td>
<td>138</td>
</tr>
<tr>
<td>8.2 Heat-pulse propagation in and around magnetic islands</td>
<td>138</td>
</tr>
<tr>
<td>8.3 Sawtooth oscillations</td>
<td>139</td>
</tr>
<tr>
<td>8.4 Outlook</td>
<td>140</td>
</tr>
<tr>
<td>References</td>
<td>142</td>
</tr>
<tr>
<td>Summary</td>
<td>143</td>
</tr>
<tr>
<td>Samenvatting</td>
<td>147</td>
</tr>
<tr>
<td>Curriculum Vitae</td>
<td>151</td>
</tr>
<tr>
<td>Acknowledgements</td>
<td>153</td>
</tr>
</tbody>
</table>
1

Introduction

1.1 Why fusion?

Energy is the basis of life. Energy is needed for a variety of technologies that greatly improve our lives such as lighting, heating, engines that allow fast and long distance traveling, food production, computers used for data processing.

Up to now humanity as a whole has consumed increasing amounts of energy as new technologies have been developed and the human population has grown. To support this increasing energy demand, sustainable energy sources need to be developed.

The current energy supply mix has some serious disadvantages. Most of the current energy (81% of the total) is derived from fossil fuels such as oil, coal and natural gas supplies [1]. There are various problems arising from the use of fossil fuels, such as the pollution associated with the mining and the burning, limited reserves and unequal distribution around the earth, which can cause conflict to secure the access. Furthermore the burning of fossil fuels produces $CO_2$ which is thought to influence climate change [2], which affects all of mankind.

Alternatives to fossil fuels that are sustainable are hydro (currently 2.2% of the world energy production), combustible renewables or bio-mass (9.8%, of which a large fraction is 'traditional' (wood fires) and waste incineration), wind and solar etc. Although solar and wind (currently 0.7%) have strong growth potential, it will be challenging to completely replace the fossil fuels since renewables such as solar and wind require large areas of land and sea and thereby preclude other use of land such as agriculture.
Another alternative energy source is nuclear energy which derives its energy from the binding energy of the atoms. Nuclear reactions can produce energy by splitting of heavy atoms (nuclear fission) or fusion of light atoms (nuclear fusion). Nuclear fission presently provides 5.9% of the world’s energy supply, mostly from splitting uranium isotopes. The uranium supply on earth is, however, fundamentally limited. By breeding fissile isotopes the supply of fission energy at current energy consumption rates can, however, be extended to a few thousand years.

Nuclear fusion energy has since long been proposed as an 'endgame' sustainable energy source. A nuclear fusion reaction between light atoms such as the hydrogen isotopes deuterium (D) and tritium (T) yields large amounts of energy. The potential for the production of fusion energy is indeed enormous, and fusion is safe, clean, zero-$CO_2$ and essentially inexhaustible. However, fusion has one serious drawback: it is very difficult, the fundamental problem being that the fuel only burns at a temperature of hundreds of millions Kelvin.

Different confinement concepts are being developed to achieve an energy gain from fusion in a controlled experiment. The main lines are inertial confinement (compression of a little ball of fusion fuel by a powerful laser pulse) and magnetic confinement. The research described in this thesis falls in the realm of the magnetic confinement concept.

In a magnetic confinement machine, the fusion fuel is confined with strong magnetic fields. This is possible, since at the required temperatures, the fuel is completely ionized and forms a plasma. Research in the past 50 years has developed the magnetic confinement concept to the point that it is now well possible to confine and sustain a plasma at the required temperature of a few hundred million K. The next requirement, i.e. that the fusion power exceeds the power needed to sustain the plasma, primarily calls for an up-scaling of the reactor. Currently the largest fusion reactor, ITER, is being built in Cadarache (south of France). ITER, a world-wide scientific collaboration, will demonstrate 10-fold power multiplication at the $500MW$ level, during pulses of 10 minutes or more. The construction of ITER will take some 10 years. After the completion of ITER, a demonstration fusion power plant (DEMO) will be built - connected to the grid - and after that the first generation of commercial plants can be expected.
1.2 The new challenge in fusion research: burn control

As has been shown, fusion research has come to the point that a power-producing reactor is being built: ITER. ITER is a tokamak [3], a device in which a hot plasma is confined by magnetic fields in a toroidal geometry (see figure 1.1). ITER will be the first reactor to produce a burning plasma, i.e. a plasma that sustains its high temperature with the power released in fusion reactions. This fact brings a whole new set of physics challenges, with ‘burn control’ as the central issue. Where previously it was sufficient to understand the transport mechanisms in a plasma so that the performance of new devices could be predicted, in a burning plasma, active, real time control of the burn process will be necessary. Whereas in present fusion reactors control of the pressure of the fuel – which is the main factor determining the burn rate – is achieved by means of the external heating systems, in a burning plasma this control tool is not available anymore. In the absence of means to control the heat source, we are looking at ways of controlling the heat loss. The occurrence of magneto-hydrodynamic (MHD) modes offers a possibility to do this, as these modes increase the heat loss. Therefore, control of MHD-modes is an important approach in the wider field of burn control.

1.2.1 Understanding transport with the aim to control it

To understand transport in a tokamak, it is necessary to have a clear picture of the topology of the confining magnetic field. This consists of two components. First: a strong toroidal field that is generated by external coils. Second: a poloidal field generated by the electric current that is run through the plasma in toroidal direction. In the ideal case, the field lines of the combined, helical field, lie on nested, toroidal surfaces. These are called flux surfaces (because the poloidal magnetic flux is constant on these surfaces). This standard picture of the topology of the magnetic field in a tokamak is sketched in figure 1.1.

1.2.2 Two-dimensional transport

In a steady state, the flux surfaces are isobaric, because there can be no pressure gradient along the field. Moreover, because thermal transport along field lines is so fast (up to 14 orders of magnitude faster than across the field) the flux surfaces are also isothermals. Thus, if the magnetic topology has the ideal form of nested toroidal flux surfaces, transport can be described as a one-dimensional (1D)
Figure 1.1: Tearing of magnetic surfaces causes the formation of magnetic islands. Shown in the figure is the formation of islands. \( q = \frac{m}{n} = \{1/1, 2/1, 3/1\} \). The magnetic field lines inside of an island form nested magnetic surfaces. The \( m \)-number corresponds to the number of crescent shaped islands in the poloidal cross-section. The center of the magnetic surfaces inside the island is a single field line called the O-point.

process: from surface to surface. This is the standard assumption for transport in a tokamak. However, there are exceptions. First, if the magnetic topology is perturbed, flux surfaces degenerate to form magnetic islands (figure 1.1). Now transport is truly two-dimensional. Second, during so-called MHD instabilities, fast plasma movements may occur, in which case the static pressure on a flux surface no longer needs to be uniform. This thesis concentrates on both these cases of two-dimensional (2D) transport, because both are of great importance for burn control. Magnetic islands enhance transport and, if they grow in an uncontrolled manner, can even lead to a total loss of pressure. Therefore, in ITER a system will be installed that can control the islands. This system either heats the island by precisely localized heat deposition, or by equally precise generation of toroidal current density just in- or outside the island. Both approaches use high-power \( mm \)-wave technology. Clearly, for this control system it is essential to understand the 2D transport in and around a magnetic island. The MHD-instability we focus on is the so-called sawtooth instability. This instability, which occurs in all standard tokamak operation modes, quasi-periodically redistributes the pressure in the central part of the plasma. It can be both beneficial - by
1.2 The new challenge in fusion research: burn control

preventing a build-up of helium (the ash of the fusion reaction) in the center of
the plasma - and detrimental: a large sawtooth instability can induce magnetic
islands that subsequently must be kept in check by the method outlined above.
For both reasons, the size and frequency of the sawtooth instability occurrence
must be controlled. Similar to the control of magnetic islands, high-power $mm$-
wave technology is also employed here. However, the actual instability, i.e. the
way in which the hot core of the plasma is redistributed in a fast MHD event, is
not unique. There are several forms with different 2D flow patterns association
with them. The way the $mm$-wave power must be employed in order to control the
instability depends on the form of the instability. Therefore, a good insight in the
2D transport process during the sawtooth instability is essential for a successful
control strategy.

In summary, the focus of this thesis is on two different cases of 2D transport in
a tokamak reactor. Both are important for burn control. Prior to the 2D transport
studies, 1D transport in the unperturbed magnetic topology is considered. This is
governed by turbulent processes, which leads to a non-linear relation between e.g.
the heat flux and the temperature gradient. This non-linearity can give rise to so-
called 'profile stiffness': the higher a gradient, the more power is needed to further
increase it. Several models have been suggested that propose a constraint on the
profiles that goes further than the profile stiffness. These models, indicated by the
generic term 'profile consistency', invoke an overarching principle that determines
the pressure or current density profile, to which the turbulent transport adjusts.
Clearly, if transport is governed by such a global constraint rather than being
locally determined, this must be taken into account in any control strategy that
aims to control the profiles.

The study of these 1D and 2D transport phenomena was only possible by
virtue of the combination of several experimental tools:

- The TEXTOR tokamak, which is equipped with a mix of external heating
tools including the high-power $mm$-waves used for the control of islands and
the sawtooth instability by means of electron cyclotron resonance heating
or current drive (ECRH/ECCD).

- The Dynamic Ergodic Divertor, a set of external magnetic perturbation coils
that can induce magnetic islands.

- A unique 2D ECE-Imaging system, that provides 2D measurements of the
plasma temperature with a high time resolution.
An advanced Thomson scattering system, providing accurate and high-resolution measurements of the pressure and temperature profiles.

The main problem addressed in this thesis is the ability to control transport in order to control the fusion burn process in a tokamak reactor. Since MHD modes affect the transport in a fusion plasma, the control of MHD modes provides a way to control the transport in a fusion plasma. The most eminent MHD modes that influence the transport in the plasma are the sawtooth ($m = 1$) instability and magnetic island $\{m = 2, 3, \ldots\}$. For the control of the stability of these modes ECRH/ECCD is applied. The efficiency of using this control method depends on the transport in the vicinity of these modes. Here it is aimed to identify how these instability modes influence the transport with the aim to control them. For this the role of transport due to instabilities is investigated for different plasma magnetic configurations:

- In a 1D magnetic configuration it is investigated whether the effect of profile shape conservation under localized heating (i.e. profile stiffness) is due to a local change in transport coefficients or if it is governed by non-local transport (profile consistency). (chapter 5)

- The 2D transport inside and around static magnetic islands is investigated with perturbative transport analysis to address the question if inside of the island the profiles also exhibit stiffness under localized heating. (chapter 6)

- The 2D transport during sawtooth crash in the plasma center is determined with the aim of identification of sawtooth instability which can lead to a stabilization strategy. (chapter 7)

Before reporting the transport studies an overview of the experimental setup at the TEXTOR tokamak and diagnostics used is given in chapter 3. In chapter 4 an extensive description the ECE-imaging system is given which is the central diagnostic for the 2D transport studies. An overview of the results and an evaluation and an outlook are provided at the end.

### 1.3 List of publications

Below a list of publications related to the work described in this thesis is given:
1.3.1 Journal publications

‘Heat pulse propagation studies around magnetic islands induced by the Dynamic Ergodic Divertor in TEXTOR’,
_Nuclear Fusion_ **48**, 115005 (2008)

‘Transport of argon and iron during a resonant magnetic perturbation at TEXTOR-DED’,
_Plasma Physics and Controlled Fusion_ **51**, 032001 (2009)

‘Particle Confinement Control with External Resonant Magnetic Perturbations at TEXTOR’,

‘Identification and analysis of transport domains in the stochastic boundary of TEXTOR-DED for different mode spectra’,
_Nuclear Fusion_ **48**, 024009 (2008)

‘The Effect of Heating on the Suppression of Tearing Modes in Tokamaks’,
Chapter 1. Introduction

‘Imaging meso-scale structures in TEXTOR with 2D-ECE’,
Plasma and Fusion Research 2 S1031 (2007)

'Self-consistency of pressure profiles in tokamaks',

'Link between self-consistent pressure profiles and electron internal transport barriers in tokamaks',

'The influence of resonant magnetic perturbations on edge transport in limiter H-mode plasmas in TEXTOR'

B.A. Hennen, E. Westerhof, J.W. Oosterbeek, P.W.J.M. Nuij, D. De Lazzari, G.W. Spakman, M. de Baar, M. Steinbuch and the TEXTOR team
'A closed-loop control system for stabilization of MHD events on TEXTOR',
Fusion Engineering and Design 84 928-934 (2009)

'Tokamak plasma self-organization and the possibility to have the peaked density profile in ITER',
Nuclear Fusion 49 065011 (2009)
References


2

1D and 2D transport in relation to burn control

2.1 Introduction

This thesis concentrates on aspects of 2D transport in tokamak plasmas, relevant to burn control. The goal and central questions of the research project are phrased in chapter 1. Here, a general introduction is given to some aspects of transport and magnetic islands that are relevant for the work in this thesis. It is by no means an exhaustive review, but may serve as background information for the chapters to follow. A review is given of the basic 1D transport, profile stiffness and profile consistency, instances of 2D transport and its importance for burn control.

2.2 Transport in a tokamak

Control of transport of heat and particles in a tokamak is important to achieve and maintain plasma conditions that are optimized for energy generation. The goal is not simply to minimize transport so as to maximise confinement. If the confinement is too good, the concentration of helium - the ‘ash’ of the reaction – will become too high and this will choke the reaction. Hence we need to minimize thermal losses while maintaining sufficient particle transport. Furthermore, control of transport is needed to optimize the pressure profile, which allows the reactor to be run close to the stability limit of magnetic confinement in a controlled manner.

The cross-field transport is dominated by turbulent processes. This holds in particular for the thermal transport via the electrons, and for the energetic alpha-
particles that are generated in the fusion reactions. These must be confined long enough to thermalize and thus transfer their energy to the bulk of the plasma, yet not so long that they choke the reaction.

In present fusion reactors, even in the largest such as the Joint European Torus (JET), the contribution of the alpha-particles to the heating of the plasma is negligible. External heating is therefore the dominant power input, and this provides the operator with a direct handle on the plasma pressure. To give a practical example: when the plasma pressure starts an uncontrolled rise due to the formation of a transport barrier, the operator can turn down the heating to avoid a disruption. Moreover, the external heating methods - in particular the Neutral Beam Injection - provide a source of toroidal momentum, which allows the operator to exert some control over the flow velocity in the plasma, which in turn affects the turbulence. And finally, with external heating methods non-inductive electrical currents can be driven, which again provides the operator with a tool to act on the magnetic equilibrium.

Clearly, in a burning plasma the bulk of the heating comes from the fusion reaction itself, so that the possibilities for control through external heating are strongly reduced. On top of that, the alpha-particle heating interacts both with the pressure (the power being proportional to the pressure squared) and the turbulence. The newly arisen interactions are schematically indicated by the red arrows in figure 2.1. The energetic alpha particles population will affect the stability of the MHD modes and turbulence which on its own will affect the population of the energetic alpha particle population. The energetic alpha particle population will also constitute a large part of the plasma energy.

As a result of the dominance of internal heating and the new, non-linear relations between pressure, alpha-particles and MHD modes and turbulence, new control schemes are needed. These will rely primarily on the control of MHD modes by local heating and current drive, where it should be noted that the localization of externally injected power in a narrow region allows this method to have a large impact at a power level that is small compared to the alpha-particle heating. The MHD modes will be controlled with high power localized mm-waves that couple to the plasma electrons in the vicinity of the MHD modes. The effect is to drive a current directly by the electron cyclotron current drive (ECCD) and/or drive a current inductively through the creation of a local temperature increase by electron cyclotron resonance heating (ECRH). The efficiency of this process depends on the transport properties in the vicinity of the modes. Because transport is important for the stabilization of MHD modes, a review of 1D transport and
Figure 2.1: Schematic diagram of the interactions between fast alpha-particles generated in fusion reactions and the magneto-hydrodynamic modes. The black arrows indicate interactions that are already acting in present experiments, and that are reasonably well understood. New in a burning plasma is the fact that the alpha particles are generated in such large quantities, that they directly affect the plasma pressure profile [adapted from application FOM-programme]

is given first. Associated with certain MHD modes are magnetic islands whose effect on transport is described in the following paragraphs.

2.3 The standard case: 1D transport

2.3.1 Classical, neoclassical and turbulent transport

As introduced in chapter 1, an idealized tokamak equilibrium configuration consists of toroidal symmetric magnetic surfaces that are isobaric and isothermal. Conservation of energy and particles requires that changes in local plasma parameters such as temperature ($T$) and density ($n$) are balanced by fluxes of heat and particles respectively. The local energy balance equation is

$$\frac{3}{2}\partial_t(nk_B T) + \nabla \cdot q = S_E$$

(2.1)
where the first term is the time derivative of the local energy density \((nk_B T = p, \text{ with } k_B \text{ is Boltzmann's constant})\), the second term the divergence of the heat flux \((q)\) and third term is energy-source density \((S_E)\) which constitutes local sources and sinks of heat. The particle balance equation has a similar form

\[
\partial_t n + \nabla \cdot \Gamma = S_p
\]

(2.2)

where the first term is the time derivative of the particle density, the second term is divergence of the particle flux \(\Gamma\) and the third term \(S_p\) is a particle source term. The fluxes depend on the gradients of different plasma parameters. These relations are conveniently expressed in terms of a transport matrix:

\[
\begin{bmatrix}
\Gamma \\
q_e \\
q_i \\
\phi
\end{bmatrix}
= -
\begin{bmatrix}
D & \ldots & \ldots & W \\
\ldots & \chi_e & \ldots & \ldots \\
\ldots & \ldots & \chi_i & \ldots \\
B & \ldots & \ldots & \sigma
\end{bmatrix}
\begin{bmatrix}
\nabla n \\
n\nabla T_e \\
n\nabla T_i \\
E_\phi
\end{bmatrix}
\]

(2.3)

As the transport coefficients themselves may depend on gradients, the relation between flux and gradient will in general not be a simple proportionality. In particular, turbulent transport itself is itself driven by gradients of e.g. temperature or pressure. This is illustrated in figure 2.2. The figure also makes clear how different ways of determining transport coefficients, i.e. from power balance or from a perturbation analysis, yield essentially different results. The power balance coefficient is defined as \(\chi_{pb} = -q/(n\nabla T)\). The perturbative transport coefficient is the partial derivative \(\chi_{pert} = -\partial q/\partial(n\nabla T)\). The first situation in figure 2.2 shows a proportionality of the heat-flux and the temperature gradient which yields \(\chi_{pert} = \chi_{pb}\). The second situation shows linear dependence with an off-set which yields a single \(\chi_{pert}\) and a \(\chi_{pb}\) that varies with the temperature gradient. This situation is encountered with off-diagonal elements in the transport matrix and results in a non-zero heat flux for zero temperature gradient. The third shows a nonlinear dependence, in which case both \(\chi_{pert}\) and \(\chi_{pb}\) are dependent on the temperature gradient. The fourth cartoon illustrates the case of a heat flux that strongly increases above a critical gradient. In such a case the concept of a heat diffusivity practically loses its meaning as the determining parameter now is the critical gradient. A brief overview of the different transport mechanisms and the transport fluxes they yield is given below.
2.3 The standard case: 1D transport

![Diagram showing possible behaviors of radial heat flux]

Figure 2.2: Schematic representation of possible behaviors of radial heat flux as a function of the temperature gradient.

**Classical transport**

In the absence of instabilities and turbulence, the absolute minimal level of transport level is due to Coulomb collisions between charged particles. In the classical picture the transport has been calculated as a random walk process. The particles make an average step size equal to the Larmor radius during an average collision time which leads to a classical diffusion coefficient

\[ D_{\text{class}} \propto \frac{\rho^2}{\tau}. \]  

(2.4)

**Neoclassical transport**

In toroidal geometry of the tokamak, the orbits of the particles are affected by the curvature and gradients of the magnetic field. As a consequence, the particles make a larger step size during a collision time, which causes an enhancement of the transport with respect to the classical transport in cylindrical geometry. Neoclassical theory also predicts off-diagonal terms in the transport matrix. The most significant are the ‘bootstrap current’ (indicated by \( B \) in equation 2.7, a toroidal current driven by the radial density gradient, which can constitute a large part of the inductively driven current; and the Ware-pinch (\( W \)) that causes an inward particle flux driven by the toroidal electric field.
2.3.2 Anomalous (turbulent) transport

Where the neoclassical theory gives a precise description of the transport phenomena due to collisions, experimental measurements commonly show that transport is much higher than the neoclassical value. In particular the cross-field thermal diffusivities exceed the neoclassical value by up to one ($\chi_i$) or even two ($\chi_e$) orders of magnitude. This enhancement of transport is due to turbulence. Typical experimental values for the cross-field diffusivities are

$$\chi_e \approx D_{neo} \approx 1 \text{ m}^2/\text{s}$$

(2.5)

The current status of turbulent transport research is that there is no first principles theory that gives a complete description of the turbulent transport processes in a tokamak plasma. Numerical models of turbulent transport are getting to the stage that they are useful tools for the interpretation of experiments, but their predictive power is still limited. The turbulence simulations are very CPU intensive, often taking weeks of computation time for the simulation of a steady state situation. Simulations that stretch out over transport time scales, in a dynamically evolving plasma, are still out of reach.

2.3.3 Scaling laws

If no detailed, local and time-dependent description of turbulent transport is needed, it is useful to consider global confinement. The most common measure of the quality of the confinement is the energy confinement time $\tau_E$, defined as the ratio of stored energy to the heating power. The $\tau_E$ is related to the transport coefficients by $\tau_E \approx a^2/\chi$, where some spatial average of $\chi$ is implied. Information on profile shapes is lost in this integral quantity, but the global energy confinement time is a very useful quantity to characterise the performance of a fusion reactor, in particular of one that is being designed. Extensive studies based on databases, which included date from many tokamak experiments of different size and shape, have yielded scaling laws which express $\tau_E$ in terms of the machine parameters. A well-known example is the ITER98H scaling law, which was developed in support of the design of ITER [8]:

$$\tau_E = 0.0365 I^{0.97} B^{-0.08} P^{-0.63} n^{-0.41} M^{0.20} R^{1.93} \epsilon^{0.23} \kappa^{0.67}$$

(2.6)

with $I[\text{MA}]$ the plasma current, $B[\text{T}]$ the magnetic field, $P[\text{MW}]$ additional heating power, $n[10^{19} \text{ m}^{-3}]$ the plasma density, $M[\text{AMU}]$ the atomic mass, $R[\text{m}]$ the plasma major radius and $\epsilon$ the aspect ratio.
2.3.4 Profile consistency

Finally, it was observed already early in tokamak research that the radial profiles of temperature, pressure and current density in many experiments appeared to be very similar, controlled mainly by the value of the magnetic winding number ($q$). Several theoretical ideas were proposed that invoked some global constraint as the cause of this self-similarity of profiles. These range from purely local 'critical gradient' models - in which the profiles fill out until they reach a locally prescribed critical value, above which turbulence kicks in - to non-local models in which a global constraint, e.g. a minimum energy state, governs the profile shape and turbulent transport is rather a vehicle that helps the profile assume its preferred shape. The possible role of such 'profile consistency' or 'profile stiffness' is the subject of chapter 5, where it is considered against the background of the requirement to control transport.

2.4 Magnetic island transport effects and stability

Magnetic islands affect the local transport properties of the plasma and thus modify the overall thermal confinement of the plasma. For the stability of the magnetic island the local transport properties are important. They play a role in the excitation threshold of a meta-stable magnetic island the neoclassical tearing mode (NTM), and in the suppression of magnetic islands with ECRH/ECCD.

2.4.1 Magnetic island topology

The unperturbed topology of the magnetic field in a tokamak is a set of nested magnetic surfaces centered on a single magnetic axis. The formation of magnetic islands causes a change of the magnetic topology at surfaces with a rational magnetic winding number $q = m/n$. Magnetic islands are due to a relatively small magnetic perturbation of the magnetic field, which is schematically shown in figure 2.3. In the left panel the helical field $B_h(r)$ is shown with respect to the magnetic field lines at the rational surface located at $r = r_s$. In the right panel a magnetic island is caused by the super position of a small periodic perturbation $B_r = \hat{B} \sin(m\theta - n\phi)$ that is resonant with the helical field at the rational surface. The magnetic island topology consists of magnetic surfaces that are centered on the island magnetic axis called the island O-point. The locally nested surfaces of the magnetic island are separated from the rest of the plasma by a separatrix.
Field lines on the separatrix either converge or diverge to the X-point, which is a singular field line. The field lines of both the X-point and the O-point have the mode-numbers of the resonant perturbation $m, n$. The island width is related to the plasma parameters of the unperturbed equilibrium and the amplitude of the radial field magnetic field perturbation: $w = 4(\dot{B}_r/msB)^{1/2}$, where $s = d\ln(q)/d\ln(r)$ is the shear of the winding number $q$.

![Figure 2.3: Schematic representation of the island magnetic field topology.](image)

In the left panel the helical field with respect to the rational surface is shown as a function of the minor radius and the helical coordinate $\theta - \phi n/m$. A magnetic island causes a small field perturbation periodic perturbation in the radial direction. The island structure consists of nested magnetic surfaces centered on the island O-point. The island is bounded by a separatrix. The field lines at the separatrix converge and diverge at the island X-point.

### 2.4.2 Effects of island geometry on transport

The transport coefficients parallel to the magnetic field are many - typically more than 12 in the hot core! - orders of magnitude larger than those perpendicular to the magnetic field. As a result, temperature and pressure can be assumed to be uniform on the magnetic surfaces in nearly all situations, exceptions being very fast transients, surfaces with divergent winding number (such as the separatrix of a magnetic island) and the outermost layer of the plasma, where the low temperature brings the parallel and perpendicular transport closer together. For diffusive transport the characteristic parallel and perpendicular transport times between two points on a flux surface are expressed as

$$\tau_\parallel \approx \frac{L_\parallel^2}{\chi_\parallel}, \tau_\perp \approx \frac{L_\perp^2}{\chi_\perp}$$

(2.7)
where $L$ is the connection length and $\chi$ the diffusion coefficient. The parameters can equilibrate on the magnetic surface if $\tau_{||} > \tau_{\perp}$. In the vicinity of the magnetic separatrix of the island the parallel connection length $L_{||}$ diverges, so that insurface gradients parameters such as temperature and density occur [3, 7, 11]. For small islands the region of incomplete profile flattening extends over the entire island width. For islands larger than a critical island width, the profiles can be uniform on the magnetic surfaces inside the island. For small islands estimations of the connection lengths for transport across the island O-point are: ($L_{\perp} \leq w$) for the perpendicular connection length and $L_{||} \geq 2\pi R_0 q / q' n w$ the parallel connection length. These length scales yield a critical island width $w_c = (8 R_0 q / q' n)^{1/2} (\chi_{\perp} / \chi_{||})^{1/4}$ above which the temperature is uniform on the magnetic surfaces near the island O-point [3, 11]. The main effect of large islands ($w > w_c$) is that the transport across the island takes place in a narrow layer in the vicinity of the island separatrix as depicted in figure 2.4. The transport in the layer is driven by gradients both parallel and perpendicular to the magnetic field. The island boundary thus provides a radial component to the fast parallel transport causing a diversion of heat past the island and a flattening of the profiles inside of the islands resulting in a reduction of the total confined plasma energy. As discussed in the following paragraph the flattening of the profiles also affects the stability of a particular instability (the neoclassical tearing mode) and the ability to control it with localized high-power $mm$-wave deposition.

2.4.3 Magnetic island stability

In this section different instability mechanisms relevant for the growth of magnetic islands in future reactors and for the experiments described in this thesis are reviewed. In particular, the instabilities responsible for the tearing mode and the instability modes associated with the sawtooth crash are treated.

Tearing mode instability

At magnetic surfaces with rational winding numbers $q$, islands can appear due to the tearing mode instability. The stability (growth rate) of tearing modes depends on gradients in the equilibrium current and resonant helical currents that can be caused by different mechanisms. A theoretical growth rate of the island width due
Figure 2.4: A schematic representation of the layer around the island separatrix where the transport across the island takes places (adapted from [3]). In this layer the total transport is due to both gradients parallel and gradients perpendicular to the magnetic field. Inside the layer the parameters temperature and density are not constant on the magnetic surfaces. The layer width is $w_c/w$ at the helical position of the island O-point and wc at the helical position of the X-point. For small island widths, smaller than the critical island width ($w < w_c$), the transport layer encompasses the entire island, which causes that gradients across the island remain.

The drive due to the tearing mode is provided by the Rutherford equation [4]:

$$0.82 \frac{\tau_R \, dw}{r_s \, dt} = r_s \Delta_0'(w) + M_{Bootstrap} + M_{ext} - M_{ECRH} - M_{ECCD} \quad (2.8)$$

where $\tau_R = \mu_0 r_s^2 / \eta(r_s)$ is the resistive time scale at the island radius. The first term on the right-hand side of the equation is the linear stability index $\Delta_0'(w)$. The second term describes the drive due to the perturbed bootstrap current which drives a neoclassical tearing mode (NTM) and which becomes important for large values of the local poloidal beta, $\beta_p = 2 \mu_0 p / B_p^2$ [3, 4]:

$$M_{Bootstrap} \approx \epsilon^{1/2} \frac{L_q}{L_p} \beta_p \frac{r_s}{w} \left[ \frac{w}{w^2 + w_c^2} \right] \quad (2.9)$$

where $L_p,q$ are the gradient length scales of the pressure and $q(r)$-profile respectively, $\epsilon = r_s/R$ is the inverse aspect ratio, $w_c$ provides the threshold as a
consequence of incomplete flattening of the pressure across the magnetic islands at small island widths due to transport effects [3]. The third term describes the drive of an external resonant magnetic perturbation field, which can be caused by resonant magnetic perturbations outside of the plasma [2]:

\[ M_{\text{ext}} = 2m \left( \frac{w_{\text{vac}}}{w} \right)^2 \cos(\Delta \xi) \]  

(2.10)

Here \( w_{\text{vac}} = 4(\hat{B}_r/msBq)^{1/4} \) is the width of the vacuum island and \( \Delta \xi \) the phase difference between the vacuum field and the actual island of the tearing mode. The vacuum magnetic island is calculated as the superposition of the external perturbation field and the unperturbed equilibrium field. For rotating islands this term is negligible since \( \Delta \xi \) is continuously varying. The last term represents the effects of the perturbed current due to localized current drive and heating inside of the magnetic island [14]:

\[ M_{\text{ECCD}} = C \frac{w^2}{w^2} \int_{r_s-w/2}^{r_s+w/2} dr \oint d\xi \delta j_{\text{CD}} \cos(m\xi) \]  

(2.11)

\[ M_{\text{ECRH}} = C \frac{j_{\text{sep}}}{w^2} \frac{T_e^{3/2}}{T_{\text{sep}}^{3/2}} \int_{r_s-w/2}^{r_s+w/2} dr \oint d\xi T_e^{3/2} \cos(m\xi) \]  

(2.12)

With ECCD the current \( \delta j_{\text{CD}} \) is driven non-inductively inside of the island which is stabilizing if the driven current is parallel to the plasma current (\( \delta j_{\text{CD}} > 0 \)) but destabilizing when driven opposite of the plasma direction. With ECRH an increase of the electron temperature \( T_e \) inside of the island with respect to the separatrix \( T_{e,\text{sep}} \) can be created which causes an inductively driven current (\( \delta j_{\text{ind}} > 0 \)) inside of the magnetic island that stabilizes the tearing mode. While pure ECRH can be applied for modifying the stability of tearing modes ECCD will always be accompanied by ECRH. This affects the efficiency by which modes can be controlled with ECCD. The stabilization of tearing mode with ECCD is enhanced by ECRH, while the destabilization of tearing modes with ECCD is reduced by ECRH. The relative importance of ECRH depends on the fraction of non-inductive current that can be driven. The inductive current that can be driven inside an island depends on the temperature peaking that can be achieved. This in turn depends on the local power-balance inside of the island. A low heat diffusivity \( \chi_e^{(PB)} = -q_e/(n_e \nabla T_e) \) will allow a large temperature peaking inside of the island. Expressions describing the efficiency theoretically are given in [10], [6]. In [1] a
heat diffusivity \( \chi_{\text{e}}^{(PB)} \approx 1m^2/s \) is experimentally found which is comparable to the ambient plasma. Based on these findings it is expected that in ITER up to 20% of the current driven inside a tearing mode with ECCD will be due to ECRH \[1\]. These prediction strongly depend on \( \chi_{\text{e}}^{(pb)} \) inside of the island but its exact dependence on transport inside of the magnetic island has not yet been fully experimentally addressed. In this thesis the electron heat transport properties inside of magnetic islands are treated in chapter 6.

**Sawtooth instability**

In the plasma center near the \( q = m/n = 1 \) radius an instability appears that is responsible for the periodic crash which is called the sawtooth-crash. The sawtooth crash appears in the form of a rapid redistribution of central plasma profiles of temperature and density \[12\]. The crash is caused by instability modes at the \( q = 1 \) radius. Different modes have been put forward (either based on theoretical models or experimental data) to be responsible for the sawtooth crash. Based on the shape of the \( q \)-profile two different classes of modes are distinguished. A schematic representation of the different modes is given in figure 2.5. With \( q_0 < 1 \) in the plasma center, a resistive internal kink instability can cause growth of a magnetic island which will occupy a large part of the core. With \( q \approx 1 \) in the plasma center, the magnetic in the plasma core can be deformed due to quasi-interchange motion.

![Figure 2.5: Deformation of the magnetic surfaces in the plasma core due to instabilities at \( q = 1 \). Depending on the shape of the \( q \)-profile at \( q = 1 \) two different instabilities can grow. If \( q_0 < 1 \) which implies a shear \( s_1 \), a kink mode can cause the growth of a magnetic island as shown in (b). If \( q_0 \approx 1 \) for a large part of the minor radius, the magnetic surfaces can become deformed due to a quasi-interchange mode as shown in (c).](image-url)
The resistive kink and the quasi interchange mode have different instability mechanisms. Control of the magnetic shear is the envisaged method to modify the stability of the sawtooth mode, although the stability depends on additional effects, like the fast particle content in the fusion core. The growth of the resistive kink is triggered by the shear at the $q = 1$ radius exceeding a critical value $s_1 = d \ln(q)/d \ln(r)|_{q=1} > s_{\text{crit}}$ [9].

The stability of the mode can be modified by control of the shear. The shear can be modified by driving current at the either side of the $q = 1$ with ECCD and ECRH. Increasing the shear will destabilize the kink mode and speeds up the sawtooth crash. Decreasing the shear will stabilize the growth of the kink modes and delay the sawtooth.

The quasi-interchange instability is driven by the pressure gradient [5]. The plasma core is instable to quasi-interchange motion if the central safety factor remains close to unity ($q_0 \approx 1$) and additionally there is low shear $s_1$ within the $q = 1$ radius. The low magnetic shear causes that the magnetic surfaces with $q = 1$ requires little energy to be deformed. At low pressure (low poloidal beta, $\beta_p$) the mode can then be driven unstable. The stability of this mode can be somewhat modified by the shape of the flux-surfaces which influences the average shear. To trigger the sawtooth instability, only a small change of the $q$-profile so that it reaches $q = 1$, is sufficient [13]. The stability of the quasi-interchange can be expected to be modifying the pressure profile with ECRH or modifying the $q$-profile with ECCD/ECRH.

\section{Conclusion}

The plasma confinement is affected by various instability modes that can cause the growth of magnetic islands. Driving currents with high power \textit{mm}-waves provides a way to modify the stability of different modes. The ability to drive a current depends on the transport properties of the plasma. But in turn the transport properties of the plasma can depend on the presence of modes like magnetic islands. The ability to control the overall transport properties in the plasma thus depends on the transport properties in the vicinity of MHD modes and islands.

With the aim of identifying how the plasma profiles depend on local heating, the general transport properties in a plasma without islands are investigated in chapter 5. In chapter 6 the 2D electron heat transport properties in the vicinity of a stationary magnetic island are investigated. The transient electron heat
redistribution due to the fast event of a sawtooth crash in chapter 7.

References


In this chapter the experimental equipment central to this thesis is briefly introduced, to provide the necessary background information for assessing the setup of the experiments.

3.1 The TEXTOR tokamak

The experiments are performed on the tokamak TEXTOR. The basic lay-out of TEXTOR is a conventional tokamak, but it has some special features that are essential for the experiments described in this thesis. TEXTOR (tokamak experiment for technology oriented research) is a medium sized limiter tokamak with a circular shaped cross-section [15]. TEXTOR has as a major radius $R_0 = 1.75m$ and a minor radius $a = 0.46m$. A photo of the TEXTOR exterior is shown in figure 3.1. The toroidal field ($B_t < 3T$) is provided by 16 coils. Inside of the vacuum vessel a toroidal plasma current is induced by the iron core transformer which consists of 6 yokes with the Ohmic coils wound around the central legs. The transformer induces a maximum toroidal plasma current ($I_p < 800kA$). The plasma discharge has a maximum duration of 10s. Also shown are a set of coils parallel to the vessel that provide a vertical field to compensate the hoop force and are used for plasma positioning.
Figure 3.1: Photo of the TEXTOR tokamak and a schematic representation of its $B$-field generating coils. Around the vacuum vessel 16 coils are wound that provide the toroidal field (four of them are visible). Parallel to the vessel 4 coils are wound that provide a vertical $B$-field (two are visible). The toroidal electric field driving the plasma current is generated by the transformer (blue). Inside of the vessel there 16 helically wound dynamic ergodic divertor (DED) coils (see next section).

3.2 The Dynamic Ergodic Divertor

For the generation of magnetic islands on TEXTOR a special tool is employed, the ‘Dynamic ergodic divertor’ (DED) [8]. With the DED the magnetic islands can be induced and positioned in a controlled way. This allows a systematic investigation of the influence of magnetic islands on transport. The DED consists of a set of helical magnetic perturbation coils that are wound on the inboard side of the TEXTOR tokamak. A schematic configuration of the DED is shown in figure 3.2. The effect of the coils is to generate a radial magnetic field $B_r$. The field can be characterized by its poloidal and toroidal mode numbers $m$ and $n$. The perturbing field is resonant at magnetic surfaces where the normalized magnetic winding number $q$ equals $m/n$. The 16 DED coils can be connected to their power supply in different configurations which yield different mode spectra. Figure 3.2a shows a schematic picture of the so called 3/1 mode DED coil configuration, which is used in this thesis. In this configuration the coils are connected in 4 groups of 4 adjacent coils. The amplitude of the mode-spectrum $B_{r,mn}$ of this configuration at the location of the $q = 3$ flux surface is shown in figure 3.2b. In the 3/1 configuration, the dominant contributions to the radial field have mode numbers $m = \{1, 2, 3, 4, 5, 6\}$. The dominant toroidal mode number is $n = 1$, modes with $n > 1$ have negligible amplitude. The main resonances of the DED
3.2 The Dynamic Ergodic Divertor

on the plasma are thus expected at rational surfaces \( q = m/n = m < 6 \). Other operating modes are the 6/2 DED mode where the mode spectrum has a dominant \( n = 2 \) mode number and the 12/4 DED mode where the spectrum has a dominant \( n = 4 \) mode number. In these modes the perturbation field penetrates less deep into the plasma than in the 3/1 mode. The DED can be operated in a static DC mode as well as a rotating AC-mode. In the DC mode the perturbation field(266,627),(361,867)(373,627),(466,867)(480,627),(574,867) is static with respect to the tokamak vessel. A unique feature of the DED operation, however, is the AC-mode, which yields a perturbation field that rotates with respect to the vessel. Two opposite rotation directions can be employed \( \{ AC^+, AC^- \} \). The maximum \( I_{DED} \) depends on the mode of operation. In the 3/1 DED configuration the maximum coil current is \( I_{DED} = 3.75kA \). The spectrum of the DED perturbation field is characterized by both amplitude and a phase. For experiments described in chapter 6 it is important to know the positions of islands that are locked to the DED perturbation field. For large islands the island positions are predicted to coincide with those in the vacuum magnetic field \[9\]. The vacuum field is the superposition of the DED perturbation field and the unperturbed equilibrium field. It does not incorporate the plasma response to the perturbation. The structure of the vacuum magnetic field is well visualized in a Poincaré–map that shows the intersection of trajectories along the magnetic field.

![Figure 3.2](image_url)

**Figure 3.2:** a: schematic representation of the DED coil configuration in its 3/1 configuration \[7\]; b: the amplitude spectrum of the radial field \( B_r(\theta) \) of the 3/1 mode DED at the \( q = 3 \) radius is shown \( (I_p = 400kA, B_t = 1.9T) \) with \( I_{DED} = 1.5kA \). In the 3/1 mode the DED the significant mode numbers are \((m < 6, n = 1)\).
with a poloidal plane [11]. A Poincaré-map for a typical vacuum field configuration in TEXTOR is shown in figure 3.3. The field line trajectories have been traced from just within the $q = 1$ radius to the plasma edge with $q_a = 4.6$. The vacuum islands appear as crescent shaped white structures. The O-points lie in the island centers. Outside the magnetic islands and close to the plasma edge the vacuum field is chaotic and the traced field fills the volume in between the islands rather than distinct closed magnetic surfaces. For islands that are locked to DED in its DC-mode of operation the island position is static with respect to a static measurement position. By adjusting the coil currents in the different coil groups the position of the phase of the vacuum island and therewith the locked island can be positioned with respect to the static measurement position.

Figure 3.3: Poincaré-map of the vacuum magnetic field in TEXTOR. The field line tracing calculations are performed between the $q \approx 1$ radius and the plasma edge. The field tracing fills the volume in between the magnetic islands. The vacuum islands with mode numbers $(m = \{1, 2, 3, 4\}, n = 1)$ appear as white crescents.

3.3 Heating systems

To modify the transport in the TEXTOR plasma, additional heating tools from different sources are available. The main heating source is the neutral beam
injection system (NBI), which injects highly energetic neutral particles into the plasma. In the thermalisation process of these particles both electrons and ions are heated. The other heating methods couple electromagnetic waves to the plasma particles. With electron cyclotron resonance heating (ECRH) microwaves are coupled to the electrons only. This ECRH heating method is used here to study electron transport. The ion cyclotron resonance heating (ICRH) couples radio frequency waves to a minority population of ions in the plasma, which then transfer their energy to electrons and ions. Besides heating the plasma these systems can also drive current, and in particular the NBI system is also a particle source as well as a source of toroidal momentum. The combination of these heating methods is a particularly apt tool for transport experiments. The specific properties of the systems are described below.

3.3.1 Neutral beam injection

On TEXTOR two neutral beam injection (NBI) sources are operational [15]. The NBI system ionizes and accelerates ionized atoms to an energy of < 60keV after which they are neutralized so they can be injected into the magnetized plasma. As beam source particles can be taken \{H, D, He\}. Most injected particles become ionized through charge exchange reactions in the plasma and transfer their energy to the thermal electrons and ions through collisions. Each of the two NBI systems can deliver a power of $P_{\text{NBI}} < 1.6\text{MW}$. The two beams inject tangentially to the major radius $R_0$ but in opposite direction to each other. This allows a decoupling of the total injected momentum from the heating power, which is particularly important as the plasma flow has a profound influence on the transport and the effect of the DED on the plasma. For local transport analysis the NBI system has the disadvantage of a relatively large heat deposition profile.

3.3.2 Electron cyclotron resonance heating

On TEXTOR, a gyrotron provides electron cyclotron resonance heating (ECRH) and electron cyclotron current drive (ECCD) [21]. The ECRH and ECCD are commonly used for heat transport experiments and for experiments on modifying the stability of MHD modes in the plasma. The microwave power is coupled to the plasma electrons at the cyclotron frequency or higher harmonics. The gyrotron at TEXTOR couples power $(P_{\text{ECRH}} < 850kW)$ at a frequency of $f_{\text{gyr.}} = 139.85\text{GHz}$ with a stability narrower that $f < 0.1\text{GHz}$ during a $t < 10s$ discharge. The
power coupled to the plasma is X-mode polarized and coupled to the electrons at the 2nd harmonic electron cyclotron frequency ($2f_{ce}$). The radial position of this resonance layer depends on the magnetic field and is approximately given by $R_{140}[m] = 0.70B_t[T]$ (see next section). The vertical and toroidal positioning of the microwave beam is achieved with a steerable mirror. A schematic representation of the deposition location in the poloidal cross-section of TEXTOR is shown in figure 3.4. The steerable mirror allows the beam to make an angle with the equatorial plane of $|\theta_{launcher}| < 30^\circ$ in the poloidal direction and direction and an angle $|\phi_{launcher}| < 45^\circ$ in the toroidal direction. Pure ECRH is achieved for $\phi_{launcher} = 0^\circ$, while $\phi_{launcher} > 0^\circ$ also yields ECCD. The width of the deposition profile typically is $\Delta r/a = 0.05$, i.e. $\Delta r = 2.5cm$ [21]. For perturbative electron heat transport studies the heating power is modulated. The propagation of the induced temperature perturbation is used to analyze the transport properties of the plasma. Modulated power output of the gyrotron is achieved by either switching on and off the beam voltage of the gyrotron, or by modulating the beam voltage [16]. The first mode of operation is referred to as gate modulation and the second mode of operation as beam voltage modulation. With gate modulation a 100% modulation depth is achieved. The characteristic times of the gyrotron installation limit the gate modulation frequency to $f_{mod} < 100Hz$. This method of modulation has the drawback that during the first few $100\mu s$ of each pulse spurious frequency modes (within a few GHz of the centre frequency) appear with a power of a few hundred kW, for which no filters are present in -wave detectors. With beam voltage modulation the output power can be reduced by $80 - 90\%$ while no spurious modes are generated. Also much higher modulation frequencies can be achieved. To minimize the perturbation on the microwave detectors used for the measurement of the heat pulse propagation, the beam voltage modulation technique has been used for the experiments described in this thesis.

3.3.3 Ion cyclotron resonance heating

With ion cyclotron resonance heating (ICRH) electromagnetic power is resonantly coupled to ion species in the plasma [13]. The frequency of the waves is at the ion cyclotron frequency and harmonics. Due to the larger mass of the ions, the ion cyclotron frequency is three orders of magnitude lower than the electron cyclotron frequency. The ICRH system on TEXTOR consists of two independent antennas that are capable of coupling 2MW of power each yielding a total power of $P_{ICRH} < 4MW$ at a frequency of $f_{ICRH} = 25 - 38MHz$. The power can be
Figure 3.4: Schematic representation of the ECRH power deposition location with respect to rational flux-surfaces in a poloidal cross-section of TEXTOR. The power is absorbed at a radially localized resonance layer, schematically represented by the dashed vertical line. The radial position of the resonance depends on the magnetic field. The vertical and toroidal deposition location is varied with a movable mirror that is positioned at \((R = 2.41\, \text{m}, Z = 0\, \text{m})\).

injected continuously up to 3s. Usually a minority heating is used such as the heating of \(H\)-ions (10\%) in a \(D\)-plasma, as is the case in chapter 5. The heating creates an energetic ion minority that through collisions causes the heating of the electrons.

3.4 Diagnostic systems

For the study of local electron heat transport, high-resolution measurements of the electron temperature are the primary requirement. At TEXTOR several Electron Cyclotron Emission (ECE) spectroscopy systems – both 1- and 2-dimensional – as well as an advanced Thomson scattering system are installed. The ECE systems offer high time resolution as well as spatial resolution, the Thomson scattering...
system measures full profiles of both electron temperature and density and is absolutely calibrated. Moreover, the overall plasma must be characterized, for which a variety of other diagnostics are available. A brief review of the most important diagnostics is given below. The ECE systems that are central for the studies in this thesis, are treated in detail in chapter 4. A general overview of the core diagnostics available on TEXTOR can be found in [6].

3.4.1 Electron Cyclotron Emission

ECE spectroscopy is based on the fact that a tokamak plasma - under normal conditions - is optically thick at the electron gyro-frequency or its second harmonic. Thus, a measurement of the intensity at this frequency is a measure for the electron temperature - following the Raleigh-Jeans (long-wavelength) approximation of Planck’s black body radiation law. Moreover, a gyro-frequency is localized in space due to its dependence on the magnetic field. Hence, ECE can be used for a spatially and temporally resolved measurement of the electron temperature. With a single viewing line, a temperature profile along a horizontal chord through the plasma is measured (1-d ECE). With suitable optics and a detection array, a 1-d measurement can be obtained.

1-dimensional ECE

On TEXTOR several 1D ECE systems are installed that are used for measurements in this thesis. A brief set description of the system parameters is given here and a more in depth description of these systems is given in [18]. The standard system employed for $T_e$-profile measurements is the 11-channel ECE-system [20]. With this system the ECE is detected from the plasma with a horn antenna at the low field side of the plasma. The antenna pattern in the plasma is $10cm$ vertically which increases beyond the plasma center. The frequencies of the channels are in the range $f_{ECE} = \{105 - 145\}GHz$. Every channel has a bandwidth of $B_{IF} = 200MHz$. The signal is routinely sampled at $2B_{IF} = 10kHz$ and can be sampled at $2B_v = 20kHz$.

The 2D ECE-Imaging system

A recently developed tool for 2D $T_e$-profile measurements is electron cyclotron emission imaging (ECEI). The first versions of ECEI have been developed on the TEXT-U tokamak [10], the RTP tokamak, [4] and the TEXTOR tokamak [3].
A schematic diagram of an ECE-I system is shown in figure 3.5. With large 2D array of sampling volumes at focal plane Detector array Local Oscillator

![Diagram of ECE-I system](image)

Figure 3.5: Schematic representation of an ECE-I. With lenses the plasma volumes are imaged on an array of detectors. The lenses determine the vertical resolution. In the detector array the ECE radiation is mixed with a tunable local oscillator source to generate a lower frequency for further processing. Next the radiation is split in eighth adjacent frequency bands which whose intensity is detected. This yields spatially resolved 2D ECE measurement.

diameter optics ECE radiation from different vertical positions in the plasma is imaged onto an antenna array. Each array element is sensitive to broad bandwidth ECE radiation and is connected to a heterodyne system that resolves the radiation in 8 frequency bands. A 16 element array and an 8 band receiver attached to each element yield 16 × 8 pixel images of $T_e$-profiles and fluctuations of the TEXTOR plasmas. The bandwidth of received frequencies is limited with respect to the ECE bandwidth of the plasma, so that only a limited area of the plasma is imaged. The received frequency band can be varied by the tuning of the local oscillator, which corresponds to a change in radial observation position. The time-resolution achieved with the system is typically 200kHz. The configuration of the ECE-I system on TEXTOR has steadily been upgraded. For this thesis different configurations are used in chapter 6 [5] and chapter 7 [22, 17]. The basic features of these configurations are discussed in chapter 4.

### 3.4.2 Thomson Scattering

A high-resolution incoherent Thomson Scattering diagnostic for $T_e$ and $n_e$ measurements [19, 12] is available on TEXTOR. With incoherent Thomson Scattering monochromatic laser-light is scattered on the electrons in the plasma. The scattered light is Doppler shifted due to the velocity of the thermal electrons. The
scattered light is spectrally resolved and measured. Since the spectral broadening results from the Doppler shift caused by the velocity distribution of the electrons, the width of the spectrum reflects the electron temperature. The number of scattered photons is proportional to the number of electrons, so that the spectral amplitude provides a measure of the electron density. Profiles of $T_e$ and $n_e$ can be acquired by measuring the spectra of scattered light at multiple positions in the plasma. At TEXTOR, Thomson Scattering takes place along a vertical chord at $R_{th} = 1.840m$ (which is $(R_{th} - R_0) = +9.0cm$ from the geometric axis). The spectra are resolved at a spatial interval of $Z = 7.5mm$ with a total of 120 measurements positions in the plasma. A schematic set-up of the system on TEXTOR is shown in figure 3.6. A special feature of this system is that it can also measure the time-evolution of the $T_e$ and $n_e$-profiles (40 consecutive measurements, currently at 5kHz). Technically, this is achieved by a multi-pulse laser. The accuracy of the measurement is limited by the number of detected scattered photons. To enhance the photon statistics, the laser path is made to pass the observed scattering volume multiple times. A double-pass implementation of the system yields an accuracy of $\Delta T_e/T_e \geq 6\%$ and $\Delta n_e/n_e \geq 3\%$. Up to 14 passes have been implemented [Kantor 2009]. Thomson scattering provides an absolutely calibrated measurement of the electron temperature. An advantage of the Thomson scattering measurement over ECE is its higher spatial resolution and that it can measure closer to the plasma edge. A disadvantage of Thomson scattering is that it cannot measure continuously during the discharge and has a relatively poor time resolution. The $T_e$-profiles measured by Thomson Scattering are in particular important in the study of profile shapes in chapter 5.

3.4.3 Interferometer

For standard density profile measurements on TEXTOR a 9-channel HCN-laser interferometer is operated on TEXTOR [14]. The plasma has a density dependent refractive index, which causes that the laser-light that propagates through the plasma experiences a phase-shift that is proportional to the line-integrated density. The interferometer determines the line-integrated electron density along vertical chords at 9 different radii in the plasma. Radial density profile are obtained by inversion of the line-integrated density. Other applications of the interferometer are feedback control of the plasma density and the radial position.
3.4 Diagnostic systems

Figure 3.6: Schematic layout of the double-pass TS diagnostic with an intra-cavity laser probing system at TEXTOR, from [12]. Thomson scattered light is spectrally resolved in a spectrometer whose elements (1-12) are: (1) fiber array output, (2) relay lens, (3) field lens doublet, (4) entrance slit, (5) littrow triplet, (6) grating, (7) two-part mirror, (8) camera objective, (9) image intensifier, (10) coupling lens, (11) beam splitter, (12) CMOS cameras. The elements (13-15) are parts of the laser system near the tokamak: (13) ruby laser table, (14) focusing lens, (15) plasma, (16) collective objective, (17) fiber array input, (18) spherical mirror.

3.4.4 Charge Exchange Recombination Spectroscopy

Besides the electron parameters also parameters of the ion species are diagnosed. A standard technique used is charge exchange recombination spectroscopy (CXRS). Neutral particles injected with the neutral beam undergo charge exchange reactions with ions in the plasma whereby an electron is transferred from the injected neutral particle to the charged plasma ion. The particles formed in this reaction are in an excited state and emit radiation when they decay to a lower state. With CXRS the emitted spectrum is resolved. The ion temperature can be determined from the Doppler broadening of the spectrum and the rotation speed can be determined from the Doppler shift of the complete spectrum. There are two CXRS systems at TEXTOR, one that use the main neutral beam (NBI1) [1] as well a system that uses a dedicated diagnostic neutral beam [2]. CXRS is routinely used for ion rotation and temperature measurements at the Carbon-VI line at 529 nm. The time resolution of the systems is 50 ms. The radial coverage
reaches from the plasma center out to \( r/a \approx 3/4 \).

References


3.4 References


Chapter 3. The TEXTOR tokamak, dynamic ergodic divertor, heating system and diagnostics


4

The Electron Cyclotron Emission diagnostic

4.1 Introduction

In Chapter 3 a brief overview of the diagnostic tools used in this thesis was given. A detailed description of the diagnostic tool central to this thesis, the ECE-Imaging system, is given. The ECE-I system allows detailed 2D measurements of the $T_e$-profile. It is used to resolve the electron heat transport at the 2D structure of a magnetic island (Chapter 6) and the electron heat-flow during the sawtooth crash (Chapter 7). For more general time resolved $T_e$-profile shape measurements 1D ECE systems are used. ECE is a powerful tool mainly because:

- The ECE intensity is directly proportional to the electron temperature under most plasma conditions, i.e. it is black body radiation.

- The frequency of the emission is directly related to the position in the plasma where the emission originates from.

- The propagation of ECE waves out of the plasma is possible for most conditions used here.

- High frequency heterodyne radiometer detection systems of microwaves are well developed and measurements are only limited by the thermal noise of the plasma.

- Advanced data analysis techniques can be employed to further reduce the thermal noise from the plasma and enhance the relevant physical informa-
• Recent development of linear microwaves arrays allows 2D measurements of the electron temperature. Here the TEXTOR system is state-of-the-art and at the forefront of the development of such 2D ECE-I systems.

These features will be detailed in the following subsections. This chapter serves as a basis for the work in later chapters. The well-established theoretical basis of the generation and propagation of ECE in the plasma is summarized first. This is followed by a treatment of the detection systems, electronics and data analysis.

4.2 Electron Cyclotron Emission

The electrons in a tokamak plasma gyrate perpendicular to the magnetic field at the cyclotron frequency \( \omega_{ce} = eB/\gamma m_e \) where \( e \) is the electron charge, \( B \) the magnetic field and \( m_e \) the relativistic electron mass. The electrons emit and absorb radiation at this frequency and its higher harmonics \((n\omega_{ce})\). In the toroidal geometry of a tokamak the magnetic field strength decreases with the major radius \((R)\): \( B \approx B_0 R_0/R \) which causes the cyclotron frequency to be radially dependent. Under certain conditions, discussed below, the ECE radiation is proportional to the electron temperature and can propagate to the plasma edge to be detected.

The intensity of the emitted radiation is described by the emissivity \( j(\omega) \), which is the rate of emission of radiant energy from the plasma per unit volume per unit frequency per unit solid angle \([2]\). Due to the velocity of the plasma electrons, the emission line for a particular \( B(R) \) is broadened. Two basic line-broadening mechanisms are identified. For a plasma with a thermal electron velocity \( v_t = (k_B T_e)^{1/2} \) the relativistic broadening leads to a line-width \( \Delta \omega_n \approx n\omega_{ce}(v_t/c)^2 \). A second broadening mechanism is due to the Doppler effect. Cyclotron radiation of electrons that are observed under an angle \( \theta \) with the \( B \)-field experience Doppler broadening of \( \Delta \omega_n \approx n\omega_{ce}(v_t/c)|N\cos(\theta)| \). Depending on the angle \( \theta \) one of the two broadening mechanisms dominates. The Doppler broadening dominates the relativistic broadening for \((v_{\parallel}/c)\cos(\theta) > (v/c)^2\) and non-relativistic plasmas \((v \ll c)\), except for perpendicular observation \((\theta \approx \pi/2)\). In practice a pure perpendicular observation yielding zero Doppler broadening is not possible, because radiation that originates from a finite volume is always observed under an angle.

A medium that emits radiation at a certain frequency also absorbs radiation at that frequency. The total measured radiation at an observation point depends on
absorption and emission processes. The intensity of radiation $l(\omega)$ per unit area per unit solid angle per unit frequency is given by [1]

$$\frac{dI(\omega)}{ds} = j(\omega) - \alpha(\omega)l(\omega)$$

(4.1)

where $l(\omega)$ is the intensity of the ray, $s$ the distance along the propagation path and $\alpha(\omega)$ the local absorption coefficient. For a wave propagating between $s_1$ and $s_2$ this equation can be solved yielding

$$I(s_2) = I(s_1)e^{\tau_2 - \tau_1} + \int_{s_1}^{s_2} j(\omega)e^{\tau - (\tau_2 - \tau_1)}ds$$

(4.2)

where $I(s_2)$ is the intensity of an incident ray and the "optical depth" is defined as

$$\tau = \int_s \alpha(\omega)dl$$

(4.3)

For a uniform $j(\omega)/\alpha(\omega)$ the solution takes the simple form

$$I(s_2) = I(s_1)e^{-(\tau_2 - \tau_1)} + \frac{j(\omega)}{\alpha(\omega)}\left[1 - e^{-(\tau_2 - \tau_1)}\right]$$

(4.4)

A plasma is called optically thick when $(\tau_2 - \tau_1) \gg 1$. In optically thick plasmas the incident radiation $I(s_2)$ is completely absorbed and the emitted radiation thus depends on the ratio of the emission and absorption coefficients. A fundamental property of any body in thermodynamic equilibrium that is perfectly absorbing (blackbody) is that it emits radiation with a well-defined blackbody intensity. For low frequency radiation in the Rayleigh-Jeans approximation $\hbar\omega \gg k_B T$ the intensity is

$$I_{BB}(\omega) = \frac{j(\omega)}{\alpha(\omega)} = \frac{\omega^2 T_e}{8\pi^3 c^2}$$

(4.5)

When the plasma is sufficiently optically thick, the emission layer is radially localized and the ECE intensity is proportional to $T_e$. The optical thickness across the resonant layer depends on several plasma parameters and is not trivial. In [2] analytical approximations of the optical thickness across the resonance layer have been calculated for the different harmonics and polarizations. Two different polarizations of the electron cyclotron waves are distinguished. The O-mode waves have the electric field parallel to the magnetic field and the X-mode waves...
that have the electric field perpendicular to the magnetic field. For O-mode waves of harmonics \( n \geq 1 \) and propagation angle \( \theta = \pi/2 \) the optical thickness is

\[
\tau_n^O = \frac{\pi n^{2(n-1)!}}{2^{n-1}(n-1)!} \left( 1 + \left( \frac{\omega_p}{n\omega_{ce}} \right)^2 \right)^{n-1/2} \langle D_n \rangle \left( \frac{\omega_{pe}}{\omega_{ce}} \right)^2 \left( \frac{v_t}{c} \right)^2 n R \omega_{ce} 2\pi c \quad (4.6)
\]

with \( \langle D_n \rangle \approx 1 \) and \( v_t = (k_B T_e / m_e)^{1/2} \). For the first harmonic X-mode the optical thickness is

\[
\tau_1^X = 5\sqrt{2}\pi^2 \left( 1 - \frac{\omega_{pe}^2}{2\omega_{ce}^2} \right)^{3/2} \langle B(z_1) \rangle \left( \frac{\omega_{pe}}{\omega_{ce}} \right)^2 \left( \frac{v_t}{c} \right)^4 R \omega_{ce} 2\pi c \quad (4.7)
\]

with \( \langle B(z_1) \rangle \) defined in [2]. For the X-mode harmonics \( n > 1 \) the optical thickness is

\[
\tau_n^X = \frac{\pi^2 n^{1(n-1)!}}{2^{n-1}(n-1)!} \langle A_n \rangle \left( \frac{\omega_{pe}}{\omega_{ce}} \right)^2 \left( \frac{v_t}{c} \right)^{2(n-1)} R \omega_{ce} 2\pi c \quad (4.8)
\]

with \( \langle A_n \rangle \approx 1 \) and \( \langle D_n \rangle \approx 1 \) for \( \omega_{pe} \ll n\omega_{ce} \). For the parameters of a typical tokamak such as TEXTOR the first and third harmonic X-mode and second harmonic O-mode are optically thin \( (\tau_1^X \approx \tau_3^X \approx \tau_2^O < 1) \). The first harmonic O-mode and the second harmonic X-mode are optically thick \( (\tau_1^O \approx \tau_2^X \geq 1) \) for the central part of the plasma. Due to the lower \( T_e \) and \( n_e \) near the plasma edge the optical thickness there becomes low which limits the applicability for \( T_e \)-measurements close to the edge. Profiles of the optical thickness of the different harmonic O and X-mode for typical plasma conditions are shown in figure 4.1b.

### 4.3 Electron Cyclotron Wave propagation

The plasma is a refractive medium for ECE radiation, which has consequences for the propagation of ECE and the possibility to detect it. The dispersion relation for waves propagating in a medium with a certain refractive index \( N \) is given by

\[
N = \frac{k c}{\omega} \quad (4.9)
\]

The refractive index \( N \) depends on the polarization of the wave with respect to the magnetic field. The O-mode polarization and X-mode polarization have the wave electric field parallel to the magnetic field and perpendicular to the magnetic field.
respectively. The refractive index in the cold plasma approximation \((v_t/c \approx 0)\) is given by the Appleton-Hartree formula [Hutchinson 1987]

\[
N_{O,X}^2 = 1 - \frac{X(1 - X)}{1 - X - \frac{1}{2} Y^2 \sin^2(\theta) \pm \left[ \left( \frac{1}{2} Y^2 \sin^2(\theta) \right)^2 + (1 - X)^2 Y^2 \cos^2(\theta) \right]^{1/2}}
\] (4.10)

with \(X = (\omega_{pe}/\omega)^2\), \(Y = \omega_{ce}/\omega\). The signs \(\{+,-\}\) correspond to the X-mode and O-mode propagation respectively. The plasma frequency only depends on the plasma density \(n_e\)

\[
\omega_{pe} = \sqrt{\frac{n_e e^2}{\epsilon_0 m_e}}
\] (4.11)

with \(\epsilon_0\) the permittivity of vacuum. A propagating wave can encounter cut-offs where the wave cannot further propagate \((k \to 0, N \to 0)\) and resonances where the wave is absorbed \((k \to \infty, N \to \infty)\). To minimize Doppler broadening most ECE measurement applications rely on perpendicular observation \((\theta = \pi/2)\). The cut-offs and the resonances for this angle are treated here. For O-mode polarization the refractive index reduces to \(N_{O}^2 = 1 - (\omega_{pe}/\omega)^2\) and wave-propagation then experiences a cut-off for \(\omega < \omega_{pe}\) which depends only on the density. For perpendicular observation X-mode polarization the refractive index reduces to \(N_{X}^2 = (\omega^2 - \omega_{UH}^2)/(\omega^2 - \omega_{ce}^2)\). The X-mode has two cut-offs and two resonances. The resonances occurs at \(\omega = 0\) and the upper-hybrid frequency \(\omega_{UH} = (\omega_{pe}^2 + \omega_{ce}^2)^{1/2}\). A higher and a lower cut-off frequency occur at \(\omega_{\pm} = \pm \omega_{ce}/2 + (\omega_{pe}^2 + (\omega_{ce}/2)^2)^{1/2}\). Above these cut-offs and resonances the waves can propagate in the plasma. The possibility of the ECE waves to propagate through the plasma and to be used for \(T_e\) measurement can conveniently be determined by plotting the resonances and cut-offs together with the ECE frequency. In figure 4.1a the resonance and cut-offs frequencies across the minor radius are plotted for a parabolic \(T_e(R)\) and \(n_e(R)\)-profile, together with the optical thicknesses in figure 4.1b.

The first harmonic O-mode is optically thick but its propagation out of the plasma will experience a cut-off for higher densities that are commonly used. The first harmonic X-mode is optically thin and not measurable. The second harmonic X-mode is optically thick and it does not experience cut-offs for commonly used \(T_e\) and \(n_e\)-profiles. The second harmonic X-mode is therefore mostly used to measure \(T_e\) on TEXTOR. Near the plasma edge the optical thickness however does
become low due to the low $T_e$ and $n_e$ there. The $T_e$ measurements from higher harmonics ($n > 2$) ECE is complicated because of the low optical thickness of the plasma for these harmonics. For low optical thickness also ECE contributions from lower harmonics from different locations may shine through the resonant layer and contribute to the measured intensity.

![Figure 4.1](image.png)

**Figure 4.1:** (a) Resonances and cut-offs together with the ECE frequency for $B_t = 2.25 T$ and $n_e(0) = 3 \times 10^{19} m^{-3}$. The O-mode polarization experiences a cut-off for $f < f_{pe}$. The X-mode has a cut-off for $f < f_+$ or $f < f_-$ and it has a resonance at $f_{\nu H}$. (b) The optical thickness for perpendicular propagation of the first three harmonics of $f_{ce}$ for X-mode propagation and O-mode propagation perpendicular to the $B$-field. The first harmonic O-mode ($\tau_1^O$) and the second harmonic X-mode ($\tau_2^X$) are optically thick ($\tau \gg 1$) in the plasma center and the ECE from these harmonics can be used for $T_e$-measurements. The higher harmonics ($n > 2$) are complicated to interpret because they are generally optically thin ($\tau < 1$) and may contain ECE contributions from different locations and or different harmonics.

### 4.4 Heterodyne ECE detection

The ECE radiation that propagates out of the plasma is detected for time resolved $T_e$-profile measurements. The radiometers on TEXTOR spectrally resolve the ECE to achieve radial and time resolved $T_e(R, t)$-measurements. The basic ECE detection scheme and the consequences it has for the radiation noise level are described here.
At the typical fields $B_t \approx 2T$ in TEXTOR the frequency of the second harmonic ECE radiation is in the order of $f_{ECE} = 100\,\text{GHz}$ ($\lambda = 3\,\text{mm}$). To detect radiation in this frequency range, there are several measurement techniques that can be used. The ECE radiometers employed on TEXTOR use the heterodyne detection technique [6]. The basic feature of heterodyne detection is that high frequency waves (RF) are down-converted to a lower intermediate frequency (IF) where further signal processing can be performed much better. A schematic representation of the basic heterodyne detection scheme is shown in figure 4.2. The

![Figure 4.2: Heterodyne detection scheme employed in ECE radiometers on TEXTOR (adapted from [6]). (b.) Schematic representation of the frequency bands that are selected by the mixing. The mixer output contains both an upper and lower sideband component. If only one band is required, a sideband needs to be filtered out of the RF signal before mixing.](image)

ECE microwaves from the plasma with frequency $f_{RF}$ are mixed with microwaves generated by a local oscillator with frequency $f_{LO}$ in a mixer that generates signals with frequency bands $f_{IF} = |f_{RF} \pm f_{LO}|$. Only the difference frequency in the intermediate frequency band is of interest $B_{IF} < |f_{RF} - f_{LO}|$. The output frequency from the mixer $f_{IF}$ consists of frequencies from lower side band (LSB) and upper sideband (USB). These correspond to different positions in the plasma. Two mixing schemes are distinguished: single sideband mixing and double sideband mixing. In single sideband mixing, one sideband is rejected by applying a side band filter to the RF waves before mixing. In double side band mixing, no side band filter is applied before mixing and both LSB and USB are retained. This can be done if $f_{LO}$ is set in the middle of both the USB and the LSB. After detection the bandwidth of the signal is further reduced with a low-pass filter of bandwidth $B_v$ and the signal is stored with ADC’s for analysis.
4.5 The 2D ECE-Imaging system

Below we give a more detailed description of the 2-d ECE-Imaging system. For convenience we repeat the diagram and basic specifications here. ECE radiation from different vertical positions in the plasma is imaged onto an antenna array with large diameter optics. Each array element is sensitive to broad bandwidth ECE radiation and is connected to a heterodyne system that resolves the radiation in 8 frequency bands. A 16-element array and an 8-band receiver attached to each element yield $16 \times 8$ pixel images of $T_e$-profiles and fluctuations of the TEXTOR plasmas. The bandwidth of received frequencies is limited with respect to the ECE bandwidth of the plasma, so that only a limited area of the plasma can be imaged. The received frequency band can be varied by the tuning of the local oscillator which corresponds to a change in the radial observation position. The configuration of the ECE-I system on TEXTOR has steadily been upgraded. For this thesis different configurations are used in chapter 6 [5] and chapter 7 [10, 8]. The basic features of these configurations are discussed.

4.5.1 ECE-I optics

The large diameter optics that focuses the ECE onto the antennas gives the ECE-I system a high spatial resolution in the vertical direction. Different optical config-

![Figure 4.3: Schematic representation of an ECE-I system. With lenses the plasma volumes are imaged on an array of detectors. The lenses determine the vertical resolution. In the detector array the ECE is mixed with a tunable local oscillator source to generate a lower frequency for further processing which is further split in frequency in frequency bands corresponding. This is then further split in frequency bands corresponding to different radial positions and these are detected.](image-url)
urations have been used. The configuration used in chapter 6 is shown in figure 4.4a. This system was designed to coexist with a microwave imaging reflectome-

![Optical layouts of the TEXTOR ECEI system.](image)

**Figure 4.4:** Optical layouts of the TEXTOR ECEI system. (a.) This system [5] was designed to coexist with a microwave imaging reflectometer, with both systems sharing two large area plasma-facing E-plane (6) and H-plane (5) cylindrical mirrors. A beam splitter serves to separate the two systems, positioned between mirror (5) and spherical lens (4). The positions and focusing properties of lenses 2-4 are designed to image the ECE layer (8) onto the mixer array which resides on the back of substrate lens (1). (b.) Vertical zoom optics ECE-I. The two plasma-facing lenses form a zoom lens and allow the vertical extent of the observation area to be varied from 20cm to 30cm. The next lens provided an independent control of the focal plane position within the plasma.

With a focusing lens the resonance layer is focused onto the antenna array. The ECE beam patterns of this system have been characterized in a laboratory [9]. The vertical inter-channel spacing of this system is $\Delta Z = 11mm$ with a total vertical plasma coverage of 17cm. The FWHM spot size is $13mm$ in the vertical direction and $9mm$ in the toroidal direction.

Substantial improvements in the optical capabilities have been achieved in the latest configuration of the ECE-I system on TEXTOR [10]. The application of substrate mini-lenses facilitates better coupling of the radiation to the antennal
array and results in improved beam patterns. An additional capability is that the vertical size of the observation area is adjustable by employing a zoom lens capability. A ray-tracing lay-out of this configuration is schematically shown in figure 4.4b. The two plasma-facing lenses form a zoom lens and allow the vertical extent of the observation area to be varied from 20 cm 'small view' to 30 cm 'wide view'. The next lens provided an independent control of the focal plane position within the plasma. The beam patterns have been characterized in the laboratory for both the small and the wide view [8]. The measured spot sizes are $9 - 15 \text{mm}$ full width at half maximum (FWHM) across all channels in the small and wide zoom configurations, respectively [8].

In experiments such as described in chapter 6, large amounts of gyrotron power ($< 800 \text{kW}$) are injected that can be $< 13$ orders of magnitude larger than the detected ECE ($< 100 \text{nW}$) from the plasma. If the gyrotron power is not completely absorbed, a fraction of this power can be reflected on the tokamak wall and be coupled to the ECEI system. This can affect the signal or even damage the mixer elements. On TEXTOR the gyrotron is close to the ECEI system with a toroidal separation $\Delta \phi = 22.5^\circ$ and can easily affect the ECEI system. To reduce the effect of the gyrotron radiation, a notch filter with notch frequency $140 \text{GHz}$ is inserted in the transmission line. In the later ECEI configuration a cascade of three notch filters with a total notch-to-pass-band rejection ratio of $69 \text{dB}$ is inserted [8]. In the earlier configuration of ECEI (used for chapter 6) only a single notch is inserted which achieved a limited rejection ratio.

4.5.2 ECEI electronics

The ECEI system uses a double down-conversion approach. The radiometer electronics of ECEI described in [5] is schematically shown in figure 4.5.

A band of ECE from the plasma (RF Signal) is received together with a LO signal for each antenna element. The LO-signal is provided by a tunable backward wave oscillator (BWO) which has a frequency range $f_{LO} \approx 80 - 140 \text{GHz}$. By the application of a high-pass filter, only the upper side-band of the ECE is mixed with the LO signal. The mixing process is wide band with a $3 \text{dB}$ loss bandwidth that extends past $30 \text{GHz}$. The RF-signal is amplified by $35 \text{dB}$ and coupled to an RF-electronics module where the band is split in 8 distinct frequency bands which are down converted (double sideband). The 8 frequency bands range from 2.4 – 8.0 GHz while for the upgraded ECE-I configuration they range from 2.4–8.7 GHz. The amplitude of each down converted signal is detected in an IF-electronics
module. The (double) bandwidth of each detected signal is $B_{IF} = 400\text{MHz}$. Before sampling the detected signal, a variable frequency low-pass filter is applied to $B_v < 400kHz$. To study localized 2D $T_e$ phenomena that are evolving relatively fast with respect to the rotation period of the plasma the ECEI system is very well suited. The high spatial resolution is achieved by large diameter optics and the

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**Figure 4.5**: Schematic layout of ECEI detection electronics (adapted from [5]). (a.) Coupling to antenna and mixing with LO signal yields a down-converted band of ECE frequencies. (b.) In the RF Electronics module each band of frequencies is resolved in 8 frequency bands and subsequently down-converted (double sideband). (c.) In the IF-Electronics module the down converted signals are low pass filtered yielding $B_{IF} = 700\text{MHz}$ before the amplitude is detected. A variable bandwidth low-pass filter with $B_v < 400kHz$ is applied to the detected signal before it is sampled.
high time resolution is achieved by fast microwave electronics. For general 1D profile evolution the 1D ECE systems are well suited. A larger 2D coverage is however be needed to get a more complete picture of large-scale 2D $T_e$ dynamics of islands or the sawtooth crash.

4.6 ECE noise-limits to resolution and signal enhancement

ECE intensity measurements contain noise that obfuscates small $T_e$ fluctuations. The noise comes from different sources. The ECE emitted by the plasma is inherently noisy. Additional noise is generated by the radiometer itself. The measured noise amplitude is characterized by its squared average amplitude $\sqrt{\langle i^2(t) \rangle}$ (standard deviation) with respect to the average intensity $I(t)$ (which is proportional to $T_e$ for optically thick plasmas). For the ECE radiation the ratio of both quantities (the signal to noise ratio: $S/N$) are related to the video bandwidth $B_v$ and the detected radiation bandwidth $B_{IF}$

$$\frac{\sqrt{\langle i^2(t) \rangle}}{I(t)} = \sqrt{\frac{\langle \Delta T_e^2(t) \rangle}{T_e(t)}} = \sqrt{\frac{2B_v}{B_{IF}}}$$

(4.12)

These bandwidths are directly related to the spatial resolution and the time resolution. The minimal radial resolution is given by $\Delta R_{IF} \geq \Delta f/|df/dR|$ where $\Delta f = B_{IF}$ and $df/dR = n(df_{ce}/dR)$. For ECE1 the layer $\Delta R_{IF} = 0.5\,cm$ for $B_{IF} = 400\,MHz$ and $R = 1.75\,m$. The maximum time resolution is given by the Nyquist-Shannon sampling theorem which states that a signal needs to be sampled at half the bandwidth of the inverse signal bandwidth $\Delta t \geq 1/(2B_v)$. The minimal resolution of the $T_e$-measurement is thus limited by both the time and radial resolution. Decreasing either the radial or time resolution will increase the signal-to-noise ratio. The radial resolution is however basically limited by the finite width of the emission layer due to Doppler broadening and relativistic broadening. For typical TEXTOR plasma conditions the width of the emission layer has been calculated to be $\Delta R_{EM} = 0.5\,cm$ [3]. The total radial resolution is the standard deviation of the widths $\Delta R = (\Delta R_{IF}^2 + \Delta R_{EM}^2)^{1/2} \approx 0.8\,cm$. While in practice, the bandwidth $B_{IF}$ is fixed by the hardware filters. The post-detector bandwidth $B_v$ is more flexible as it can also be decreased by software filter operations. For certain applications the signal-to-noise ratio ($S/N$) of the ECE needs to be enhanced to detect useful information. Depending on the characteristics of the signal different filtering methods are used. If only a specific bandwidth of the
The \( T_e(t) \) signal is of interest, the rest of the bandwidth can be filtered out by whereby \( B_v \) in eq. 4.12 is effectively reduced. For example, if only the slowly varying part of the ECE signal is of interest, the high frequency components can be filtered out by applying a low-pass band filter.

### 4.6.1 Fourier transform filtering

In this thesis periodic or nearly periodic \( T_e \) fluctuations are investigated. In chapter 6 the amplitude and phase are determined of a periodic \( T_e \) perturbation induced by modulated ECRH. The standard way to determine the amplitude and phase of a signal is by Fourier transform. For a discretely sampled signal usually a fast Fourier transform (FFT) algorithm is employed. In the presence of noise this allows the amplitude and the phase of a signal to be recovered. This is illustrated by determining the amplitude and phase of a simulated noisy ECE signal. The simulated signal is signal is of the form

\[
T_e(t) = T_{e,0} + T_{e,1} \cos(2\pi f_1 t) + \Delta T_e,
\]

with average values \( T_{e,0} = 100\,\text{eV} \) and an oscillation of frequency \( f_1 = 50\,\text{Hz} \) and amplitude of \( T_{e,1} = 2.6\,\text{eV} \). The noise has a standard deviation equal to the amplitude of the oscillation. The signal is sampled at a frequency \( f_s = 2\,\text{kHz} \) (= \( 2B_v \)). The noisy signal together is shown for an interval of \( 1\,\text{ms} \) (200 samples) and the amplitude of a Fourier transform spectrum are shown in figure 4.6. The signal without noise is depicted by the thick blue line. The thin blue line is the same signal with white noise added. From the Fourier transform the amplitude and the phase of the reconstructed signal are determined (thick red line). The uncertainty in the amplitude and the phase of the signal depends on the noise on the amplitude and phase which can be determined from the spectrum. The standard deviation of the noise level is depicted by the dashed green line. This is also the uncertainty in the amplitude \( \Delta T_{e,j}/T_{e,j} \) and phase of the different harmonics. The amplitude of the noise level in the spectrum decreases relative to the amplitude of the Fourier components with an increasing time interval (proportional to the inverse square root of the time interval). The amplitude and phase of an arbitrarily small harmonic temperature oscillation can be resolved with a Fourier transform over a long enough time interval if it is larger than the noise level. The accuracy increases with the length of the integration time of the Fourier transform. The Fourier transform can very well be used to reduce the noise level of a periodic signal whose spectrum consists of a few harmonics. The FFT transform is less effective in reducing the noise levels of a signal whose spectrum consists of many harmonics, such as a signal exhibiting a large discontinuity.
Chapter 4. The Electron Cyclotron Emission diagnostic

![Simulated ECE signal](image)

**Figure 4.6:** Simulated ECE signal (top panel) and its amplitude spectrum (lower panel). The signal has an average value $T_{e,0} = 100\, eV$ and an oscillation of frequency $f_1 = 50\, Hz$ and amplitude of $T_{e,1} = 2.6\, eV$. The noise has a standard deviation $T_{e} = 2.6\, eV$ equal to the amplitude of the oscillation. The signal is sampled at a frequency $f_s = 2\, kHz$ ($= 2\, Bv$). The noisy signal together is shown for an interval of $1\, ms$ (200 samples).

### 4.6.2 Singular value decomposition

For multi-channel ECE data the signals from the plasma can be strongly redundant, meaning that the same signal features are present in the measurements sequences of multiple channels. The ECE-Imaging system measures with 128 channels in a relatively small area of the plasma. In case the scale of the temperature dynamics is large compared to the ECE-Imaging observation area, similar $T_e$ features are present at the same time in the different channels and the set of measurement vectors is redundant. The redundancy in set of data vectors can be very well analyzed by making a singular value decomposition (SVD) [4]. In a SVD of a set of $T_e(t_i)$ vectors measuring at different position $x_j$ in the plasma, the
signal is decomposed as

$$T_e(t_i, x_j) = \sum_{k=1}^{K} s_k u_k(x_j) v_k(t_i)$$  \hspace{1cm} (4.13)

where the base vectors \( \{u_k\} \) and \( \{v_k\} \) form orthonormal bases in space and in time, respectively. The basis vectors \( \{u_k\} \) and \( \{v_k\} \) are linear combinations of the \( T_e(t_i, x_j) \)-vectors. Each pair of space and time vector have a weight \( s_k \). In the SVD the weight vectors are ordered in descending order. An approximation of the ECE signals is made by retaining only \( L < K \) components in the reconstruction

$$T_e(t_i, x_j) = \sum_{k=1}^{L<K} s_k u_k(x_j) v_k(t_i)$$  \hspace{1cm} (4.14)

This is the best approximation of the set of data vectors in the least square sense on an \( L \)-dimensional basis. The approximation becomes better for an increasing number of components (\( L \)). In the case of a strongly redundant set of \( T_e \)-vectors, the data set is well approximated with only a few components. The thermal ECE noise on the different measurement vectors is uncorrelated and thus is not redundant. The noise energy is distributed over all the pairs of basis vectors \( \{s_k\}, \{u_k\} \) and \( \{v_k\} \). By retaining only a limited number of eigenvectors, the noise energy is reduced while the energy of the signal is retained. As an example the reduction by means a SVD of the noise level of a 1D ECE signal is simulated which resembles a sawtooth precursor oscillation as described in chapter 7. The results are shown in figure 4.7. The signal contains \( j \leq 100 \) spatial channels and \( i \leq 200 \) time points. The simulated original signal figure 4.7a resembles a temperature perturbation due to a sawtooth precursor oscillation that is followed by a sawtooth crash of the central \( T_e \) and simultaneously a sudden increase of \( T_e \) further outward. White noise is added to the original signal figure 4.7c and a reconstruction of a limited number of eigenvectors is shown in figure 4.7e. The singular values of the SVD are shown in figure 4.7b. The blue line denotes the singular values of the original signal. The signal is well described by the first 6 components of the SVD. The green line represents the singular values when noise is added. The noise causes the singular values to be shifted by the noise level [4]. The noise level is represented by the near horizontal line. Only the first five components (denoted by the colored circles) are larger than the noise level. The corresponding eigenvectors \( \{u_k\} \) called topos and \( \{v_k\} \) called chronos,
Figure 4.7: Singular value decomposition analysis of redundant multi-channel $T_e$ measurements.
The simulated original signal is shown in panel (a). In panel (c) white noise is added to the original signal. Panel (c) shows the reconstructed signal. In panel (b) the singular values of the original signal (blue line) are shown together with the singular values of the noisy signal (green line). The original signal is strongly redundant; most of the signal is contained in the first six components of the SVD. The first five components are above the noise level which is observed as a near horizontal green line. The corresponding eigenvectors $\{u_k\}$ and $\{v_k\}$ are shown in panel (d) and panel (f) respectively. The first five components are retained for the reconstruction shown in panel (e). In the reconstruction the noise level is clearly reduced compared to panel (c) and the contours similar to the original structure shown in panel (a) can be identified.

are shown in figure 4.7d and figure 4.7f, respectively. For larger $k$ the relative
noise level on the eigenvectors increases. For \( k > 5 \) the eigenvectors appear completely dominated by noise. The first five components are retained for the reconstruction which is shown in figure 4.7e. In the reconstruction the noise level is clearly reduced compared to figure 4.7c and the contours similar to the original structure shown in figure 4.7a can be identified. The enhancement is due to the fact that most of the redundant signal is retained while most of the noise which is not redundant has been rejected. While an FFT is of limited applicability in reducing the noise level of signals with a large discontinuity, the above example shows that the SVD can effectively reduce the noise level of redundant signal with a large discontinuity.

### References


5

Pressure profile consistency in TEXTOR

5.1 Introduction

Understanding and controlling the transport of heat and particles in a tokamak device is important for predicting and controlling the fusion performance of future fusion devices. The fusion performance in a tokamak depends on the efficiency by which the plasma pressure can be sustained in the reactor, while the efficiency depends on electron and ion transport in fusion plasmas. The transport is dominantly due to turbulent processes. The turbulent electron thermal transport in particular exceeds the collisional transport by one to two orders of magnitude, whereas the turbulent ion thermal transport is a few times larger than the collisional transport. Additional heating increases the turbulent transport, so that the thermal energy content of the confined plasma increases less than proportionally to the additional heating power.

There is no single, first-principles theoretical model that predicts the turbulent transport. Various types of turbulence, in which many plasma parameters are involved, may contribute to the total transport. The best numerical turbulence models are achieving increasingly better descriptions of experimental observations, but the run time of these codes is very long and detailed predictions are not yet in reach. For this reason, scaling laws are commonly used to predict the performance of future tokamaks, such as ITER. But these predict rather global aspects of the confinement and are not concerned with details of the profiles of temperature, pressure and current density. Yet, these profiles are very important
for the control of a burning plasma.

There is another approach to the description of transport that does consider the profiles. This approach came forth from the observation that the temperature profiles have similar shapes for a wide variety of plasma conditions. This was first observed for the electron temperature ($T_e$) profiles of Ohmic heated plasmas [4], where the relation between the electron temperature and current density profile is imposed by the (electron temperature dependent) resistivity, so that in turn the Ohmic dissipation profile is determined. The term 'profile consistency' was introduced to describe the interrelationship of these profiles. Several different principles were later invoked to explain the observation of these 'preferred' profile shapes [3, 11, 22]. These theories have the character of minimum energy states and the idea is that the turbulent transport adjusts itself in such a way that the plasma finds these states. This implies a non-local aspect of the transport in a tokamak.

Later it was also observed that in additionally heated plasmas the electron and ion temperature ($T_i$) profiles react weakly to changes of the heating power deposition [13, 16, 21, 25]. This phenomenon is known as temperature 'profile resilience' or 'profile stiffness'. For the plasma electrons a strong increase in heat transport is observed when the temperature gradient exceeds a critical value. Critical gradient models explain why the temperature profile is stiff, while being fully locally determined [16, 17].

Additional to observations of stiff temperature ($T$) profiles it is also observed that the electron pressure ($p_e$) profiles have very similar shapes for both Ohmic and additionally heated plasmas. Both quantities are related since the electron pressure is the product of the electron density ($n_e$) and the electron temperature $p_e = n_e T_e$. The $n_e$-profile appears to adapt in such a way that the $p_e$-profile is conserved. In particular, it has been observed that when the plasma is heated locally using electron cyclotron resonance heating (ECRH), the density reduces at the heat deposition location, a phenomenon referred to as 'pump-out' [5, 6, 7, 14, 19, 24]. It has been proposed that temperature profile stiffness is a special case of pressure profile conservation [15, 5, 19]. However in ASDEX-Upgrade it is observed that while the electron temperature profile is stiff, the flattening of the density profile does not always occur for central ECRH [1]. Only for a large $T_e/T_i$, a flattening of the central density profile due to ECRH is observed to yield a stiff pressure profile [1]. In off-axis ECRH experiments on TEXTOR, the formation of off-axis maxima in the $T_e$-profile is accompanied by a local reduction of the $n_e$-profile [5]. The time-scales of the $n_e$-profile modification is however larger than
the time-scale of the $T_e$-profile modification, implying that the pressure profile shape was at least not transiently maintained.

For the development of a control strategy for a burning plasma, it is important to know whether profile consistency applies in its strong form, in which case a local description of turbulence induced transport appears to be of limited value, or that the profiles are merely stiff to some extent. In this chapter it is investigated if electron temperature profiles stiffness is a special case of electron pressure profile consistency. Different models exist that describe the stiffness of $T$- and/or $p$-profile. A comparison between the models is made to identify the conditions where the predicted profile shapes coincide or differ. Experiments are performed on the TEXTOR tokamak to create conditions where the $T_e$-profile and $p_e$-profile shapes are expected to coincide or differ according the different theories. In addition to ECRH used in the previous experiments on TEXTOR [5, 15], ion cyclotron resonance heating (ICRH) and neutral beam injection (NBI) is applied, which causes heating of both ions and electron species. The measurements of the electron temperature and density profile shapes are done with the advanced high-resolution Thomson Scattering system (TS) at TEXTOR [23, 12].

5.2 Profile models

In this section the basic models that describe temperature profile stiffness and the conservation of the pressure profile are discussed. At the end of the section the compatibility of the models with respect to the profile shapes they predict is discussed.

5.2.1 Critical Temperature gradient transport

A critical gradient model accounts for the onset of increased transport above a critical gradient. Electrostatic instability modes such as trapped electron modes (TEM), electron temperature gradient modes (ETG) and ion temperature gradient modes (ITG) could be responsible for turbulent transport and these modes do exhibit critical temperature gradient dependence in their onset. In [8] this kind of behavior is modeled in terms of the heat conduction equation $Q = -\chi n \nabla T$ with a heat-diffusivity:

$$\chi = \chi_g B \chi_s (R|\nabla \ln T| - \kappa_c) H(R|\nabla \ln T| - \kappa_c) + \chi_0$$  \hspace{1cm} (5.1)
in which $\chi_{gB} = q^\nu(T/eB)\rho_s/R$ is a heat-diffusivity based on scaling of gyro-Bohm turbulent transport. In these equations $B$ is the magnetic field strength, $q$ is the local safety factor, $\rho_s = \sqrt{m_iT/eB}$ is the Larmor radius and $H(...)$ is the Heaviside function. The dimensionless ‘stiffness parameter’ $\chi_s$ stands for the enhancement of the turbulent heat-diffusivity when the normalized temperature gradient, $R|\nabla \ln T|$, becomes larger than a critical value $\kappa_c$. Similarly the factor $\chi_0$ indicates the lower level of turbulent transport if the temperature gradient is smaller than the critical one. It should be noted that the factor $q^\nu$ is introduced ad-hoc, as means to obtain the experimentally observed increase of the total kinetic energy content of the plasma with the plasma current. The value $\nu = 3/2$ gives the best result as a compromise between various experiments. The gyro-Bohm diffusivity increases strongly with temperature ($\chi_{gB} \propto T^{3/2}$) and therefore the stiffness of the profile increases with the temperature.

A typical $T$-profile described by this model is shown in figure 5.1. The $T$-profile can be divided in three regions. A core region where the $T$-profile is non-stiff $R|\nabla \ln T| < \kappa_c$ due to a low $Q$. A stiff region where $R|\nabla \ln T| \approx \kappa_c$. An edge region where $R|\nabla \ln T| \gg \kappa_c$ but the profile is not stiff since $\chi_{gB}$ is low and other loss mechanisms start to play a role. The transition from the non-stiff core region to the stiff region is well localized near the center. The transition from the stiff region to the edge region is smooth. The edge region becomes smaller with increasing $T(a)$.

When a large amount of heat is deposited in the plasma center, this will cause high $T$ and high $Q$ and therefore this model predicts stiff $T$-profiles with $R|\nabla \ln T| \approx \kappa_c$ for a large part of the minor radius. If $\kappa_c$ is independent of the plasma conditions then the $T$-profiles have similar shapes. Modeling of $T_e$-profiles on different machines yields $\kappa_{ce} = 3 – 8$ with a relatively large variability of the stiffness parameter $\chi_s = 0.3 – 4.5$ and the background diffusivity $\chi_0 = 0.02 – 0.3$ [9]. The large variability together with the arbitrarily introduced $q$-dependence are weaknesses of this model. There is no guideline what values should be chosen for a given experiment.

5.2.2 Profile consistency

In Ohmic tokamak plasma discharges it has often been observed that besides a self-consistent or stiff $T_e$-profile, also the $p_e$-profiles is conserved [7]. Furthermore the normalized profile shapes of the electron pressure and the plasma current are
5.2 Profile models

![Graph showing modeled T-profile and its logarithmic gradient](image)

**Figure 5.1:** Modeled $T$-profile (solid blue top) and its logarithmic gradient (solid blue bottom) (adapted from [8]). The model parameters are $\{X_0 = 0.3, X_s = 5, \kappa_c = 5, T_a = 0.35 keV\}$. The dashed red line represents profile with $-R \nabla \ln(T) = \kappa_c$. In the core the heat flux is low the profile is not stiff $-R \nabla \ln(T) < \kappa_c$. In the stiff region the $-R \nabla \ln(T) \approx \kappa_c$. In the edge region $-R \nabla \ln(T) \gg \kappa_c$ due to the relative low $T$ there. The transition between the stiff region and the edge region is smooth.

observed to coincide very well [5, 18, 19]:

\[
\frac{p_e(r)}{p_{e0}} \approx \frac{j(r)}{j_0}
\]  
(5.2)

Kadomtsev and others [3, 10, 11] have calculated the ‘natural’ shape of the $j$-profile under the assumption that plasma relaxes to a state of minimal free energy:

\[
\frac{j(r)}{j_0} = \left[1 + \left(\frac{qa}{q_0} - 1\right) \left(\frac{r}{a}\right)^2\right]^{-2}
\]  
(5.3)

A similar form for the $j$-profile was derived in [22] assuming the most likely statistical current distribution. The $j$-profile induces a poloidal magnetic field ($B_p(r) =$...
\[ \mu_0 \int_0^r j(r') 2\pi r' dr' / 2\pi r \] that gives the safety-factor profile \( q(r) = r B_t / R_0 B_p(r) \).

The natural \( q \)-profile can be written as:

\[ q(r) = q_0 + (q_a - q_0) \left( \frac{r}{a} \right)^2 \] (5.4)

with \( q_0 = 2 B_t / R_0 \mu_0 j_0 \) and \( q_a = 2 \pi a^2 B_t / \mu_0 l_p R_0 \). The assumption that the plasma adapts its \( j \)-profile such that the contribution of the poloidal magnetic field energy to the total free energy is minimal, together with the fact that the \( j \)-profile induces via the Maxwell equations the poloidal field, makes that eq. 5.3 leaves no free parameters other than \( q_0 \) for an experimentally chosen \( q_a \). This does not hold for the contribution of the \( p \)-profile to the free energy as its influence on the total magnetic energy is minute. This makes that in the \( p \)-profile another free parameter can play a role without influencing the state of minimal free energy. This parameter is an additive pressure which is constant over the plasma cross-section. The \( p \)-profile can therefore be written as:

\[ p(r) = p_1 \left[ 1 + \left( \frac{q_a}{q_0} - 1 \right) \left( \frac{r}{a} \right)^2 \right]^{-2} + p_2 \] (5.5)

The edge pressure is often observed to increase with an order of magnitude when an edge barrier is formed (the so-called ‘H-mode’). In an H-mode the shape of the radius dependent part of the pressure is preserved and it is the same as when the edge pressure is low (the so-called ‘L-mode’). This is in agreement with eq. 5.5: in L-mode \( p_2 \) is negligible or even slightly negative, whilst in H-mode \( p_2 \) can be as large as \( p_1 \). In the early days when eq. 5.2 was observed, additional heating above Ohmic heating was non-existent or too low for H-mode and therefore \( p_2 \approx 0 \), bringing the profile-functions of \( p \) and \( j \) together. In this form the shapes of the profiles only depend on the ratio \( q_a/q_0 \). The edge safety factor \( q_a = 2 \pi a^2 B_t / \mu_0 l_p R_0 \) is readily determined as \( a, B_t, \) and \( l_p \) are externally controlled operational quantities. The central safety factor \( q_0 \) is limited by the ‘sawtooth’ instability. Polariometry measurements of Ohmic sawtoothing plasmas in TEXTOR have yield central safety factors of \( 0.7 \leq q_0 \leq 1 \) [20]. The evolution and the average value of \( q_0 \) during the sawtooth period still a matter of discussion.

Although the value of the central safety factor is still somewhat hard to measure, the position of the sawtooth inversion radius \( r_{inv} \) is close to the \( q = 1 \) radius \( (r_1) \) which is readily determined from profile measurements. According to eq. 5.4
one can derive a good estimate of \(q_0\) from the measured \(r_{\text{inv}}\):

\[
q_0 = \left[ 1 - q_a \left( \frac{r_{\text{inv}}}{a} \right)^2 \right] / \left[ 1 - \left( \frac{r_{\text{inv}}}{a} \right)^2 \right]
\]  

(5.6)

In Ohmic heated plasmas without additional heating, it is commonly observed that the sawtooth inversion radius increases with the inverse of the edge safety factor \(r_{\text{inv}} / a \approx 1 / q_a\) [2, 26]. With the identification \(r_{\text{inv}} = r_1\) the central safety factor becomes \(q_0 \approx q_a / (q_a + 1)\) and the natural profile takes the simple form [19, 26]:

\[
\frac{p(r)}{p_0} = \left( \frac{j(r)}{j_0} \right) = \left[ 1 + q_a \left( \frac{r}{a} \right)^2 \right]^{-2}
\]  

(5.7)

The shape of these Ohmic profiles then only depends on \(q_a\).

At the plasma edge \(j(a)\) and \(p(a)\) have a value of only a few percent of the central value. For instance, at \(q_a = 4\) the edge pressure is only 4% of the central pressure. An edge pressure of typically 400 Pa is within the error bars of measured the values. Moreover it is very unlikely to find an exact agreement between models based on turbulence and experiments since at the very edge other loss mechanisms than turbulent heat-diffusivity play a dominant role: charge-exchange losses due to penetrating neutrals or, outside the last closed magnetic surface (LCMS), parallel transport to the limiter or divertor.

### 5.2.3 Coincidence between profile consistency and temperature profile stiffness

The obvious question that arises is whether these two models, the profile consistency model and the critical temperature-gradient transport model which are based on different theories (electrostatic turbulence and minimal free energy or statistical current distribution), are contradicting, or are a different representations of the same effects.

The inductively driven current in a tokamak is a strong function of the electron temperature through the Spitzer conductivity [27]:

\[
j = \sigma E_{\text{tor}}
\]  

(5.8)

with \(E_{\text{tor}}\) the toroidal electric field and \(\sigma\) the electrical conductivity of the plasma:

\[
\sigma = T_e^{3/2} f(\epsilon, Z_{\text{eff}}, v_e^*)
\]  

(5.9)
 Chapter 5. Pressure profile consistency in TEXTOR

The $j$-profile is to a certain extent modified by neo-classical effects such as the bootstrap current and enhanced resistivity due to particle trapping. For small pressure gradients and large aspect ratio ($R/r \gg 1$), the bootstrap current contribution is negligible. The enhanced resistivity is described by $f(\epsilon, Z_{\text{eff}}, \nu_e^*)$ which depends on the inverse aspect ratio ($\epsilon = r/R$), the effective ion charge $Z_{\text{eff}}$ and the electron collisionality $\nu_e^*(T_e, n_e, Z_{\text{eff}}, q, \epsilon, R)$. If the $j$-profile depends weakly on the modification $f(\epsilon, Z_{\text{eff}}, \nu_e^*)$, then the $T_e$-profile is well prescribed by profile consistency.

Under certain conditions the $n_e(r)$ and the $p_e(r)$-profiles are also prescribed by profile consistency. The total $p(r)$-profile constitutes the sum of the electron and ion pressures ($p(r) = p_e(r) + p_i(r)$). In the case $p_e(r) \propto p_i(r)$ or $p_e(r) \gg p_i(r)$, the $n_e(r)$-profile is also determined. If one can neglect the neoclassical effects on resistivity and bootstrap current contribution, one can immediately derive a $T_e$-profile from eq. 5.7:

$$\frac{T_e(r)}{T_{e0}} \approx \left[ 1 + q_a \left( \frac{r}{a} \right)^2 \right]^{-4/3}$$ (5.10)

and since $j \propto p \propto n_e T_e$:

$$\frac{n_e(r)}{n_{e0}} \approx \left[ 1 + q_a \left( \frac{r}{a} \right)^2 \right]^{-2/3}$$ (5.11)

The normalized gradients of $j$, $p_e$, $T_e$ and $n_e$ can in this simple case be calculated:

$$-R \nabla \ln(p) = -R \nabla \ln(j) = 4q_a R \frac{r}{a^2} \left[ 1 + q_a \left( \frac{r}{a} \right)^2 \right]^{-1}$$ (5.12)

and

$$-R \nabla \ln(T_e) = \frac{8}{3} q_a R \frac{r}{a^2} \left[ 1 + q_a \left( \frac{r}{a} \right)^2 \right]^{-1}$$ (5.13)

$$-R \nabla \ln(n_e) = \frac{4}{3} q_a R \frac{r}{a^2} \left[ 1 + q_a \left( \frac{r}{a} \right)^2 \right]^{-1}$$ (5.14)

These normalized gradients all have a maximum at $r_{\text{max}}/a = [(q_a/q_0) - 1]^{-1/2}$ (for Ohmic cases $r_{\text{max}}/a \approx q_a^{-1/2}$). The values at this maximum are:

$$-R \nabla \ln(X)_{r_{\text{max}}} = 0.5 C_x \frac{R}{a} \left( \frac{q_a}{q_0} - 1 \right)^{1/2}$$ (5.15)
with \( C_x = \{4, 4, 8/3, 4/3\} \) respectively for \( x = \{p, j, T_e, n_e\} \). Eq. 5.12 is depicted in the right panel of figure 5.2 for 3 values of \( q_a \). The corresponding \( j \)-profiles are shown in the left panel. Shown in this same picture is the \( j \)-profile, which is derived from the \( T_e \)-profile that is shown in figure 5.2, taking into account the neoclassical effects on conductivity. For the derived profile it is found that \([(q_a/q_0) - 1] = 6\). If one compares the value of \(-R \nabla \ln(j)\) at \( r/a = 0.4 \) between the two models, than one can conclude that in the stiff zone the two models (green line and black dashed line) give the same maximum value of \(-R \nabla \ln(j) = 14\), however in the non-stiff central zone and edge zone the behaviour is quite different. Note that other choices of parameters \( \{\kappa, \chi_0, \chi_s\} \) for a critical temperature gradient model will yield different \( T_e \)-profiles and related \( j \)-profiles with different \([(q_a/q_0) - 1] \). A clear difference in the critical temperature gradient model and the profile consistency model is the large variability in the profile shapes that is allowed in a critical gradient model. In [5] a modification of the critical \( T \)-gradient transport model is used which yields high transport when the profiles deviate from the natural profiles prescribed by profile-consistency. Consequently the profile shapes are close to the natural profile and their logarithmic gradients have a \( q_a \)-dependence.

### Figure 5.2
Comparison of shapes from natural \( j \)-profiles (\( j_N \) solid lines with \( q_a = \{2, 4, 6\} \)) with the shape of a \( j \)-profile predicted derived from a critical gradient model (\( j_{cr} \) thick dashed line) which includes neoclassical resistivity enhancement factor (\( f \) dashed thin line). The natural \( j \)-profile shapes become broader for lower \( q_a \). The natural profiles have a maximum at \( r/a = 1/\sqrt{q_a} \) with the maximum proportional to \( \sqrt{q_a} \). In the right panel the logarithmic gradients of the profiles are shown. The shapes (logarithmic gradients) of \( j_N \) and \( j_{cr} \) differ near the plasma center and near the plasma edge.
5.3 Comparison between models and experiments

5.3.1 Experimental tests of model validity

To distinguish between profile consistency and temperature profile stiffness, further measurements are performed on the TEXTOR tokamak. TEXTOR is a circular limiter device like T-10 and RTP. It is equipped with three different heating methods, ECRH, ICRH and NBI. Additionally, a high resolution Thomson scattering system provides detailed profile measurements of $n_e$ and $T_e$. Whereas this study is not meant to be an exhaustive test of profile consistency, still it is intended to perform some critical tests on the hypothesis of pressure profile consistency:

1. Is the $j$-profile consistency as given by eq. 5.3 also valid for strongly additionally heated plasmas at various $q_a$ values?

2. Is the $T_e$-profile stiffness given by the critical temperature gradient model or by eq. 5.13 also valid for additionally heated plasma?

3. If the answer to test 1 is yes, is then also the $p$-profile conform the $j$-profile for additionally heated plasmas as it has been seen in Ohmic plasmas?

4. What are the answers to the above questions if the additional heating power is deposited off-axis?

5.3.2 Experimental setup

A description of TEXTOR, its heating sources and diagnostics is given in chapter 3 and is summarized here. The limiter tokamak has a major radius $R_0 = 1.75m$ and a minor radius of $a = 0.46m$. The localized ECRH is delivered by a $140GHz$ gyrotron with a maximum plasma heating power of $P_{ECRH} \approx 800kW$ that is deposited with a typical width $\Delta r/a \approx 0.05$ [28]. The neutral beam system acts as the main heating source on TEXTOR. It heats both the electrons and the ions and it is also a particle source. Its deposition width is comparable to the minor radius. Another heating source is provided by ICRH which can heat both electrons and ions. The deposition width of ICRH is also larger than the width of ECRH but somewhat smaller than that of the NBI. The $T_e$ and $n_e$-profiles are measured with a Thomson scattering system [12, 23] which has a high resolution in space and time. During one discharge the Thomson scattering system delivers a burst of $\sim 30$ laser-pulses at a rate of $5kHz$. The light is Thomson scattered along a
vertical chord in the plasma and it is spectrally resolved for 120 positions along this chord. A schematic overview of the measurement positions that are projected on a poloidal cross-section of the TEXTOR tokamak, is shown in figure 5.3. The minimal vertical resolution is $\Delta Z \geq 0.9 \text{cm}$. Several consecutive measurements are averaged to enhance the signal to noise ratio ($S/N$). A further increase in $S/N$ is achieved by spatial averaging of neighboring measurement positions whereby the spatial resolution is reduced.

For comparison of the profiles of different plasma discharges, mappings to a magnetic flux surface coordinate $r$ are made:

$$r^2 = [R_{TS} - R_{LCFS} - \Delta R_{Sh}(r)]^2 + [Z_{TS} - Z_{LCFS}]^2$$  \hspace{1cm} (5.16)$$

The mapping assumes circular shaped nested flux-surfaces of radius $r$ whose centers are displaced with respect to the center of the last closed flux surface ($R_{LCFS}, Z_{LCFS}$). The radial position $R_{LCFS}$ is determined with the density interferometer. The vertical position $Z_{LCFS}$ is determined from the Thomson scattering measurements. The magnetic surfaces experience an outward shift, called the Shafranov shift [27], that is due to the pressure and the plasma current profile. An approximation of the Shafranov shift is $\Delta R_{Sh}(r) = \Delta R_{Sh,0}[1 - (r/a)^2]$, where
Table 5.1: List of shots with plasma parameters, heating powers of NBI, ICRH and ECRH for the measured profiles. The last four columns show the time of Thomson scattering measurements relative in the sawtooth cycle, the sawtooth inversion radius measured with Thomson scattering and the central safety factors $q_0$ calculated for Ohmic ($q_{0,\Omega}$) heating and $q_0$ derived from eq. 5.6

<table>
<thead>
<tr>
<th>Shot#</th>
<th>$q_a$</th>
<th>$P_{NBI}$ [kW]</th>
<th>$P_{ICRH}$ [kW]</th>
<th>$P_{ECRH}$ [kW]</th>
<th>$t_{TS} - t_{saw}$</th>
<th>$r_{inv}$ / $a$</th>
<th>$q_{0,\Omega}$</th>
<th>$q_0$</th>
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<td>6.3</td>
<td>1040</td>
<td>1430</td>
<td>0</td>
<td>1.0</td>
<td>0.18</td>
<td>0.86</td>
<td>0.82</td>
</tr>
<tr>
<td>106225</td>
<td>6.3</td>
<td>1025</td>
<td>0</td>
<td>410</td>
<td>1.0</td>
<td>0.22</td>
<td>0.86</td>
<td>0.73</td>
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<tr>
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<td>6.3</td>
<td>1040</td>
<td>0</td>
<td>540</td>
<td>1.0</td>
<td>0.20</td>
<td>0.86</td>
<td>0.78</td>
</tr>
<tr>
<td>106230</td>
<td>3.9</td>
<td>0</td>
<td>0</td>
<td>820</td>
<td>0.4</td>
<td></td>
<td>0.80</td>
<td></td>
</tr>
<tr>
<td>106231</td>
<td>3.9</td>
<td>0</td>
<td>0</td>
<td>500</td>
<td>1.0</td>
<td>0.31</td>
<td>0.80</td>
<td>0.69</td>
</tr>
<tr>
<td>106232</td>
<td>3.9</td>
<td>1060</td>
<td>970</td>
<td>0</td>
<td>1.0</td>
<td>0.34</td>
<td>0.80</td>
<td>0.62</td>
</tr>
<tr>
<td>106235</td>
<td>3.9</td>
<td>1060</td>
<td>730</td>
<td>0</td>
<td>0.5</td>
<td></td>
<td>0.80</td>
<td></td>
</tr>
</tbody>
</table>

$\Delta R_{Sh,0}$ is calculated from the plasma equilibrium parameters. The radial position of the Thomson scattering chord is located at $R_{TS} = 1.84m$. In standard Ohmic discharges with $\beta_p \approx 0.3$ the plasma center is at $R_{LCFS} + \Delta R_{Sh,0} = 1.79m$ and TS thus does not measure at the plasma center. The main inaccuracies in the mapping are due to uncertainties in the determined radial position of the plasma. For the mapping of the Thomson Scattering positions this yields the largest inaccuracies in the plasma center.

5.4 Profile consistency and profile stiffness features from TEXTOR plasmas

5.4.1 Central heated plasmas

A series of discharges with central heating is analyzed on features of profile consistency and $T_e$-profile stiffness. The parameters of these discharges are listed in table 5.1. The discharges have different edge safety factors ($q_a = 3.9$ and $q_a = 6.3$). The centrally deposited heat comes from different sources.

General features of the $T_e$ and $n_e$-profiles

The general features of the $T_e$ and $n_e$-profile measurements are discussed first. In figure 5.4 the $T_e$ and $n_e$-profiles are shown on a logarithmic scale, together
with the logarithmic gradients of the $T_e$ and $n_e$-profiles profiles.

![Graphs of $T_e$, $n_e$, and their logarithmic gradients](image)

**Figure 5.4:** Shapes of the $T_e$, $n_e$, (left column) and their logarithmic gradients (right column). Dashed lines represent $q_a = 6.3$ and solid lines represent $q_a = 3.9$. The spatial resolution of the $T_e$ and $n_e$-profiles in the left panels is $\Delta r/a = 0.02$. The logarithmic gradients have a spatial resolution of $\Delta r/a = 0.17$. The vertical bars at the top of the figures indicate the positions of the sawtooth inversion radii.

The $T_e$ and $n_e$-profiles are the time-average of $\sim 10$ profiles just before the sawtooth crash. The spatial resolution of the $T_e$ and $n_e$-profiles is $\Delta r/a = 0.02$. The $S/N$ of the logarithmic gradients of the $T_e$ and $n_e$-profiles is further enhanced by averaging over 8 neighboring points, whereby the spatial resolution is reduced to $\Delta r/a = 0.17$. The profiles shown for $r/a > 0.7$ are extrapolations of the central $T_e$ and $n_e$-profiles because the plasma edge measurements are affected by Thomson scattering stray light.
Chapter 5. Pressure profile consistency in TEXTOR

The upper left panel shows the $T_e$-profiles on a logarithmic scale. The amplitudes of the profiles are different due to different heating conditions. Logarithmic gradients of the $T_e$-profiles are shown in the upper right panel. Profiles with the same shape have equal logarithmic gradients. The logarithmic gradients equal zero in the plasma center and they increase further outwards. In the region $0 < r/a < 0.5$ two groups of logarithmic gradients are identified. The two different groups correspond to the two different $q_a$. The $T_e$-profiles from the group with $q_a = 6.3$ have larger logarithmic gradients than $T_e$-profiles from the group with $q_a = 3.9$, which means that the $T_e$-profiles from the former group are more peaked. At $r/a = 0.5$ the logarithmic gradients of the different profiles approximately coincide with $-R \nabla \ln T_e = 11$.

The $n_e$-profiles show a relatively larger variation than the $T_e$-profiles. Hollow $n_e$-profiles are observed for discharges with on-axis ECRH (106230 and 106231). No hollow central $n_e$-profiles are observed for the discharges 106232 and 106235 which both have $q_a = 3.9$ but have different central heating (NBI+ICRH). The logarithmic gradients of $n_e$-profile are nearly equal at $r/a = 0.5$ with $-R \nabla \ln(n_e) \approx 4$. This is approximately half the logarithmic gradient of the $T_e$-profile which is compatible with the profile consistency predictions if neo-classical resistivity effects can be neglected.

The profile shapes in the plasma center are affected by the sawtooth crash that causes a periodic flattening of the central profiles. In figure 5.5 the flattening of the central $T_e$ and $n_e$-profiles due to the sawtooth crash in discharges 106220 and 106232 is shown. The discharges have different $q_a$ but comparable heating power levels (see table 5.1). The different $q_a$ of the discharges yield different inversion radii. The sawtooth crash causes a flattening of the central $T_e$ and $n_e$-profiles to just beyond the sawtooth inversion radius. The change in profile shape is reflected in the change in the logarithmic gradient which is reduced by the crash to just outside of the inversion radius. Well beyond the sawtooth inversion radius the profile shape is little affected by the sawtooth crash. The appearance of the sawtooth crash may prevent that the central profile shape ($r < r_{inv}$) can reach a steady state.

Profile consistency features

The questions regarding profile consistency can now be addressed. The $j$-profiles and the $p_e$-profiles are derived from the measured $T_e$ and $n_e$-profiles. The $j$-profiles are calculated using a flat $Z_{eff}(r)$-profile that fits the constraints of total
5.4 Profile consistency and profile stiffness features from TEXTOR plasmas

Figure 5.5: Change in $T_e$ and $n_e$ profile shapes due to the sawtooth crash for discharges with $q_a = 3.9$ (106232, yellow) and $q_a = 6.3$ (106220, blue). The inversion radius is determined from the intersection of the flat central $T_e$-profile after the crash and the steep $T_e$-profile before the crash. The different $q_a$ yield different inversion radii.

plasma current and loop voltage. The normalized $j/j_0$ and $p_e/p_{e,0}$-profiles and the logarithmic gradients of the $j$ and $p_e$-profiles are shown in figure 5.6. Two different groups of nearly coinciding $j/j_0$-profiles and $p_e/p_{e,0}$-profiles are identified that correspond to the different $q_a$. The profiles of the group with $q_a = 6.3$ are more peaked than the profiles of the group with $q_a = 3.9$.

The logarithmic gradients of the profiles are compared with logarithmic gradients of the natural $p(r)$ and $j(r)$-profiles according to equation 5.12. The positions of the maxima of the logarithmic gradients of the natural profiles are indicated with colored error bars in the figure. The size of the error bars is derived from the photon statistics of Thomson scattering and asymmetries in the $T_e(z)$ and $n_e(z)$-profiles. The maxima of both the logarithmic gradients of the natural $j$ and $p$-profiles coincide with the logarithmic gradients of the measurements.

The dashed and solid grey lines indicate the average logarithmic gradients of the natural profiles according equation 5.12, with respectively $(q_a = 6.3, q_0 = 0.78)$ and $(q_a = 3.9, q_0 = 0.66)$. Logarithmic gradients profiles from both groups of $q_a$ nearly coincide with the average logarithmic gradients for the given errors.
Chapter 5. Pressure profile consistency in TEXTOR

5.4.2 Off-axis heating

The effect of different central heating methods on the profile shape has been analyzed in the previous section. The analysis of the effects of central heating on the central profile shape is however hampered by the sawtooth instability. By application of off-axis ECRH well outside of the sawtooth inversion radius it is aimed to reduce the effect of the sawtooth instability on the profile shapes. Off-axis...
ECRH will cause a broadening of the $T_e$-profile and consequently a broadening of the conductivity profile and the $j$-profile. The profile consistency model and the critical $T_e$-gradient model do however predict different behavior for off-axis ECRH.

From a profile consistency point-of-view one can expect that the $j(r)$ and $p(r)$-profiles remain close to what is predicted by profile consistency (eq. 5.7). For a given $q_a$ the broadening of the shape of the natural $j$ and $p$-profiles can then only be achieved by an increase of $q_0$ which can be observed in a reduction of sawtooth inversion radius $r_{\text{inv}}$.

From the critical temperature gradient model one can expect that off-axis ECRH increases the electron heat-flux and causes stiff $T_e(r)$-profile outside the deposition radius. Inside the heat deposition radius the $T_e$-gradient can be below the critical gradient if the residual heat-flux (from e.g. Ohmic heating) is low enough. For off-axis heating, a critical gradient model thus allows a much larger variety of profiles than profile consistency.

A comparison is made between discharges with off-axis ECRH and Ohmic discharges. The parameters of the discharges are listed in table 5.2. The response of the near central $T_e$ and $\langle n_e \rangle$ to off-axis ECRH is shown in figure 5.7. The off-axis ECRH, that is applied from $t = 2.000\text{s}$ to $t = 2.200\text{s}$, is more than a factor two larger than the total Ohmic heating power ($P_\Omega 200 \pm 100\text{kW}$). The off-axis ECRH causes a rise in both the near central $T_e$ and central $\langle n_e \rangle$. The $T_e$ time trace exhibits MHD crashes after $\sim 30\text{ms}$ of ECRH (this effect has been found to be reproduced in several discharges).

A comparison between the Ohmic discharge and the off-axis ECRH discharge 25$\text{ms}$ after switch-on, is shown in the figure 5.8. The $T_e$-profile of discharge 108942 shows that an off-axis maximum has formed at $r/a = 0.26$ after 25$\text{ms}$ of off-axis ECRH. The $T_e$-maximum is somewhat closer to the plasma center than the calculated deposition radius at $r_{\text{ECRH}}/a = 0.30$, but it is still outside the measured (Ohmic) inversion radius $r_{\text{inv}}/a = 0.20$. Outside of the heat deposition location the logarithmic $T_e$-gradient is equal to that of the Ohmic plasma. This shows that $T_e$-

### Table 5.2: List of shots with plasma parameters, for off-axis ECRH.

<table>
<thead>
<tr>
<th>Shot#</th>
<th>$q_a$</th>
<th>$P_{\text{ECRH}}$ [kW]</th>
<th>$\Delta t_{\text{ECRH}}$ [ms]</th>
<th>$r_{\text{inv}}/a$</th>
<th>$r_{\text{ECRH}}/a$</th>
</tr>
</thead>
<tbody>
<tr>
<td>108931</td>
<td>4.7</td>
<td>0</td>
<td>0</td>
<td>0.21</td>
<td>0.30</td>
</tr>
<tr>
<td>108942</td>
<td>4.7</td>
<td>600</td>
<td>25</td>
<td>0.21</td>
<td>0.30</td>
</tr>
</tbody>
</table>
Chapter 5. Pressure profile consistency in TEXTOR

Figure 5.7: Measurement times Thomson Scattering in reproducible discharges. The upper panel shows time dependent response of 135GHz ECE-channel measuring just inside the sawtooth inversion radius. The central ($R = 1.75m$) line-averaged densities are shown in the lower panel. The Thomson scattering times are represented by the vertical dashed lines with colors that correspond to the colors of the $T_e$ and $n_e$ time traces.

profile shape is little affected by the additional ECRH power ($-R\nabla \ln(T_e) = 11\pm3$ for $0.4 < r/a < 0.7$). At the off-axis $T_e$ maximum the logarithmic $T_e$-gradient is reduced compared to the Ohmic case: $-R\nabla \ln(T_{e,ECRH}/T_{e,\Omega} = -8)$.

The $n_e$-profile shape is less affected by the off-axis ECRH than the $T_e$-profile shape. Over the entire profile $R|\nabla \ln(n_{e,ECRH}/n_{e,\Omega})| < 2$. The resulting $p_e$-profile and its logarithmic gradient are shown in figure 5.9. The local modification of the $T_e$-profile and the lack of modification of the $n_e$-profile cause that also the $p_e$-profile is locally modified by off-axis ECRH. The $p_e$-profile also shows an off-axis maximum at maximum at $r/a = 0.26$. The MHD crash cause the disappearance of the off-axis maximum and the appearance of a tearing mode. A comparison of the profiles without island is complicated because the tearing mode causes a local two-dimensional modification of the $T_e$ and $n_e$-profiles.

The $j(r)$-profiles, shown in figure 5.9, are derived from the measured $T_e(r)$ and $n_e(r)$-profiles, assuming a steady state neo-classical conductivity. The $j$-profile calculation for the Ohmic plasma phase was found to converge with a realistic value of $Z_{eff} > 1.3$. For the plasma with off-axis ECRH the profile was only found to converge with the unrealistic low value $Z_{eff} < 1$. The $j$-profile shown has $Z_{eff} = 1$ which consequently yields a too low plasma current. The missing current is driven by other mechanisms such as the generation of supra-thermal
5.5 Summary and conclusions

In this chapter the question was addressed to which extent local heating and current drive is an effective tool to locally modify current and temperature profiles.

electrons by ECRH that enhances the conductivity. At the onset of the MHD-crashes, the current-profile may not have reached a steady state diffusive profile but the appearance of the crashes is a clear indication that the current profile has changed so that instabilities can grow causing the growth of the modes that cause the crashes.

Figure 5.8: Comparison of $T_e$ and $n_e$-profiles and their logarithmic gradients for off-axis ECRH and Ohmic heating. The calculated ECRH deposition location is at $r/a = 0.30$. The sawtooth inversion radius as determined from $T_e$-profile for the Ohmic discharge is indicated at $r/a = 0.21$. The ECRH deposition location is at $r/a = 0.30$. The black dashed vertical line indicates the radius beyond which an exponential profile is fitted with a function of constant logarithmic gradient. The profiles are averaged over two radial neighboring points yielding a radial resolution $\Delta r/a = 0.04$. 

5.5 Summary and conclusions

In this chapter the question was addressed to which extent local heating and current drive is an effective tool to locally modify current and temperature profiles.
Chapter 5. Pressure profile consistency in TEXTOR

Figure 5.9: The Effect of off-axis ECRH on the $p_e$-profile and the $j$-profile shape in comparison to profiles of purely Ohmic heated plasmas and natural profile shapes (grey). The $j(r)$-profiles are calculated assuming neo-classical resistivity.

The motivation for this investigation is that these tools are foreseen for the control of MHD-modes in fusion reactors such as ITER. In particular, they will be used to control the size of magnetic islands such as neo-classical tearing modes, and the amplitude and period of the sawtooth instability. Local transport, i.e. diffusion of heat and current density, will affect the efficiency of such control schemes. In the following chapters the two-dimensional aspects, mode structure and 2D transport in and around islands, are investigated. In the present chapter it was investigated to which extent a possible non-linear response of the temperature profile to local heating, or a non-local response, would affect the possibilities to control the MHD-modes. In particular, two classes of models were considered: ‘profile consistency’ and a ‘critical gradient model’. The former model is essentially non-local, where the transport adapts in such a way that a particular ‘preferred’ profile shape is maintained. In the latter model the local transport coefficients are
a function of the temperature gradient, which leads to a reduced response of the profile to local heating.

Using the combination of heating methods and high-resolution diagnostics available on TEXTOR, several experiments were done to test whether local heating results in local adaptation of the profiles. Experiments with central heating were not conclusive. The results were compatible with profile consistency as well as a critical gradient model, but it should be added that this is not a sensitive test of these models. In fact, with central heating the profiles are dominated by the sawtooth instability, which flattens the central profiles out to the inversion radius, leaving little freedom for the profile to evolve. For such plasmas profile consistency does give a fair description of the profile shapes of electron temperature, electron density, electron pressure and plasma current. The critical temperature gradient model, too, is capable of reproducing these profiles, but with the number of free parameters in the model this was to be expected.

More testing were experiments with off-axis electron heating. Local off-axis ECRH deposition was found to cause a pronounced local effect on the local electron temperature profile shape with the creation of an off-axis maximum. No modification of the local electron density profile was observed so that the resulting electron pressure was not conserved. Outside the ECRH deposition zone the $T_e$-profile shape is not affected, which is a behavior that can be well described by critical temperature gradient model. In conclusion, these experiments confirmed that off-axis heating can indeed be used to locally modify temperature and pressure profiles.

References


Chapter 5. Pressure profile consistency in TEXTOR


Chapter 5. Pressure profile consistency in TEXTOR


Heat pulse propagation studies around magnetic islands induced by the Dynamic Ergodic Divertor in TEXTOR

Abstract Since the efficiency of the tearing mode suppression by heating depends on the electron heat diffusivity it is important to know if the electron heat transport coefficients inside the island are reduced compared to the ambient plasma. With that aim, modulated ECRH has been employed for heat pulse propagation studies in and around magnetic islands at the TEXTOR tokamak. The combination of its special hardware tools of the Dynamic Ergodic Divertor to generate tearing modes, the ECRH system for producing heat pulses and the ECE-Imaging diagnostic for its analysis offered a direct view of the perturbed two dimensional heat flow in around the magnetic island. Inside $m/n = 2/1$ and $m/n = 3/1$ islands with a flattened temperature profile, the electron heat transport is shown to be strongly reduced with respect to the surrounding plasma. Inside the islands a heat pulse diffusion coefficients $\chi_e \approx 0.4 m^2/s$ was derived, while outside the island it is an order of magnitude larger $\chi_e > 3 m^2/s$. In contrast, power balance calculations of strongly heated islands shows that the electron transport is similar to the surrounding plasma. These results suggest that the heat transport inside a magnetic island is also governed by a critical gradient like behaviour, similar to the bulk plasma.*

*The work presented in this chapter is published in this form in Nuclear Fusion under the title Heat pulse propagation studies around magnetic islands induced by the Dynamic Ergodic Divertor in TEXTOR. G.W. Spakman et al., Nucl. Fusion 48, 115005 (2008)
6.1 Introduction

The formation of tearing modes in fusion plasmas is often observed to degrade the plasma energy confinement. With such a tearing mode, an island structure is formed which consists of locally nested flux surfaces. The cross field transport in and around magnetic islands is important for the stability and the evolution of (neoclassical) tearing modes. For high temperature plasmas the heat transport parallel to the field on the flux surfaces is orders of magnitude larger than the transport perpendicular to the field. This causes the island boundary to act as an effective short-circuit for heat transport across the island. In the absence of heat sources inside the island this leads to the flattening of the temperature profile over the island, resulting in the degradation of total plasma energy confinement. Profile flattening in high pressure regimes induces a loss in the bootstrap current over the island which increases the island growth resulting in neoclassical tearing modes [16].

It has been shown that, depending on the finite ratio of parallel to perpendicular transport, only islands of sufficient width will lead to flattening of the profiles [9]. The cross field transport inside an island also determines the possible effect of a local heat or particle source. The electron temperature peaking inside an island induced by a local heat source has been shown to have a stabilizing effect on a tearing mode [3, 22].

Results on measurements of transport coefficients inside the island are ambiguous. On the one hand, data from various machines indicate that they are lower compared with the ambient plasma [4, 15, 21]. Recent perturbative transport experiments by means of cold pulse propagation at the LHD heliotron [10] showed that the electron heat diffusion coefficient $\chi_e$ inside an $m/n = 1/1$ island is reduced by an order of magnitude with respect to the $\chi_e$ of the surrounding plasma. Here $m$ is the poloidal mode number and $n$ is the toroidal mode number. On the other hand, power balance analysis of recent experiments on island suppression at the TEXTOR tokamak showed that $\chi_e$ of large $m/n = 2/1$ islands is comparable to $\chi_e$ of the bulk plasma, in case there is a large temperature gradient due to strong heating inside the island [3].

The power balance analysis that yields the electron heat diffusivity $\chi_{e}^{(pb)} = -q_e/n_e \nabla T_e$ depends on measurements of the electron heat flux $q_e$, the electron temperature gradient $\nabla T_e$ and the electron density $n_e$. Both $q_e$ and $\nabla T_e$ are generally hard to determine inside an island which results in inaccurate estimates of the electron heat diffusivity. Another method is perturbative transport
analysis which relies on the analysis of transient perturbations. With electron heat pulse propagation analysis the perturbed electron heat diffusivity \( \chi^{(hp)}_{e} \) can be determined from the perturbed temperature gradient due to the perturbed heat flux. The measured quantity is different in nature and it can yield additional information [2]. In the flat temperature region of an island the perturbative analysis can yield accurate results if the temporal temperature evolution can be accurately measured [10].

In tokamaks the perturbative transport analysis of islands is usually complicated by the natural rotation of the island with respect to the measurement frame. However, the TEXTOR tokamak is equipped with a set of external perturbation field coils called the Dynamic Ergodic Divertor (DED) [7] which can induce and lock magnetic islands. Under the application of the DED islands with mode numbers \( m/n = 2/1 \) and \( m/n = 3/1 \) have been observed [13].

The aim of this paper is to measure directly if there is a difference in the perpendicular electron heat diffusion coefficients inside and outside islands of a tokamak plasma.

The difference in transport inside and outside the different islands has been investigated by means of heat pulse propagation. The results for both \( m/n = 2/1 \) and \( m/n = 3/1 \) islands are compared. The heat pulses are generated by means of modulated electron cyclotron resonance heating (MECRH) [23] and the response of the electron temperature profile to the heat pulses is measured with a two-dimensional electron cyclotron emission imaging system (ECE-Imaging) [5].

In section 6.2 the experimental setup for the generation of islands, the heat pulse generation and the measurements is described. The derivation of the transport coefficients for island geometry is described in section 6.3. Experimental results of the heat pulse propagation in and around a large \( m/n = 2/1 \) island are presented in section 6.4, followed by the heat pulse propagation in and around a large \( m/n = 3/1 \) island in section 6.5. A discussion and conclusions are given in section 6.6.

6.2 Experimental setup

The TEXTOR tokamak (circular limiter tokamak, \( R_0/a = 1.75m/0.46m \); toroidal field \( B_t < 3T \); plasma current \( I_p < 800kA \)) is equipped with a controllable external perturbation field, the DED. The DED consists of 16 helical coils that are poloidally localized on the inboard high field side of the vessel [7]. For the ex-
experiments reported in this paper the DED is operated in the $m/n = 3/1$ mode configuration. This mode of operation causes large resonant magnetic perturbations at the rational magnetic surfaces with toroidal mode numbers $n = 1$ and low poloidal mode numbers $m$. The mode spectrum of the DED is well visualized in vacuum field calculations. The vacuum magnetic field is defined as the superposition of the applied DED perturbation field and the unperturbed plasma equilibrium field. Figure 6.1 shows the contours of the major $n = 1$ magnetic islands in the form of a Poincaré mapping of the vacuum field. The Poincaré

![Diagram](image)

Figure 6.1: Poincaré mapping of the DED induced vacuum magnetic field at the poloidal cross section of the ECE-Imaging observation area. Indicated are the deposition radius of ECRH (solid circle, $R = 162cm$) and the ECE-Imaging observation area (deposition at three different radial settings, rectangles). The O-point of the $m/n = 2/1$ island and the X-point region of the $m/n = 3/1$ island are near the bottom of the ECE-Imaging observation area. The O and X-point positions of the $n = 1$ islands become swapped in the vacuum field calculations by reversal of both the toroidal field and plasma current.

mapping is calculated with the field line tracing Gourdon code. A good correlation between experimental plasma edge structures and the modelled magnetic topology, using the Gourdon code, was revealed in [17]. For large excited tearing modes that are located well inside the plasma and are locked to an external per-
6.2 Experimental setup

turbation field, the phases of the island O-points are predicted to correspond to those of the vacuum field [8]. The actual island width also depends on the classical tearing mode stability parameter $\Delta'$ and other effects such as mode coupling with other modes in the plasma. Therefore the measured island width is likely to be different from the vacuum field island width [8]. When a large $m/n = 2/1$ island is excited in TEXTOR, $n = 1$ modes at different rational surfaces can become coupled to this island. In particular it is often found that the $m/n = 1/1$ internal kink mode locks to the $m/n = 2/1$ island and that the sawtooth crash cycle stops. With the locking, the toroidal rotation profile between the $q = 1$ and the $q = 2$ surfaces flattens [1]. It was previously found that the $m/n = 3/1$ island becomes excited at a higher DED perturbation amplitude than the $m/n = 2/1$ island [13]. For certain plasma conditions, as described in this paper, an $m/n = 3/1$ island becomes excited at a lower DED perturbation field amplitude than needed for the excitation of the $m/n = 2/1$ island. The heat transport properties of this island are also investigated with and without the $m/n = 2/1$ island being present at the same time.

For the generation of electron temperature perturbations localized MECRH is applied. The heat is deposited by means of a 140 GHz gyrotron which heats the electrons at the second harmonic X-mode [23]. The MECRH deposition location is near the equatorial midplane at the high field side of the magnetic axis. The typical radial deposition width (FWHM) is $1 - 2 cm$. The vertical deposition location is set with the ECRH launcher. The temperature perturbations are measured at the low field side with lower frequency ECE systems. With a high resolution 2D ECE-Imaging system which employs 128 ECE channels (16 vertical $\times$ 8 horizontal) [5] the propagation of the heat pulses is measured in the vicinity of the islands. The system has an observation area of 17 cm (vertical) $\times$ 10 cm (radial) which corresponds to 10% of the poloidal angle at the $q = 2$ surface and 20% of the minor radius. Different parts of the plasma minor radius, indicated in figure 6.1, are covered in reproducible discharges by tuning the local oscillator of the system. Additional measurements are made with a conventional 1D 11-channel ECE system which covers a larger radial range [20]. The ECE systems are calibrated with a high resolution Thomson scattering system [18]. The additional diagnostics that are used for the measurement of some basic plasma parameters are described in [6]. The vacuum field calculations show that with normal toroidal field, plasma current and DED configuration, the O-point of the $m/n = 2/1$ island and the X-point region of the $m/n = 3/1$ island are near the bottom of the ECE-Imaging observation area (see figure 6.1). To compare the O-point with the
X-point region of an \( n = 1 \) island a toroidal phase shift of the plasma column by \( \pi \) rad is required. The static dc operation of the DED practically only allows a toroidal phase shift of \( \pi/2 \) rad. For \( n = 1 \) islands an interchange of O- and X-point positions can also be achieved by reversal of both the directions of the toroidal magnetic field and the plasma current.

### 6.3 Derivation of \( \chi_e \) inside and outside the island

The propagation of heat pulses is determined by diffusive as well as non-diffusive processes. In [11] analytical expressions relating the transport coefficients to the propagation of periodic temperature perturbations have been derived incorporating non-diffusive effects such as damping, density gradients and cylindrical geometry. In this section the effect of the island geometry on the heat pulse propagation inside the island is described and compared with the propagation outside the island. The propagation of the temperature perturbations is modeled as a diffusive process. The heat flux is taken to be governed by the energy balance equation

\[
\frac{3}{2} \partial_t (n_e T_e) + \nabla \cdot \vec{q}_e - S_e = 0 \tag{6.1}
\]

with \( T_e \) the perturbed electron temperature, \( q_e \) the perturbed electron heat flux and \( S_e \) the perturbed electron heating power density which constitutes local electron heating and loss terms. The perturbed electron heat flux is taken to be governed by perturbed temperature gradient driven diffusion

\[
\vec{q}_e = -n_e \chi_e \nabla T_e \tag{6.2}
\]

with \( \chi_e \) the local heat pulse diffusivity. The perturbed electron temperature and density are taken to be flux surface functions in the island vicinity. If the temperature is constant on the flux surfaces, then the temperature gradient varies on a flux surface because the distance between the flux surfaces varies on a flux surface. To account for this equation 6.1 is averaged over the flux surfaces using the following flux surface averaging operation:

\[
\langle ... \rangle_S = \frac{1}{V'(r)} \frac{\partial}{\partial r} \left( \int_{V(r)} \langle ... \rangle dV \right) \tag{6.3}
\]

where \( r \) is a flux surface coordinate, \( V(r) \) is the volume enclosed within the flux surfaces and \( V'(r) = dV(r)/dr \). This operation does not affect the first term of
6.3 Derivation of $\chi_e$ inside and outside the island

equation 6.1 since it is taken constant on the flux surfaces. The last term is set zero outside the deposition area of the perturbed heat source. For the flux surface averaged divergence it is assumed that $\chi_e$ is constant on the flux surfaces. Using Gauss’ theorem the second term of equation 6.1 takes the form

$$\langle \nabla \cdot \vec{q}_e \rangle_S = -\frac{1}{V'(r)} \frac{\partial}{\partial r} \left( n_e \chi_e \frac{\partial T_e}{\partial r} \int_{\partial V(r)} \nabla r \cdot dS \right)$$

(6.4)

The flux surface averaged divergence depends on the flux surface geometry and can be calculated. For the coordinate $r$ the distance along the minor radius is taken which corresponds to the radial direction in which the plasma is observed. Inside the island $r$ is the defined as the distance to the island O-point. Outside the island $r$ is defined as the distance to the plasma centre. Inside the island around the O-point the flux surface geometry is strongly elliptical. The calculated divergence of the heat flux for strongly elongated ellipses is

$$\langle \nabla \cdot \vec{q}_e \rangle_S = -\frac{1}{2} \frac{\partial}{\partial r} \left( n_e \chi_e \frac{\partial T_e}{\partial r} r \right)$$

(6.5)

This form of the divergence operator corresponds to a cylindrical geometry with a larger effective radius $r_i = \sqrt{2} r$. Close to the island boundary, the separatrix, the flux surfaces deviate from the ellipse-like geometry and $V'(r)$ becomes large compared with the flux surface integral of $\nabla r$. This will result in a local increase in the effective minor radius. Outside the island the effective geometry is close to a circular cylindrical. There the average divergence of the electron heat flux (Equation 6.4) is a factor 2 larger.

The MECRH generates periodic temperature perturbations in the plasma and can be written in the form of a Fourier series

$$T_e = \langle T_e \rangle_t + \sum_{k=1}^{\infty} A_k(r) \exp [i2\pi k ft - i\phi_k(r)]$$

(6.6)

General expressions relating the gradients of the phase delay $\phi$ and the relative gradient in the amplitude $A'/A = \ln(A)'$ of the different harmonics $k$ of the modulation frequency $f$ to perturbative transport coefficients can be found in [11]. In the limit of high modulation frequencies the propagation of the heat pulses is dominated by diffusion. Assuming constant heat diffusivity inside the island, the following relation holds for diffusive heat pulse propagation near the island
where $\Delta \tau = \Delta \phi/2\pi f$ is the time difference with respect to the O-point. For a larger radius a more general analytical form can be used which takes into account effects due to damping, cylindrical geometry and density gradients [11].

$$\chi_e = \frac{3}{8} \frac{r_i^2}{\Delta \tau} = \frac{3}{4} \frac{r_i^2}{\Delta \tau}$$  \hspace{1cm} (6.7)

where $n'/n$ is the relative gradient of the density. Large gradients in both the phase delay and logarithmic amplitude correspond to low diffusive transport coefficients and shallow gradients to high heat diffusion coefficients. In the limit of high modulation frequencies, diffusion dominates the propagation of perturbations which causes $\phi' = A'/A$ and the effect of geometry and gradients in density to become negligible.

6.4 Heat pulse propagation in and around a large $m/n = 2/1$ island

6.4.1 Experimental results $m/n = 2/1$ island

Typical parameters of a TEXTOR discharge used to excite a large $m/n = 2/1$ island are shown in figure 6.2. A toroidal field of $B_t = 2.3T$ and a plasma current of $I_p = 320kA$ are used, resulting in an edge safety factor $q_a = 4.5$ in the cylindrical approximation. After the discharge reaches a quasi-steady state, the DED perturbation field is increased by ramping up the currents in the DED coils as is shown in figure 6.2a. At a critical current of $0.8kA$ a large $m/n = 2/1$ island is excited. This is followed by the disappearance of the sawtooth oscillation as observed in the central $T_e$. The disappearance is attributed to the locking of the $m/n = 1/1$ mode to the $m/n = 2/1$ mode [1]. The line-averaged density is kept at a constant value of $\langle n_e \rangle = 2 \times 10^{19}m^{-3}$ by feed back control. After the plasma reaches a steady state again, MECRH is applied from 1.8 to 2.8 s with a high power level of $540kW$ and a low power level of $140kW$. A modulation frequency $f_{mod} = 40Hz$ and a duty cycle of 50% are found to achieve large enough a temperature modulation for the first harmonic well inside the $m/n = 2/1$ island. A series of reproducible discharges are made in which the radial position
6.4 Heat pulse propagation in and around a large $m/n = 2/1$ island

Figure 6.2: Scenario of a discharge where a large $m/n = 2/1$ is generated. (a) Amplitude of a DED perturbation coil current, (b) central line-averaged $\langle n_e \rangle$, (c) central $T_e$, (d) heating power of MECRH (solid line), Ohmic (dash dotted line) and neutral beam (dashed line). The time of excitation of the $m/n = 2/1$ tearing mode is indicated by the vertical dashed line.

of the ECE-Imaging observation area is varied over the minor radius in order to cover the island vicinity as indicated in figure 6.1. Three discharges near the $m/n = 2/1$ island O-point (normal $B_t$ and $I_p$), three discharges near the island X-point (reversed $B_t$ and $I_p$) and a reference discharge without island. The time traces of ECE-channels measuring near the deposition location, near the island O-point and outside the island beyond the $q = 2$ radius are shown in figure 6.3. It is found that the temperature modulation inside the island is much smaller than it is outside the island, at both sides of the island.

6.4.2 Fourier analysis $m/n = 2/1$ island

The Fourier transform of single radial rows of ECE-Imaging channels (at the bottom of its observation area) for measurements near the magnetic island O-
Figure 6.3: Temperature response to MECRH near the deposition location ($R = 1.97m$), inside the island near the O-point ($R = 2.06m$), outside the island outside the $q = 2$ radius ($R = 2.12m$). The measured square wave modulated ECRH power is also shown.

The time averaged temperature $\langle T_e \rangle$ profiles are shown in figure 6.4a. The presence of a locked $m/n = 2/1$ tearing mode decreases $\langle T_e \rangle$ near the island’s O-point compared with the situation without the tearing mode. The temperature profile near the magnetic island’s O-point is flattened over a width of $w \approx 11cm$ which is interpreted as the island width. Near the island X-point no clear flattening is observed. Note that the global confinement in the presence of an island is decreased (for instance lower central temperature), which is attributed to the fast parallel transport around the island or at the boundary which does not allow large temperature gradients to be build up on the flux surfaces near the island boundary. The net effect of this is a larger temperature gradient outside the island (see figure 6.4a, green line), leading to increased radial heat transport.
6.4 Heat pulse propagation in and around a large $m/n = 2/1$ island

Figure 6.4: Profiles of (a) Average $T_e$ during power modulation, (b) Phase delay of the first harmonic of $T_e$ and (c) Amplitude of the first harmonic, near the $m/n = 2/1$ tearing mode O-point (green O-symbols), X-point (red X-symbols) and in a reference discharge without island (blue triangles) as measured by a single radial row at the bottom of the ECE-Imaging observation area. For the reference discharge additional data from a 1D ECE-system is included. The orange bar indicates the calculated ECRH-deposition region. The green arrow denotes the island width obtained from the flattening of the temperature profile. The ECE-Imaging measurements near the O-point and the reference were taken at $z = -3\, \text{cm}$ and near the X-point they are taken at $z = -7\, \text{cm}$ and mapped to the equatorial midplane. Vertical black arrows indicate calculated rational surface positions of the unperturbed plasma.

The phase delay profile (figure 6.4b) exhibits a local minimum at the deposition radius ($r_{dep}$) and thus increases away from the deposition region. Without an island the increase is monotonic, both towards the plasma centre and towards the plasma edge which is typical for diffusive propagation [11]. Near the magnetic island X-point the increase in the phase delay is also nearly monotonic. However, near the O-point the phase delay has a pronounced local maximum. The width and the edges of the peak correspond to that of the flattened area of the average $\langle T_e \rangle$. At the outer boundary of the island the phase delays near O- and X-point coincide
again. This shows that the heat pulses are propagating fast along the boundary of the island, while they propagate relatively slow, with a large delay, towards the island's O-point. Outside $R = 2.10m$ the phase delay does not increase any more. This could be due to the presence of a smaller $m/n = 3/1$ island near the plasma edge which is discussed below. Moreover, it has to be noted that the ECE measurements close to the plasma edge contain small suprathermal contributions which decrease the phase delay.

The amplitude of the propagating heat pulses, shown in figure 6.4c, has reversed but similar behaviour as the phase, as is typical for diffusive transport: its maximum is at the deposition radius and decreases in both directions. Without an island the amplitude decay is monotonic. With an island it is nearly monotonic near the X-point. Inside the island a large drop of $A$ is observed whose width and position correspond to that of the peak in $\Delta \phi$ and the width of the flattening $\langle T_e \rangle$ profile.

The local 2D heat pulse propagation is most clearly observed in the phase delay near the island's X-point which is shown in a composite picture (figure 6.5). The red arrow indicates the radial row that is also shown in figure 6.4b. The phase delay has a local maximum in the radial direction at the island position. The height and the width of the local maximum vary in the vertical direction (poloidal angle), being largest at the top and lowest at the bottom of the ECE-Imaging observation area. A dashed white line which connects points with a phase delay of $\sim 75^\circ$ indicates the inner edge of the island. Since the $m/n = 2/1$ island structure displaces the flux surfaces more towards the plasma centre than towards the plasma edge, the outer edge of the island (outer dashed white line) should be close to the $q = 2$ flux surface, as calculated for an unperturbed plasma. Indeed this line is seen to connect points with a nearly equal phase delay of $\sim 75^\circ$, as at the inner edge. Both lines converge towards the bottom of the observation area. A complete vanishing of the maximum as expected for the X-point is not observed. According to the vacuum field calculations, shown in figure 6.1, the X-point is situated at the bottom of the ECE-Imaging observation area. This corresponds well given the accuracy of the measurements.

Inside the deposition radius also a large increase in the phase delay is observed. This is considered to be due to a locked $m/n = 1/1$ island. The vacuum field calculations for this field configuration predict the O-point of the $m/n = 1/1$ mode to be positioned near the equatorial midplane at the low field side of the magnetic axis (see figure 6.1). Note that for the situation of figure 6.5 the islands are poloidally rotated by $180/m$ degrees, due to the reversed current and magnetic
Figure 6.5: 2D phase delay of a tip of the island near the X-point. The image is a composition of three images at adjacent radial positions obtained in reproducible discharges near the estimated $m/n = 2/1$ island X-point position (black double arrows indicate the individual observation areas). The black dots denote evaluated channels which had sufficient signal to noise ratio; the other data points were obtained by means of 2D linear interpolation. The dashed white lines connect points with equal phase delay ($\sim 75^\circ$) with the outer line representing a calculated flux surface near the outer island boundary. The orange line indicates the calculated ECRH deposition radius. The thick red arrow indicates the radial row of the phase-delay near the X-point which is also plotted in figure 6.4b.

field). An alternative suggestion for the increased phase delay towards smaller radii is the decrease in $\nabla T_e$ due to the heat pulse. This however is discarded since the phase delay measurement for normal field and plasma current (when the vacuum field $m/n = 1/1$ X-point area is located at the low field side), do not show a large increase towards the plasma centre. Also in the reference discharge a smaller phase delay is observed.

6.4.3 Heat transport coefficients inside the island $m/n = 2/1$ island

The measurements of the amplitude and the phase delay of the heat pulse propagation across the island O- and X-point indicate that the $T_e$ perturbations are functions of nested closed flux surfaces, both inside and outside the island. The heat transport properties inside the island are different from outside the island. The heat transport around the island through the X-point region is fast in com-
comparison with transport towards the island O-point.

Under the assumption that the transport coefficient inside the island is constant and that the temperature perturbations are flux surface functions the heat pulse diffusion coefficients are derived. The phase delay near the island O-point as shown in figure 6.4b, is taken to be parabolic as expected for diffusive propagation in a cylindrical geometry. Using the cylindrical form of equation 6.7, with $r = w/2$ and $\Delta \phi$ the phase difference between the island edges and the O-point, $\chi_{e}^{(hp)} = 0.4m^2/s$ inside the island is obtained. A typical relative density peaking $\Delta n_e/n_e \approx 15\%$ between island centre and edge as measured with Thomson scattering yields an approximate inverse density gradient length $\Delta n_e/n_e w \approx 1.4m^{-1}$ which is negligible compared with approximate inverse gradient length of the amplitude $2\Delta \ln(A)/w = 26m^{-1}$. A density gradient is therefore neglected in the derivation of the heat diffusivity. The propagation of the heat pulses inside the island is therefore taken to be dominated by diffusion perpendicular to the flux surfaces.

For the plasma without an island a heat pulse diffusivity of $\chi_{e}^{(hp)} = (5 \pm 2)m^2/s$ was derived at the island position, using the slab geometry with cylindrical corrections. Most of the heat is deposited inside the $q=2$ surface. Power balance analysis yields $\chi_{e}^{(pb)} = (3 \pm 1)m^2/s$ around the $q = 2$ radius. The electron heat diffusivity determined by heat pulse propagation is generally found to be equal or higher than that determined by power balance analysis [2].

6.5 Heat pulse propagation in and around an $m/n = 3/1$ island

6.5.1 Experimental results $m/n = 3/1$ island

To corroborate the results obtained for the $m/n = 2/1$ island, the heat transport in and around an $m/n = 3/1$ island is investigated in another set of experiments. The discharge described here has reversed toroidal field and plasma current: $B_t = -2.25T, I_p = -305kA$ and $q_a = 4.6$. This results in a similar $q$-profile and vacuum field topology as for the discharges of the previous section. The radial position of the ECE-Imaging observation area is set to measure from the estimated $q = 2$ surface to the outer plasma edge at $R = 2.21m$. The field line tracing calculations indicate that the O-point of the $m/n = 3/1$ vacuum island is in front of the ECE-Imaging observation area.

During the entire DED-phase MECRH is applied. The deposition location on
the high field side is just inside the \( q = 2 \) surface which corresponds to \( R \approx 2.04m \) on the low field side. A modulation frequency of \( f_{\text{mod}} = 40Hz \) and a duty cycle of 70% are used; these choices cause both the first and second harmonic of \( T_e \) to be above the noise level in the region of interest. A line-averaged density \( \langle n_e \rangle = 3 \times 10^{19}m^{-3} \) and an electron temperature \( T_e = 500eV \) at the deposition location yield an optical thickness of \( \tau \approx 3 \) which ensures a sufficiently high single-pass absorption not to disturb the measurements. Strong tangential beam injection in the counter plasma current direction is applied, which increases the threshold for \( m/n = 2/1 \) island excitation by the DED [12].

Typical discharge parameters are shown in figure 6.6. During the ramp-up of the DED perturbation field current (\( I_{\text{DED}} \), figure 6.6a) the toroidal rotation of the plasma (figure 6.6b) which is measured near the plasma centre and just inside the \( q = 3 \) surface, decreases with increasing \( I_{\text{DED}} \). At \( t = 2.5s \) with \( I_{\text{DED}} = 2.4kA \) the \( m/n = 3/1 \) island is excited. This is accompanied by a sudden drop in the central line-averaged density. The toroidal rotation profile shape remains peaked, implying no locking of a large \( m/n = 2/1 \) island with a \( m/n = 1/1 \) mode. The sawtooth oscillation as observed in the central electron temperature (figure 6.6d) also continues. The increase in central electron temperature can be attributed to the decrease in the central density. The electron temperature near the plasma edge, as measured with a single radial row of ECE-Imaging, decreases at certain radii near the estimated \( q = 3 \) surface (figure 6.6e). Only at \( t = 3.32s \) the \( m/n = 2/1 \) island is observed to become excited due to an external perturbation of the plasma. A strong decrease in the toroidal rotation and a flattening of the rotation profile are observed. This is accompanied by the disappearance of the sawtooth oscillation and a drop in the central electron temperature (figure 6.6d). Also another drop of the central line-averaged electron density is observed. Time averaged radial electron temperature profiles, measured during the three different phases of the discharge are shown in figure 6.7. Without an excited island, during the DED ramp-up phase, the temperature gradient at \( R = 2.16m \) near the \( q = 3 \) surface is nearly constant. The \( T_e \) time traces also remain constant when overlooking 40Hz oscillations due to heat pulses. After \( t = 2.5s \) electron temperature becomes flattened from \( R \approx 2.14m \) to \( R \approx 2.19m \). After \( t = 3.32s \), when the \( m/n = 2/1 \) island becomes excited, the flattening stays at \( R = 2.16m \) but the width decreases to about 4cm.
Figure 6.6: Discharge with a $m/n = 3/1$ island being excited before a $m/n = 2/1$ island; (a) DED coil current; (b) Toroidal rotation angular frequency near the plasma centre (solid blue $R = 1.85m$) and just inside the $q = 3$ surface (dashed green $R = 2.13m$); (c) Central line-averaged density, (d) Central electron temperature; (e) Electron temperature measured with a single radial row of ECE-Imaging at different radial position between $q = 2$ and the plasma edge; (f) Heating Power of MECRH (solid orange), Ohmic (dash dotted blue) and Neutral Beam (dashed red). The times at which an $m/n = 2/1$ island and an $m/n = 3/1$ island are excited are indicated by dashed black vertical lines.

6.5.2 Fourier analysis $m/n = 3/1$ island

A single radial row of ECE-Imaging channels that measures near the magnetic island O-point is evaluated at the first and at the second harmonic. The time
traces of these channels are also shown in figure 6.6e. Its amplitude and phase profiles are shown in figure 6.8. As in the case of the \( m/n = 2/1 \) island, the \( \Delta \phi \)-profiles show local maxima and the \( A \)-profiles show local minima at the position of the flattened temperature gradient due to the \( m/n = 3/1 \) island. The gradients inside the island are large compared the gradients outside the island and the case without island. Comparison of the first and second harmonic can tell something about non-diffusive features in the electron thermal transport. With pure diffusion one would expect that the gradients of the second harmonic for both \( \Delta \phi \) and \( \ln(A) \) to be a factor \( \sqrt{2} \) larger than those of the first harmonic. The present data show that in the island the gradient of the second harmonic \( \Delta \phi \) is more than twice the gradient of the first harmonic; in contrast, the gradient of the second harmonic \( \ln(A) \) is slightly smaller than the gradient of the first harmonic. This indicates the
Figure 6.8: Amplitude and phase delay of the first and second harmonic temperature perturbations from different plasma states. In the left panels the first harmonic phase delay (a) and the amplitude (c) are shown. In the right panels the second harmonic phase delay (b) and the amplitude (d) are shown. Blue circles are from the DED ramp-up phase without an excited island, green squares are from the phase with only an excited $m/n = 3/1$ island, red crosses are from the state with an excited $m/n = 3/1$ island in the presence of a large excited $m/n = 2/1$ island.

presence of a damping term. Damping is known to affect the propagation of the lower frequencies more than the higher frequencies and thus becomes relatively important at low modulation frequencies resulting in shallower gradients of $\Delta \phi$ and larger gradients of $\ln(A)$ at the lowest harmonics [11]. Damping could be caused by strong coupling of the electrons to the ions due to a higher collisionality near the plasma edge.

When the $m/n = 2/1$ island (not in the observation area) becomes destabilized, the $m/n = 3/1$ island stays but the local island width becomes smaller. Also the peaks of the amplitude and the phase delay become smaller. Figure 6.9 shows a 2D pictures of $\Delta \phi$ in the vicinity of the predicted $m/n = 3/1$ island O-point.
for three phases of the discharge: (a) during the DED ramp-up and before the \( m/n = 3/1 \) island is excited, (b) after an excited \( m/n = 3/1 \) island and (c) with the \( m/n = 3/1 \) island with the \( m/n = 2/1 \) island being excited. When the \( m/n = 2/1 \) island has not yet been excited, a pronounced local maximum due to the island O-point is observed at \( R = 2.16m \). After the \( m/n = 2/1 \) island is excited, the \( m/n = 3/1 \) island becomes smaller; the width of the peak reduces to less than 4\( cm \) at the bottom part of the observation area. Near the top, no clear local maximum is measured anymore; the phase delay profile is nearly flat. Either the local island width has become smaller than the radial resolution of an ECE-Imaging channel (\( \sim 1.5cm \)) or the heat pulse propagation is hardly influenced by the island due to the perpendicular transport dominating the parallel transport across the island.
6.5.3 Heat transport coefficients inside and outside the \( m/n = 3/1 \) island

The large gradients of the amplitude and the phase delay inside the island compared with outside the island show that also inside the \( m/n = 3/1 \) island \( \chi_e^{(hp)} \) is reduced compared with the ambient plasma. From the phase delay the heat pulse diffusivity inside the island can be determined, using the same method as for the \( m/n = 2/1 \) island, from the island width and the phase difference between the island centre and edges. This yields for the \( m/n = 3/1 \) island without the presence of the \( m/n = 2/1 \) island for the first (second) harmonic \( \chi_e^{(hp)} = 0.5(0.4)m^2/s \). For the smaller \( m/n = 3/1 \) island, when the \( m/n = 2/1 \) island has been excited, \( \chi_e^{(hp)} = 0.3m^2/s \) and \( \chi_e^{(hp)} = 0.4m^2/s \) are found for the first and second harmonic, respectively. Thus for islands of different sizes the heat pulse diffusion coefficient is reduced compared with the ambient plasma and they have nearly the same value.

Outside the island the gradients of the phase delay are too shallow for an accurate heat diffusivity determination. For the state without an island a crude estimation is made from the phase delay and the amplitude decay of the second harmonic between \( R = 2.10m \) and \( R = 2.16m \), using a slab model with cylindrical corrections. This yields a heat diffusivity of \( \chi_e^{(hp)} > 10m^2/s \). Power balance analysis yields \( \chi_e^{(pb)} = 3.0 \pm 0.5m^2/s \) around the \( q = 3 \) radius. Despite the higher heat pulse diffusion coefficients of the ambient plasma, the heat pulse diffusion coefficients inside the island are of a comparable low magnitude as those measured for the large \( m/n = 2/1 \) island with lower ambient transport coefficients.

6.6 Discussion

6.6.1 Heat diffusivity in island

The main point addressed in this paper is the question whether the heat diffusivity inside an island is considerably different from the bulk plasma. Several different situations have been analyzed: the heat diffusivity has been calculated using either the method of power balance or the heat pulse propagation. Values have been obtained for the island O-point, the island X-point, the bulk plasma or the reference case without an island. Results from [3] can be directly compared with the present results. The collection of data points are plotted in figure 6.10. If it is hypothesized that the heat diffusivity can be described in a simplified model by two
6.6 Discussion

Figure 6.10: Tentative model of the heat diffusivity in these experiments. Two branch model capturing the measurements of the electron heat diffusivities of a large $m/n = 2/1$ island. The green points are the measurements. The thick blue line represents a two-branch stiff profile model. The gradients of the red dotted line represent the power balance heat diffusivities. The gradients of the blue lines represent the heat pulse diffusivities. Please note that in case of the X-point, the gradient could not be exactly determined (X-point is just outside view of ECE-Imaging in figure 5), but is extrapolated.

Linear branches, then our data points can be captured by two straight lines. This two branches model is inspired by a 'stiff profile' model. For tokamak plasmas it is generally found that the electron heat diffusivity increases strongly with the electron heat flux for a temperature gradient larger than a critical gradient, leading to 'stiff' profiles [14]. This behaviour is represented by the blue lines in figure 6.10. The gradients of the blue line yield the heat pulse diffusivity $\chi_{e}^{(hp)} = -\partial q_{e}/n_{e}\partial \nabla T_{e}$. The gradients of the red lines yield the power balance heat diffusivity $\chi_{e}^{(pb)} = -q_{e}/n_{e}\nabla T_{e}$. The model shows that the heat pulse diffusivity is larger than the power balance diffusivity, which is commonly observed. Please note that we also include a data point for the X-point heat diffusivity in the plot.
It is just an extrapolation and has not been directly measured. It is based on the assumption that the heat flux from the plasma centre is channelled through an area which is smaller than for the case without an island, assuming that the island is effectively blocking the heat transport. From the Poincaré plot of figure 6.1, one could estimate that the heat is transported only through the stochastic region around the X-point, which has an area of about half of the total flux surface. This transport along the stochastic field lines can be an extra transport channel in addition to the electrostatic turbulence. The gradient of the X-point could not be accurately determined but is extrapolated (X-point is just outside view of ECE-Imaging, figure 6.5). When one compares the measurements of the \( m/n = 2/1 \) island case with the \( m/n = 3/1 \) island case, strikingly, inside the islands the transport coefficients are of comparable magnitude, irrespective of the much stronger heating in the \( m/n = 3/1 \) island case (1.2MW instead of 0.3MW of NBI). Due to the stronger heating in the plasma centre, the electron heat flux is estimated to be a factor of \( \sim 1.5 \) larger (in spite of the larger radial position of the \( m/n = 3/1 \) island). Hence it is very likely that the \( T_e \)-profile is driven more unstable in the \( m/n = 3/1 \) island case, i.e. is driven more above a critical gradient for some (e.g. TEM) instability. The experimental footprint of this is a larger ratio of \( \chi_e^{(hp)} \) over \( \chi_e^{(pb)} \) \[14\].

6.6.2 Critical island width

Another point that can be addressed is the critical island width for which the temperature is constant on the magnetic surfaces inside the island. It has been shown that for islands smaller than a certain width the temperature is no longer constant on the magnetic surfaces inside the island such that a temperature gradient at the island O-point can exist \[9, 19\]. The incomplete flattening at the island O-point is due to the competition of fast heat transport parallel to the field and slow heat transport perpendicular to the field. The temperature is a magnetic flux surface function inside the island when the heat transport time from one side of the island parallel to the field is smaller than the heat transport time perpendicular to the field through the island

\[
\tau_\parallel \approx \frac{L^2_\parallel}{\chi_\parallel} < \tau_\perp \approx \frac{w^2}{\chi_\perp} \tag{6.9}
\]

The transport time parallel to the field \( \tau_\parallel \) depends on the parallel heat conductivity \( \chi_\parallel \) and on the parallel connection length from one side of the island to the
other side of the island $L_{||}$. At the island O-point the parallel connection length is minimal $L_{||} = 2\pi R_0 q/q'nw$ [9], with $q$ the local safety factor. Close to the island boundary the parallel connection length becomes very large due to field line stagnation in the X-point region. There the temperature is predicted not to be constant on the magnetic surfaces. The same arguments can be used for perturbations of the temperature. Islands larger than a critical island width will yield non-monotonic radial phase and amplitude profiles for the heat pulse propagation across the island O-point. For islands smaller than the critical island width the profiles will be monotonic. An expression for the critical island width can be found in [9]:

$$w_c = \left( \frac{8R_0 q}{nq'} \right)^{1/2} \left( \frac{\chi_{\perp}}{\chi_{||}} \right)^{1/4}$$

(6.10)

A minimal critical island width can now be estimated. The parallel heat diffusivity is taken to be governed by conduction With an effective ion charge $Z_{eff} = 2$ and a Coulomb logarithm $\ln(\Lambda) = 15$ this yields for the $m/n = 2/1$ island $\chi_{||} \approx 2 \times 10^8 m^2/s$ ($T_e = 250eV$, $n_e = 2.2 \times 10^{19} m^{-3}$) and for the $m/n = 3/1$ island $\chi_{||} \approx 7 \times 10^7 m^2/s$ ($T_e = 170eV$, $n_e = 2.2 \times 10^{19} m^{-3}$). The measured heat pulse diffusivities of the islands of different sizes are of comparable magnitude. A similar observation was made in [10]. The experimentally determined heat pulse diffusivity is used for the perpendicular heat conductivity $\chi_{\perp} \approx 0.4 m^2/s$. This is done under the assumption that the measured heat pulse diffusivity is not much larger than the minimum heat diffusivity inside the island and that its magnitude is independent of the island size, down to the critical island width. For the $m/n = 2/1$ and $m/n = 3/1$ island the minimal critical widths $w_c = 1 - 2 cm$ are derived. This width is comparable to the radial resolution of the ECE-Imaging channels ($\sim 1.5 cm$). A more accurate estimate of the critical island width requires knowledge of transport in the island boundary of islands near the critical island size. There the perpendicular heat pulse transport coefficients may strongly vary since they are small inside the island and large outside the island.

### 6.6.3 Conclusion

The excitation of large $m/n = 2/1$ and $m/n = 3/1$ islands in the TEXTOR Tokamak by the DED is accompanied by a decrease in the energy confinement. The heat pulses are mostly diverted along the island boundary and propagate with a delay towards the island O-point. Inside the $m/n = 2/1$ and $m/n = 3/1$ island $\chi^{(hp)}_e$
is reduced, respectively, by 1 and 2 orders of magnitude compared with outside the island or compared with the plasma without an island. The values of the heat pulse diffusivity $\chi_\perp = 0.4 m^2/s$ inside the $m/n = 2/1$ island and the $m/n = 3/1$ island are comparable despite the different levels of $\chi_{e,(hp)}$ outside the island, which differ by an order of magnitude due to the different level of background heating. Nevertheless $\chi_{e,(hp)}$ inside the island is still an order of magnitude above the neoclassical level $\chi_{e,\text{neo}} \approx 0.01 m^2/s$. Some authors, e.g. [4, 15, 10], have claimed that the heat confinement is improved inside an island. Our results support this, but it is noted that in all cases no (large) temperature gradients were present and therefore these measurements all fall in the lower branch of figure 6.10. Under strong local ECRH heating inside the island, a strong temperature peaking inside the island has indeed been measured [3]. Power balance calculations however yielded $\chi_{e,(pb)} = 1 - 2m^2/s$, comparable to the ambient plasma. This can be understood by a stiff profile which governs the transport in the ambient plasma also being applicable inside the island. Low perpendicular heat diffusion coefficients inside an island are beneficial for the suppression of tearing modes by ECRH and electron cyclotron current drive (ECCD) [22]. Besides the non-inductive current driven inside the island by ECCD, an inductive current can be driven by the creation of a temperature peaking inside the island. The effect of heating is always present during ECCD. The efficiency by which a peaking can be created depends on the electron heat diffusivity inside the island, which in turn depends on the peaking.

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7

Determining the sawtooth mode structure with 2D ECE-Imaging

7.1 Introduction

The core of a tokamak plasma is commonly affected by a regular, periodic reorganization which is known as the sawtooth instability [24]. Initially in a sawtooth cycle, the core temperature increases slowly due to heating. This is followed by a rapid crash in which energy is ejected from the core towards the edge. In a fusion device, the occurrence of sawtooth oscillations has both desirable as well as undesirable effects [8]. The plasma can easily survive small sawteeth with small temperature drops and mixing radii. Sawteeth may also by preventing impurity accumulation and removing helium ash from the plasma centre. However, sawtooth crashes with large amplitudes and mixing radii can also cause the excitation of neoclassical tearing modes (NTM's) which degrade plasma pressure and thereby the fusion performance of the plasma [19]. To mitigate these effects accurate control of the instability leading up to the crash is needed.

In experiments a wide variety of observations of sawtooth phenomena have been made [7]. The sawtooth crash is associated with instabilities at the $q = 1$ radius. Before the sawtooth crash a pre-cursor mode with dominant mode numbers $m = 1, n = 1$ usually is observed to grow. The measured shape of the plasma deformation due to the precursor also varies from a growing 'cold' $m = 1, n = 1$ island to a 'cold' bubble that is drawn into the plasma centre. The type of sawtooth crash may depend on the particular plasma conditions. Measurements indicate
that plasma shaping effects such as triangularity can change the type of sawtooth instability [11]. Depending on the various contributions one or the other instability mode is expected to dominate. It may therefore be concluded that the sawtooth crash is not a physically unique phenomenon [28]. However, in general, two main classes of instability have been put forward to explain the sawtooth crash: the resistive kink mode or the quasi interchange instability.

For the control strategies it is desirable to know what the instability leading to the crash is. While lowering the local shear \( s_1 \) at the \( q = 1 \) surface stabilizes for the resistive kink mode, it destabilizes the quasi-interchange instability. Therefore, electron cyclotron current drive and heating (ECCD/ECRH) is envisaged as an ideal control tool, since it can modify the local shear and the temperature gradient independently. However, for a reliable control strategy, both the type of sawtooth instability as well as the actual position of the \( q = 1 \)-radius is a prerequisite. A diagnostic to provide this information (in real time) is not readily available. Accurate \( q \)-profile measurements could in principle be used to determine what kind of instability mode is responsible for the sawtooth crash, but these are difficult to perform. Alternatively, one could try to determine the mode-structure of the instability mode. A common technique to achieve this is based on tomography of soft X-ray (SXR), which is cumbersome due to the tomographic inversion of the limited number of line integrated measurements to determine the mode structure.

The mode structure due to the sawtooth instability can also be recognized from the electron temperature profile as determined from electron cyclotron emission (ECE). However, as most conventional ECE-systems are 1D, it is hard to resolve the fast evolving 2D mode structures during the crash. Only under the constraint that the mode is rotating, and slowly evolving during a rotation, a 2D reconstruction can be made to identify the mode structure [13].

An ECE-Imaging system with a unique 2D observation area has been developed and used on the TEXTOR tokamak [14]. With the 2D system, a direct 2D mapping of the temperature evolution during the sawtooth crash has been reported [16, 15]. These measurements found that the temperature fluctuations of the sawtooth precursor oscillations resemble a circular hot spot and a growing cold island and not a crescent shaped hot core as expected for the quasi-interchange mode. During the crash, a fast poloidally localized heat flow across the sawtooth inversion radius is observed with the sawtooth inversion radius being defined as the radius where the temperature change is minimal before and after the crash. A drawback of these measurements was, that due to the limited extension of the
observation area, only a small region of the total inversion radius was covered, which makes it hard to get a complete view of the dynamics going on or to identify the shape of the large mode structures. Recently the observation area of the 2D ECE-Imaging system has been increased [Zhang 2008] which allows for identification of the large scale mode dynamics across the entire \( q = 1 \) radius during the crash.

The goal in this chapter is to develop a method based on these 2D ECE measurements to identify the mode structure during the sawtooth crash on TEXTOR and determine whether it is due to a quasi-interchange mode or a resistive internal kink mode. Furthermore, having determined the mode structure, a comparison between the experimental observations and the expectations following the corresponding theoretical model can be made. The insights provided by the measurements can aid in developing better control methods of the sawtooth instability. Such a control method is beyond the scope of this chapter.

The plan of the chapter is as follows; in section 7.2 the basic sawtooth models found in literature are described and the specific features of the modes which allow them to be identified with ECE measurements. In section 7.3 the 2D-ECE-Imaging system as it is implemented on the TEXTOR tokamak is introduced. As the ECE-Imaging is normally not calibrated, an effort is made in section 7.3 to find suited methods for data normalization to interpret the measurements. In section 7.4 experimental results from a scan of the low-field side to the high-field side of the \( q = 1 \) radius are presented and analyzed to determine the mode structure. Finally, in section 7.5 conclusions are made on the ability to determine mode structure of the sawtooth crash with 2D ECE-Imaging.

7.2 Sawtooth models

Aiming mainly at the identification of the type of sawtooth instability, this section serves as a summary of the main features of the two classes of sawtooth models: the resistive internal kink and the quasi interchange mode. It is not a complete treatment of the theory behind these models, but it focuses on the features, that can be experimentally detected to determine the type of instability. Also it is investigated if one or the other model provides an adequate description of the experimental situation.

The main features of the sawtooth instability are clear, the sawtooth period can be subdivided in three distinct phases: a precursor phase, crash phase and post-
crash phase. Often (but not always!) a phase is observed in which an \( m = 1, n = 1 \) mode grows, referred to as the precursor mode. Sometimes the \( m = 1, n = 1 \) mode is locked to the laboratory frame instead of rotating toroidally. In that case the 1-D diagnostics does not show temporal variation which can be interpreted as the absence of the mode. This precursor phase is followed by the crash phase in which part of the central heat is expelled from the core causing a flattening of the central temperature profile and an increase of the temperature across the sawtooth inversion radius. The inversion radius is often associated to be close to the \( q = 1 \) radius. Note however that they do not necessarily coincide because the evolution of the \( q \)-profile during the crash is not clear.

From other experiments it has become clear that also part of the current-density \[25, 20\] and toroidal momentum are expelled \[21\]. Finally, in the post-crash phase the heat pulse propagates outward and the core normally is rearranged into a poloidal symmetric pressure profile. These phases are ingredients to the sawtooth models, but the models differ in the way they describe the development of the precursor mode develops and how it induces the final crash.

### 7.2.1 Sawtooth models bases on resistive internal kink instability

One of the earliest and certainly most influential theoretical model of the sawtooth crash has been Kadomtsev’s complete reconnection model \[10\]. In this model the sawtooth crash is caused by the growth of an \( m = 1, n = 1 \) resistive internal kink mode which causes the crash. The evolution of this mode has been numerically calculated in \[22\] and is schematically shown in figure 7.1. Due to the good thermal confinement, the central \( T_e \) increases and hence the current density peaking. When the central safety factor becomes smaller than unity, reconnection occurs at the \( q = 1 \) radius.

The mode is called resistive kink because the \( m = 1, n = 1 \) displacement of the core is accompanied by the reconnection and the formation of a magnetic island with the same helicity. During the growth, helical magnetic flux from both sides of the \( q = 1 \) radius is reconnected at the \( m = 1, n = 1 \) island X-point which causes \( q > 1 \) inside of the island. With the reconnection the radial position of the X-point moves outward \[10\]. The reconnection process continues until the central safety factor \( q_0 > 1 \) and the newly formed island has completely replaced the former core. The mixing radius is the maximum radius of the X-point where full reconnection is reached. Poloidal symmetry is regained after full reconnection.

Focusing on the question what this model would predict for the temperature
evolution, it is found that hot plasma from the core and cooler plasma from outside the core are mixed in the newly formed island. When prior to the sawtooth crash the $T_e$-profiles (as well as electron density profile ($n_e$)) are peaked, then after full reconnection the profiles are flattened within the mixing radius. During the reconnection finite parallel transport processes near the island separatrix will influence the temperature distribution [3]. For large islands the temperature is expected to be a flux-surface function inside of the island. For small islands widths the temperature will not be a flux-surface function and no complete flattening will be observed. Compared to experimental measurements, Kadomtsev’s full reconnection model however has a number of shortcomings related to the onset of the mode, the crash, the typical crash time and the current distribution in the plasma core after the crash:

- The resistive internal kink becomes destabilized as soon as the central safety factor $q_0 < 1$. Measurements however show a sudden onset of sawtooth precursor modes and the sawtooth crash indicating that the mode is initially stabilized while $q_0 < 1$.

- The typical reconnection time calculated by Kadomtsev is $\tau_K = \sqrt{\tau_A \tau_R}$

\[\begin{align*}
\text{(a)} & \quad \text{Poloidal symmetric configuration with } q > 1. \\
\text{(b)} & \quad \text{Central safety factor has become } q < 1 \text{ prior to the growth of the instability.} \\
\text{(c)} & \quad \text{A 'cold' island has formed and replaced } \sim 50\% \text{ the former core.} \\
\end{align*}\]
where \( \tau_A = r_1/(B_t \sqrt{\mu_0 \rho}) \) is the Alfvén time and \( \tau_R = \mu_0 r_1^2/\eta \) is the characteristic resistive timescale near the \( q = 1 \) radius. With \( \tau_A \ll \tau_R \) the reconnection time scale predicted by Kadomtsev’s full reconnection model is much shorter than the resistive timescale. In the small-scale tokamak plasma, \( \tau_K \) is comparable to experimental values. However in the larger and hotter tokamak plasmas such as JET have larger \( \tau_R \) yielding \( \tau_K \) much larger than the observed collapse times.

- A third inconsistency between the Kadomtsev’s model and experiments is found in the value for the central safety factor. Whereas Kadomtsev’s model predicts that the central \( q_0 > 1 \) after the crash, detailed investigations at TEXTOR [20] showed that \( q_0 < 1 \) after the crash, showing that full reconnection does not occur during the crash. The observation of the so-called snake instability [5, 25] in sawtoothing plasmas is consistent with the observation that the central safety factor remains below unity during the sawtooth crash. The snake appears as a persistent density peaking with an \( m = 1, n = 1 \) island structure that is caused by the evaporation and ionization of injected frozen hydrogen pellets on the \( q = 1 \) surface. The snake can survive sawtooth crashes while its finite radius to the centre is reduced with typically with 20 to 30%. This implies that there is a \( q = 1 \) radius with finite shear which is consistent with \( q_0 < 1 \) after the sawtooth crash.

To cope with the shortcomings of Kadomtsev’s model, modifications have been proposed by others:

- In [18] different stabilizing and destabilizing mechanisms such as fast particles effects are incorporated in determining the instability of the resistive internal kink mode. The trigger condition for the growth onset of the \( m = 1, n = 1 \) is given when the shear at the \( q = 1 \) radius exceeds a critical shear \( (s_1 > s_{1, \text{crit}}) \). Many experiments report on the triggering of the sawtooth with local current drive near the \( q = 1 \) radius. These confirm that the sawtooth crashes are triggered faster when the current drive increases the shear around \( q = 1 \) radius and vice versa, e.g. [1, 12, 29].

- A secondary instability mechanism that is proposed by Gimblett and Hastie [4] could be responsible for the rapid outflow of the heat during the crash and the central safety factor being \( q_0 < 1 \) after the crash. In [4] it is assumed
that the relaxation mechanism is due to resistive 'g'-modes [6] which are
driven by the unfavorable field line curvature. It is proposed that these
modes become destabilized at a stage during the growth of the resistive
\( m = 1, n = 1 \) mode. The resistive \( m = 1, n = 1 \) kink mode causes a large
pressure gradient \((dp/dr)\) just outside of the island separatrix. When the
pressure gradient close to the separatrix exceeds a critical value, nearby
tearing modes with \( q = m/n \approx 1 \) (e.g. \( m/n \approx 9/10 \)) can rapidly grow
causing the collapse process.

### 7.2.2 Sawtooth model based on Quasi-interchange instability

An alternative model proposed to explain the sawtooth crash is Wesson’s quasi-
interchange model [26]. The quasi-interchange instability is driven by the pressure
gradient in contrast to a resistive kink which is also driven by a current gradient.
The plasma core is instable to quasi-interchange motion if the central safety factor
remains close to unity \( (q_0 \approx 1) \) and additionally there is low shear \( s_1 \) within the
\( q = 1 \) radius. The low magnetic shear causes that the magnetic surfaces with
\( q = 1 \) require little energy to be deformed. At low pressure (low poloidal beta, \( \beta_p \))
the mode can then become unstable. The stability of this mode can be somewhat
modified by the shape of the flux-surfaces which influences the average shear. An
unstable deformation of the core magnetic surfaces as determined in [26] is shown
in figure 7.2. The plasma core with \( q \approx 1 \) is deformed by the flow of cold plasma

![Figure 7.2: Sketch of Hot core deformation into a crescent shaped core due to a quasi-interchange instability, adapted from [26]. The arrows denote the direction of the plasma flow leading to the deformation. In the color-scaling red denotes high temperature and green lower temperature.](image)

from outside of the core towards the plasma centre and the accompanying back
flow of hot plasma out of the plasma centre.

To trigger the sawtooth instability, only a small change of the \( q \)-profile until it reaches \( q = 1 \), is sufficient [26]. Once destabilized the growth rate can increase on a time-scale faster than the resistive time-scale since no reconnection is necessary. The original model [26] is based on ideal MHD motion and no reconnection or relaxation mechanism is defined during the crash.

A quasi-interchange deformation of the hot core yields a flux surface pattern at the \( q = 1 \) radius that is distinctively different from the rigid \( m = 1, n = 1 \) shift due to the resistive internal kink. The difference is due to the shear near \( q = 1 \). The finite shear of the resistive internal kink results in flux surfaces which have a near-circular cross-section embedded in a clear \( m = 1, n = 1 \) island structure. Prior to the crash the quasi-interchange motion causes a strong crescent shaped deformation of the hot core with respect to the unperturbed equilibrium.

Experimental evidence for this interpretation of the sawtooth crash exists as well. The original papers reporting a crescent shape prior to a sawtooth crash on JET as measured by SXR [2] have been disputed by TdeV results. This interpretation appears to be an artifact of the tomographic reconstruction due to the limited number of SXR cameras [9]. Other experimental support for this model has been reported from different devices: A crescent shaped hot core has been observed in \( T_e \) image reconstructions of sawtooth oscillations on the Tore Supra tokamak and on TEXTOR [23]. On DIII-D [11] the effect of a quasi-interchange mode on the electron temperature profiles is measured.

In summary, the first obvious experimental goal will be to discriminate the shape of the precursor mode being either a hot rotating core embedded in a cold island or a crescent hot shape surrounding a cold bubble.

### 7.3 Experimental setup

The reported measurements are performed on TEXTOR, a circular limiter tokamak with the dimensions in major/minor radius of \( R/a = 1.75m/0.47m \), on which a 2D ECE imaging system is operational. Other features of relevance for the sawtooth investigations are the flexibility to control the sawtooth amplitude and frequency by additional beam heating. Bi-directional beams are used, which allows the power and the torque to be varied such that the precursor frequency can be varied. For the experiments discussed here, the total beam power is chosen to be 1.3\( MW \) in the co-current direction and 0.8\( MW \) counter-current direction which
yields large amplitude sawtooth-crashes that are preceded by a rotating precursor mode of constant frequency.

The principle diagnostic is the 2D ECE system [31] and a detailed treatment of it is already given in chapter 4. Its main features are summarized here. The ECE-imaging system measures the 2D ECE intensity of an area in the poloidal plane with $8 \times 16$ radiometer channels. The vertical size of the observation area can be adapted by the optical system. In the "wide zoom" configuration, as used here, the ECE-imaging system has a vertical coverage of $\Delta Z = 30\,cm$ and a radial coverage of $\Delta R \approx 10\,cm$. For every channel the radial resolution is $\Delta R \approx 1\,cm$ and the vertical resolution is $\Delta Z = 1\,cm$. The measured second harmonic X-mode ECE intensity from the plasma centre is proportional to $T_e$ because the plasma there is optically thick ($\tau^{2X} \gg 1$). The ECE is measured at a sample frequency of $50\,kHz$ which yields a minimum ECE noise level of $\Delta T_e/T_e \geq 1\%$ [31].

Since the coverage of the ECE system is limited, care should be taken to image the relevant area. For the sawtooth investigations, it is obvious to position the system around the inversion radius. To a reasonable approximation, this can be estimated as $r_{inv}/a = 1/q_a$. For the cases under investigation it is aimed to have $r_{inv} \approx 10\,cm$. This allows sufficient resolution and full coverage of mode dynamics in the vertical direction but in the radial direction there is however no full coverage. To achieve full coverage, the radial position is scanned by adjusting the magnetic field in successive discharges (see figure 7.3). The main plasma parameters chosen are a magnetic field of $B_t = 2.15 - 2.4 T$, and plasma current of $I_p = 310\,kA$, resulting in an edge safety factor $q_a \approx 4.2 - 4.7$. The line averaged plasma density is taken $\langle n_e \rangle \approx 3 \times 10^{19}\,m^{-3}$.

### 7.3.1 Noise reduction with Singular Value Decomposition

A point of concern is the intrinsic noise on the ECE measurements, which could disturb the interpretation of the spatial features during the sawtooth. The relative noise level of the measured ECEI signals $\Delta T_e/T_e$ decreases with the time resolution. To reduce the influence of high-frequency correlated and uncorrelated noise, the measurements are filtered with a low-pass frequency filter at $50\,kHz$ which allows several harmonics of the sawtooth precursor frequency $\sim 5\,kHz$ to be identified. Further noise reduction is achieved by application of singular value decomposition (SVD) of the multi-channel data. This method is illustrated in chapter 4.

The measurements presented here are approximated by a reconstruction in
Figure 7.3: Calculated positions of the ECE-Imaging system with respect to the estimated sawtooth inversion radius (dashed blue circle). The high field side (HFS), centre and low field side radial measurement positions observed in the different discharges are indicated by the black squares. For the centre position the individual channel positions are indicated by blue crosses.

which only the $\sim 15$ largest singular value pairs are retained, meaning that the dynamics contained in the picture is a linear combination of $\sim 15$ different pictures (spatial eigenvectors). In this way the intrinsic noise is in principle reduced by a factor $\sqrt{15/128} = 0.34$.

7.3.2 Data normalization

The ECE intensity signals from the ECE-Imaging system are not standardly calibrated (due to system settings that are difficult to reproduce after they have been changed). For an accurate calibration however a detailed and accurate $T_e$-profile measurement is needed which is not available for the discharges of interest. Instead the shapes of normalized profiles $(T_e(R, Z, t) - T_{e,norm}(R, Z))/T_{e,norm}(R, Z)$ are determined. Such a normalized profile can be readily determined because the
detected ECE intensity of each ECE-imaging channel \((i,j)\) is proportional to the electron temperature

\[
\frac{T_e(Z_i, R_j, t)}{T_{e,norm}(Z_i, R_j)} - 1 = \frac{I_{ECE}(Z_i, R_j, t)}{I_{ECE,norm}(Z_i, R_j)} - 1 \tag{7.1}
\]

Knowledge of the shape \(T_e(Z, R, t)\) then depends on knowledge of the general shape of \(T_{e,norm}(Z, R)\). Profiles \(T_{e,norm}(Z, R)\) of which the general profile shape is known should be used for the normalization.

The question is still open of what are suited profiles \(T_{e,norm}(R, Z)\). Two different normalizations are explored here for their applicability to distinguish between structures due to resistive internal kink mode or to the quasi-interchange mode:

- In [15] the mode structure is determined from the normalization of the time averaged profiles over a sawtooth period: \(\langle T_e(Z, R, t) \rangle_{T_{sawt}}\). This profile is close to the average of the peaked profiles before the sawtooth crash and the flat profile after the sawtooth crash. This method thus highlights temperature variations with respect to the average. The shape of the normalized profile is not directly proportional to the shape of the \(T_e\)-profile which is a disadvantage for a direct interpretation.

- The other normalization that is explored is \(T_{e,norm}(Z, R) = \min[T_e](Z, R)\) with \(\min[T_e](Z, R)\) being defined as the minimum \(T_e\) attained at a position \((Z, R)\) during a sawtooth period. An advantage of this normalization is that it yields profiles that are proportional to \(T_e(Z, R)\) inside of the inversion radius because there the minimum \(T_e\)-profile, which is attained right after the sawtooth crash, is flat. Outside the sawtooth inversion radius the minimum \(T_e\) is usually attained just before the sawtooth crash. There the \(\min[T_e](Z, R)\)-profile has lower values than inside of the inversion radius. The effect of the lower \(\min[T_e]\) outside of the inversion radius on the normalization is an enhancement of the \(T_e\)-profile features outside of the inversion radius such as the heat redistribution after the sawtooth crash. It is noted that in chapter 5 has been found to be of the order \(-R\nabla T_e/T_e \approx 10\) for a variety of plasma conditions and that this knowledge could be used to acquire a normalized profile that also outside of the proportional to \(T_e(Z, R, t)\).

To study the sensitivity of these normalizations with respect to the ability to distinguish a quasi-interchange perturbation from a resistive internal kink perturbation, artificial \(T_e(R, Z)\)-profiles of these modes are generated and normalized.
Figure 7.4: The 2D normalizations of a simulated $m = 1, n = 1$ resistive internal kink mode are shown in the upper row. A quasi-interchange mode displacement is shown in the second row, a centrally flattened profile after the crash in the third row and radial cross-sections of the profiles are shown in the fourth row. The type of normalization is different for every column of the figure. The first column shows the simulated $T_e$-profiles. The second column shows the normalized profiles $T_e/\min[T_e]$. The different methods allow distinguishing between a circular shaped displaced hot core of the resistive internal kink and the deformed hot core of the quasi interchange.

The resulting normalizations are depicted in figure 7.4. Here the 2D normalizations of a simulated $m = 1, n = 1$ resistive internal kink mode are shown in the upper row, a quasi-interchange mode displacement in the second row, a centrally flattened profile after the crash in the third row and radial cross-sections of the
profiles are shown in the fourth row. Each column of the figure represents a different normalization procedures, and the first column shows the absolute temperature contours.

The assumptions made in the simulation and pictures are the following: The resistive kink has a flat $T_e$-profile inside the island. The island separatrix position is indicated by dashed lines. The outer separatrix coincides with the inversion radius (note that in Kadomtsev's model this is nearly correct for small islands widths). The $T_e$-profile within the inner separatrix (hot core) is quadratic. For the quasi-interchange mode a crescent shaped hot core is simulated. For the flat $T_e$-profile after the sawtooth-crash a flat $T_e$-profile with equal $T_e$ as in the island is taken. Outside the plasma core the $T_e$ is increased after the crash.

The second column shows the $T_e/\min[T_e]$-profile. The profile shape of the hot core is preserved under this normalization. The crescent shaped core of the quasi-interchange and the circular shape of the resistive internal kink are clearly distinguishable with this normalization. The island edge however is not distinguishable.

The third column shows $(T_e/\langle T_e \rangle)_{T_{sawt}}$. For the resistive internal kink the normalization yields nearly circular shaped surfaces in the plasma centre. The circular shaped surfaces are positioned radially further outward when compared to the $T_e$-profile. For the island structure the normalization yields a decrease in the relative temperature profile. The normalization appears to enhance the quasi-interchange mode deformation. This normalization thus also allows distinguishing the circular shaped hot core of the resistive kink from the strongly deformed hot core of the quasi-interchange. The advantage of $T_e/\min[T_e]$ however is that its shape is directly proportional to $T_e$ within the inversion radius.

In the bottom row the 1D radial cross-sections of all the profiles are plotted in case only a 1D radial ECE system would be available. Clearly the difference in all these profiles is only marginal (also taken into account that noise is not simulated) so that it will be a tough task to distinguish between the two modes. This justifies the use of the 2D system.

Inspection of figure 7.4 justifies determining the curvature of the normalized profiles in the plasma centre to distinguish both modes. For the quasi-interchange mode the radius of curvature of the contours is directed away from the hot core and for the resistive internal kink it is directed towards the hot core. When the direction of the radius of curvature can be measured, a clear distinction can be made between both the modes.

In summary, the normalization procedures of the ECE data should allow a
direct identification of the simulated 2D shapes in the poloidal plane within the plasma centre by visual inspection. To obtain the exact shape of the profile within the inversion radius the normalization $T_e/\min[T_e]$ is best suited, provided the $T_e$-profile is indeed flat after the sawtooth crash. Although the normalization $T_e/\langle T_e \rangle_{\text{sawt}}$ does not directly represent the $T_e$ contours, it also does allow the identification of a circular shaped hot core of the resistive kink and the crescent shaped hot core due to the quasi-interchange mode. The mode structures even appear enhanced by the normalization. Using either one of the normalizations, the quasi-interchange deformation can be clearly distinguished from the resistive internal kink by determining the curvature of the normalized profiles. In this chapter the normalization $T_e/\min[T_e]$ is applied to the experimental data.

7.4 Experimental results

7.4.1 General observations

To introduce the sawtooth behavior measured in TEXTOR, first the gross features of a representative sawtooth oscillation are discussed. In figure 7.5a the time traces of ECEI data from three different minor radii are shown together with the temporal spectral evolution figure 7.5b. From these data, already some general observations can be made regarding the sawtooth behaviour and the application of normalizations. Sawtooth crashes occur at $t = 1.3417$s and at $t = 1.3811$s as observed in figure 7.5a. Inside the sawtooth inversion radius, the minimum $T_e$ is indeed achieved just after the sawtooth crash. Outside the inversion radius the minimum $T_e$ is achieved at half the the sawtooth period and remains nearly constant until the crash.

In the evolution of $T_e$ a coherent oscillation is observed (at $t = 1.362$s). The central $T_e$ saturates with the appearance of the coherent oscillation. The evolution of amplitude and frequency of the $T_e$-oscillation due to the mode is shown with a spectrogram in figure 7.5b. The amplitude of the $T_e$-oscillation which can be inferred from the color in the spectrogram, is a measure for the size of the mode. The mode oscillation persists up to the sawtooth-crash (half a sawtooth period $\sim 20ms$) and it attains a rotation frequency of $f_{\text{prec}} \approx 6kHz$ and amplitude. The crash appears suddenly with regard to the nearly constant mode amplitude and saturated central $T_e$. After the crash, no clear coherent oscillations are resolved in the electron temperature which indicates that no rotating mode is left and implies that the $T_e$-profile has regained poloidal symmetry. The central $T_e$-profile is
Figure 7.5: Panel (a) shows three temperature time traces of ECE-channels measuring at minor radii close to the plasma centre (blue), just inside the inversion radius (green) and outside of the inversion radius (red). Sawtooth crashes occur at $t = 1.3417$ s and at $t = 1.3811$ s. After the crash $T_e$ increases in the plasma centre and decreases outside the centre. The first sign of a rotating precursor-mode is observed at $t = 1.36$ s which coincides with the saturation of the central $T_e$ increase. Panel (b) shows the time resolved spectral content of the green time trace from panel (a). A mode with a frequency $f_{\text{prec}} \approx 6\text{kHz}$ is observed. The amplitude of the mode saturates. After the crashes no coherent modes oscillations are observed in the time resolved spectra implying poloidal symmetry.

The main focus is on the identification of the mode structure leading to the sawtooth crash. From the 2D ECE-Imaging measurements the normalized $T_e/\min[T_e]$-profiles are determined, which should result in an undistorted image of the structure inside the inversion radius, provided the $T_e$-profile after the crash is flat. The entire area of the sawtooth inversion radius is covered by radially scanning the ECE-Imaging observation area in reproducible discharges as described in figure 7.3. The measured sequence of profiles $T_e/\min[T_e]$ of the plasma centre during a precursor oscillation at half the sawtooth period and a precursor oscillation
preceding the sawtooth crash is shown in figure 7.6a and b respectively. Figure

![Figure 7.6](image)

**Figure 7.6:** Contour plot of the normalized electron temperature $\Delta T_e / \min[T_e]$ prior and during the sawtooth crash (TEXTOR #108341). (a) The upper series of frames are during an early phase of the precursor oscillation around $t = 1.3647s$ with the detection frame focused at the centre of the plasma. (b) The lower series of frames are around the crash phase $t = 1.3811s$. The time difference between the frames is $\Delta t = 20\mu s$. The contour separation $\Delta T_e / \min[T_e] = 0.02$. The color scaling goes from green, through yellow up to red. The small solid blue circle denotes the estimated plasma centre around which the hot core rotates. The dashed blue line denotes the estimated inversion radius. The small black crosses denote channels that are evaluated for the analysis. Collective heat flow during the sawtooth crash is observed in the 10th frame of series (b).

7.6a shows a hot (red) core performing $\sim 1.5$ pre-cursor oscillations around the plasma centre which is denoted with the small blue circle. In the centre, the precursor oscillations in $T_e$ are minimal. The excursion of the core with respect to the centre is $1 - 2cm$ ($\sim 2$ channels radially and $\sim 1$ channel vertically). During the precursor oscillations the shape of the hot core does not clearly deviate from circular when considering the deformations due to the noise in the measurements. No clear distinction between a quasi-interchange and a resistive kink is made in this phase.

In figure 7.6b it is focused on the $T_e$-distribution in the plasma centre. The hot
7.4 Experimental results

core makes an increasingly larger excursion from the plasma centre up to the crash. In the $10^{th}$ frame the plasma centre is at the edge of the hot core. The contours through the plasma centre have a radius of curvature that is directed towards the hot core. Opposite the hot core and inside the inversion radius the normalized $T_e$-profile is flattened. In the $11^{th}$ frame the $T_e$-profile has flattened over the inversion radius area that is within the ECE-I observation area. The central observations of the $T_e$-distribution inside of the inversion radius of a displaced circular shaped hot core and a crescent shaped flattened $T_e$-profile are consistent with a growing resistive kink mode and inconsistent with those of a quasi-interchange.

Next, it is focused on the heat-redistribution across the inversion radius to see if this phase of the crash can also be described by the growth of a resistive internal kink mode. The heat from the core collectively passes the sawtooth inversion radius in the $10^{th}$ frame of figure 7.6b. Outside the inversion radius a large gradient in the poloidal direction is observed but the extent of this $T_e$-structure is not resolved. In the subsequent frames the heat becomes more spread in the poloidal direction. A better view of the $T_e$ redistribution and the mode structure during the crash is obtained by focusing of the ECE-Imaging system to the region around the inversion radius, either on the low field side (figure 7.7) or high field side (figure 7.8).

Because the plasma density for low field side measurements is somewhat higher than the density for the central and the high field side measurements ($\langle n_e \rangle = 4.0 \times 10^{19}m^{-3}$ compared to $3.0 \times 10^{19}m^{-3}$), the sawtooth precursor frequency is lower for this discharge. It is assumed that the higher density has no influence on this type of the sawtooth instability. The hot core is observed to perform one precursor oscillation during the frames of figure 7.7a and the first five frames of figure 7.7b. From the $4^{th}$ until $8^{th}$ frame of figure 7.7b a localized temperature increase across the inversion radius is observed. In the $7^{th}$ and $8^{th}$ frame of figure 7.7b a the heat is observed to collectively flow across the inversion radius. The poloidal extent of the temperature increase on the inversion radius is $\sim 8cm$ wide ($\sim 4$ vertical channel spaces) corresponding to $13\%$ of the poloidal circumference of the inversion radius. In the $8^{th}$ frame most of the heat has passed the inversion radius. Just within the inversion radius there is still a small part of the hot core with a radius of curvature of its $T_e$-contours still directed toward the centre of the hot core.

Looking at the high field side in figure 7.7, the core performs one precursor oscillation from the $1^{st}$ to the $10^{th}$ frame. Around the $10^{th}$ frame the heat from the core collectively passed the inversion radius. The poloidal extension of the hot
Figure 7.7: Contour plot of distribution of the normalized electron temperature $\Delta T_e / \min[T_e]$ prior and during the sawtooth crash at the low field side of the inversion radius (TEXTOR #108342, $t = 1.3755s$). The time difference between the frames is $\Delta t = 20\mu s$. The sawtooth crash occurs around the 7th frame of figure 7.8b. The increase in $T_e$ outside of the inversion radius is poloidally and radially localized initially and spreads in the poloidal direction. In the 8th frame a small part of the hot core is still inside of the inversion radius. The radius of curvature of the $T_e$-contours is still directed towards the hot core which is consistent with a resistive kink displacement and inconsistent with a quasi-interchange deformation.

The observations inside of the inversion radius for both low field side and high field side are consistent with the central ECE-I observations from figure 7.6. It can therefore be assumed that the type of sawtooth in both discharges are the same. The heat redistribution outside of the inversion radius due to the sawtooth however appears to be different at the low field side compared to the high field side.
7.4 Experimental results

Figure 7.8: Contour plots of the normalized electron temperature $T_e/\min[T_e]$ prior and during the sawtooth crash at the high field side of the inversion radius (TEXTOR #108337, $t = 1.7233s$). The time difference between the frames is $\Delta t = 20\mu s$. The sawtooth crash occurs around the 10th picture. In the 12th picture when most of the heat has passed the inversion radius, the radius of curvature of the $T_e$-contours is directed towards the hot core which is consistent with a resistive kink displacement and inconsistent with a quasi-interchange deformation.

comparison between two different magnetic flux-surface shapes due to resistive kink-modes is made in figure 7.9 to explain this. In figure 7.9a the resistive kink as predicted by Kadomtsev’s theory has grown to large amplitude with the outer island separatrix extending beyond the inversion radius. Also the hot core extends to beyond the inversion radius. The resulting normalized $T_e$-profile is shown in figure 7.9b; it shows a poloidal symmetric increase of $T_e$ outside of the inversion radius at the position of the island and a localized $T_e$-increase at the position of the hot core. The normalization of panel (b) is similar to what is observed for the normalized $T_e$-profiles at the high-field side shown in figure 7.8.

In figure 7.9c a localized $T_e$-bulge at the island X-point position is produced. The resulting normalization of this profile, shown in figure 7.9d, is similar to what is observed in the measured profiles from the low-field side shown in 7.7b. The observed asymmetry may be caused by finite pressure effects of the hot core that cause a more pronounced deformation at the low magnetic field side than at the high magnetic field side [17].

It is concluded that the 2D ECE-I $T_e/\min[T_e]$-profile measurements within the sawtooth inversion radius resemble $T_e$-profiles expected from the evolution of a resistive internal kink mode. The observed $T_e/\min[T_e]$-profile evolution outside of the inversion radius may be explained to be due resistive kink modified with finite pressure effects causing a deformation of the magnetic surfaces that is stronger at the high field side than at the low field side.
Figure 7.9: Normalizations of modeled resistive kink modes. The top figures show 2D modeled profiles, the bottom panels show the corresponding cross-sections at \( Z = 0 \). In panel (a) a resistive internal kink mode with an island whose outer separatrix extends beyond the inversion radius. The corresponding \( T_e / \min[T_e] \)-profile is shown in panel (b). A \( T_e \)-profile with a hot core that causes a \( T_e \)-bulge beyond the outer separatrix is shown in panel (c) and the corresponding \( T_e / \min[T_e] \)-profile is shown in panel (d).

7.5 Summary

The main aim of this chapter was to develop a method based on 2D ECE measurements to identify the mode structure during the sawtooth crash. The mode structure determines the control strategy for the sawtooth instabilities. Depending on the mode, current drive or heating is preferred. The first hurdle to tackle was the identification of a suitable normalization procedure. Relying on an individual calibration of the numerous ECE channels to construct a \( T_e \)-profile, is not regarded as feasible. In this work it was shown that even with uncalibrated data the essential features of the different possible modes at play during the sawtooth crash could be unambiguously distinguished. To accomplish this, a general assumption on the temperature profile after the sawtooth crash (flat profile inside the inversion radius) has to be used. Accepting this, two methods have been analyzed: one based on a normalization with respect to the minimum temperature during the sawtooth phase and one based on the average temperature.
The next step was to investigate if these methods could be used to distinguish the mode structure of two distinctly different classes of sawtooth instabilities: the quasi-interchange mode and the resistive internal kink mode. It turned out that the clearest criterion to discriminate between the two instabilities is the direction of the radius of curvature of the $T_e$-profile-contours at the plasma centre: for a quasi-interchange the radius of curvature at the plasma centre is directed away from the displaced hot core and for the resistive internal kink mode it is directed towards the hot core.

As an application of the method, an example of sawteeth at TEXTOR were presented. For the measured $T_e/\min[T_e]$-profile structure the radius of curvature which suggests a resistive internal kink mode. The measured evolution of the $T_e/\min[T_e]$-profile outside of the sawtooth inversion radius could be explained in terms of a resistive kink mode whereby the high pressure of the displaced hot core causes a bulging of the further outward magnetic surfaces. The pressure effect is larger at the low field side than at the high field side of a magnetic surface. This is due to the the stabilizing effect of the magnetic shear which is smaller at the low magnetic field side compared to the high magnetic field side.

Several other methods have been proposed and tested in the past to determine this mode structure. They relied either on 1D ECE data or SXR tomography. The 1D ECE method used a mapping of the time to the poloidal angle (for instance:[30, 13, 23]) to obtain a quasi 2D picture. The main shortcoming is that this method can only be applied for relatively slow sawtooth crashes. Interpretation of the results is therefore not unambiguous, as illustrated in the example of [23], where a hot crescent shape was concluded for TEXTOR, whereas our results undoubtedly identified the true 2D character, being a resistive internal kink for this specific case. For the SXR based technique it has been shown that the conclusions may by falsified by a limited number of channels [9]. Only with a sufficient large set of detectors and a cumbersome tomographic reconstruction, similar results as in the more straightforward method proposed here can be obtained.

The main conclusion of the work reported here is that a method is proposed using a 2D ECE system, that allows the identification of the mode structure during a sawtooth crash (which is essentially 2D). The real time availability of this information might have benefits in the sawtooth control strategies, since it determines the power deposition location and injection angle of the microwave system envisaged for this purpose in fusion experiments like ITER.
Chapter 7. Determining the sawtooth mode structure with 2D ECE-Imaging

References


Chapter 3: Mhd stability, operational limits and disruptions, Nuclear Fusion 47 (2007), no. 6, S128.


Chapter 7. Determining the sawtooth mode structure with 2D ECE-Imaging


The next generation of fusion reactor (ITER) is projected to produce fusion energy. Achieving the burn criteria is within reach but controlling the burn adds another dimension to the operation. Manipulating the transport properties of the plasma is the envisaged method to do so. Magneto hydrodynamic (MHD) modes can be used for this. They can appear spontaneously, degrading the confinement, requiring a suppression of these modes. Alternatively, by local manipulation of the current profile or pressure profile, these modes can be triggered to reduce the confinement if needed. A control scheme based on this (i.e. suppression or triggering of the modes by local profile manipulation) still has to be developed and demonstrated. Knowledge on the changing transport properties of the plasma in the vicinity or in the presence of these modes is a prerequisite to come to a control of the fusion performance. This is the main theme of this thesis. To address this issue, three different questions were posed in the beginning of this thesis:

- Can the profile shape be locally modified, to modify the transport properties?
- How does the transport behave in the vicinity of an MHD mode, an island?
- Can the instability mode responsible for the sawtooth crash be resolved from the two-dimensional transport during a sawtooth-crash?

First a summary is given here to answer these questions, after which the relevance and the outlook of this work are discussed in a broader perspective.
Chapter 8. Evaluation and outlook

8.1 Profile-consistency

The main question posed in chapter 5 was whether localized deposition of heating could cause a local modification of the current density and pressure profiles in a plasma without large scale MHD modes. For centrally heated plasmas from different heat sources, the current and the pressure profiles derived from the measured electron temperature and density profiles were found to be similarly shaped for a large part of the plasma radius. These profile shapes of the electron pressure, current, temperature, density are well described by the model of profile consistency that only depends on a single variable. For off-axis \( mm \)-wave deposition it has been found that the \( T_e \)-profile shape and the closely related \( j \)-profile shape could be locally modified. While outside of the heat-deposition region the profile shape was found to be unaffected, inside of it the profile shape is flattened. Such local current profile modifications are not expected from profile-consistency, but can be well described by a less restrictive critical temperature gradient model. In conclusion, these studies confirmed that despite the profile stiffness, it is well possible to locally modify the temperature and current density profiles by local deposition of \( mm \)-waves.

8.2 Heat-pulse propagation in and around magnetic islands

The control of magnetic island growth depends on the ability to drive a current inside the magnetic island. The efficiency by which localized \( mm \)-wave deposition can drive a current depends on the electron heat-transport properties inside the magnetic island. The electron heat transport properties of magnetic islands are however not well known. The main question posed in chapter 6 is whether the electron heat transport properties inside a magnetic island give rise to electron temperature gradient dependent transport, such as is observed for plasmas without magnetic islands. Statically locked \( m/n = 2/1 \) islands and \( m/n = 3/1 \) islands have been created with the DED in 3/1 configuration. Localized ECRH has been deposited outside the islands to create electron heat-pulses. The propagation of the electron heat-pulses in and around the island has been measured with 2D ECE-Imaging. The heat-pulse diffusivity has been determined from the propagation speed in the flat
temperature region of the island. Fast propagation around the island across the island X-point has been observed, whilst the propagation towards the island O-point was relatively slow. The corresponding heat-pulse diffusivities inside the 2/1 and 3/1 islands are \( \chi_e^{(hp)} \approx 0.2 - 0.3 m^2/s \), while outside the island diffusivities were found to be an order of magnitude larger, with \( \chi_e^{(hp)} > 2 m^2/s \).

From previous measurements of peaked temperature profiles due to ECRH deposition inside the island, a power balance heat-diffusivity of \( \chi_e^{(pb)} > 1 m^2/s \) has been found. Since it is generally observed that the heat-pulse diffusivity is larger than the power balance heat diffusivity, this implies that the heat diffusivity inside of the island increases with the electron temperature gradient. These results confirm that the heat-transport inside magnetic islands is temperature gradient dependent, similar to observations without magnetic islands.

In small to medium-sized limiter tokamaks such as TEXTOR and T-10 the effect of heating dominates the direct current drive in stabilizing the magnetic island growth [2, 5]. For large tokamak reactors like ITER, the contribution of current drive efficiency is expected to be significantly larger than the effect of heating [3].

In conclusion, the low heat-diffusivity in a magnetic island with a flattened temperature profile allows to efficiently create a temperature peaking inside a magnetic island with local mm-wave deposition. The efficiency will however decrease as the temperature profile becomes more peaked inside of the island.

### 8.3 Sawtooth oscillations

In chapter 7 it is assessed whether a sawtooth instability mode can be resolved with a 2D ECE-I system, so that such a system may eventually be used for sawtooth control. Two instability modes are generally identified, a quasi-interchange mode and a resistive internal kink mode. For both modes the parameters (pressure and shear) at the \( q = 1 \) radius determine the stability of the respective modes differently. Control of these modes requires a local modification of these parameters which can be achieved with ECRH/ECCD. Efficient control requires the identification of the instability modes. It has been assessed whether these mode could be identified using an uncalibrated 2D ECE-I system.

Temperature normalization procedures have been developed and tested on simulated 2D \( T_e \)-structures due to a resistive internal kink mode and a quasi-interchange mode. It has been found that with the tested normalizations the \( T_e \)-profiles due to a resistive kink or due to a quasi-interchange mode could be
well distinguished from each other, based on the curvature of the normalized $T_e$-surfaces inside of the inversion radius. Two different normalizations have been applied to TEXOR data. The mode preceding the sawtooth crash causes a displaced circular shaped hot core. During the crash of the central $T_e$, the hot core still consists of circular shaped magnetic surfaces inside the inversion radius. A poloidal localized $T_e$ increase across the inversion radius and it has been observed at the low field side of the inversion radius and it has been explained as a deformation of the magnetic surfaces. These $T_e$-structures resemble a resistive internal kink mode. No temperature structures resembling a quasi-interchange deformation of the hot core were found in the investigated discharges.

The ability to identify 2D structures in the $T_e$-distribution gives confidence that 2D ECE-I can be used for MHD-mode identification.

8.4 Outlook

The relevance of the work can be classified in three categories:

a relevance on island suppression

b relevance sawtooth control

c relevance of the 2D ECE diagnostic

Ad a) The finding that the heat-conductivity inside a magnetic island is temperature-gradient dependent means that the efficiency by which $mm$-waves can create a temperature peaking inside a magnetic island decreases with the applied heating power. The deposition of $mm$-waves can also directly drive a current inside a magnetic island. For low temperature islands such as found in TEXTOR, the effect of non-inductive current-drive is negligible compared to inductive current-drive (i.e. due to the temperature rise) [3]. For high temperature tokamak plasmas such as ITER, the direct (non-inductive) current-drive is expected to be much larger than the inductive current-drive but this latter term (due to heating) will, however, be non-negligible [2, 5]. Nevertheless, increasing the applied $mm$-wave power will decrease the efficiency of inductive current-drive, but not the efficiency of the non-inductive component [3]. Thus at high heating power the heating term will become negligible. The requirement on the necessary heating power for island suppression will thus not change appreciable due heating inside the islands.
The method developed to measure with the combination of heat pulse propagation and 2D ECE might have further relevant applications. It could be used to determine the position and the size of stationary islands or locate transport barriers. The traditional method to determine the island position is by looking to the temperature perturbations (on 1D ECE) as induced by the island rotation or magnetic coil measurements, related to changing currents (due to the island rotation). Clearly these methods will not work in case of stationary islands, but the 2D ECE – heat pulse technique will do. The same technique can be applied in case of transport barriers.

Ad b) For control of the sawtooth-instability the mode structures needs to be determined in real-time. Although the underlying instability mode structure could be resolved from the measurements in this work, the measurements were still hampered by the limited observation area of the measurement system. An increase of the ECE-Imaging observation will help resolve the mode structure in real-time. An increase of the observation area will also be of use to resolve the a-symmetries between low and high-field side of a kink-like sawtooth crashes. ECE-Imaging systems with increased observation areas that can simultaneously measure at the low and high-field side of the sawtooth inversion radius have recently been installed at DIII-D [4] and K-Star [6], inspired by the results on the sawtooth measurements on TEXTOR.

Ad c) Also on the diagnostic side, the results of this work might have further reaching implications. The demonstration that 2D information is important and the 2D ECE is a powerful tool for these investigations, opens the possibility to apply it for various other subjects. For a fusion reactor two specific problems should be addressed by this:

- The identification of the instabilities triggering the edge localize modes (ELMs) [1]. This is one of the hot issues in fusion research, since ELMs will reduce the wall lifetime (and thus the reactor availability) considerably. Mitigation of ELMs is therefore required.

- The localization of Alfvén modes. These modes can occur if fast particles are present in the reactor, which are produced in a burning plasma in the form of alpha particles. The interaction of these modes with the alpha particles, might reduce their confinement and thus leads to a reduced heating of the plasma.
In conclusion: 2D effects are important to control the burning plasma, 2D ECE imaging is a powerful tool for measuring these local 2D effects and should therefore be further exploited. These measurements can assist in the control strategy of different instability modes and thus the plasma performance.

References


A very promising source for future energy is nuclear fusion. Achieving energy production from fusion reactions on earth however is very challenging. The fusion fuel, a plasma of hydrogen isotopes, needs to be brought to high temperatures and densities, to produce a net energy gain.

Fusion machines of the Tokamak concept have already approached the required fusion conditions by means of magnetic confinement of the plasma fusion fuel. The fusion energy production is degraded by magnetic instabilities such as island formation and sawtooth crashes.

For the control of the fusion process, it is important to control the growth of these magnetic instabilities. The growth of magnetic instabilities depends on the plasma current profile. The current profile can be modified by the deposition of high power millimeter waves. The millimeter wave deposition can drive currents directly and/or inductively. Inductive current drive is due to a local modification of the electron temperature profile. The efficiency of inductive current drive depends on the electron heat transport properties of the plasma. On future tokamak machines such as ITER, the direct current drive is expected to be larger than the inductive current drive.

The effect of inductive current drive on the stabilization of magnetic instabilities depends on the electron transport properties in the vicinity of these instabilities. The transport properties in the vicinity of instability modes such as magnetic islands are however not well known.

With the aim of characterizing these properties, transport experiments were carried out on the TEXTOR Tokamak. TEXTOR is well equipped with a set of experimental tools that allows a deeper investigation of the transport properties near the magnetic instabilities.

First it has been investigated whether the electron temperature-profile and
closely related electron current-profile could be modified with heating in plasmas without magnetic islands. The profiles of the electron temperature and density have been measured with high-resolution Thomson scattering. Central heating of the plasmas from different sources yielded similarly shaped electron temperature profiles and current-profiles outside the plasma center. The profiles of the plasma center were however found to be strongly affected by the sawtooth instability. A local modification of the electron temperature profile and the inductive current profile was created by local electron heating outside the plasma centre.

Secondly it has been investigated how magnetic islands locally modify the heat transport properties of the plasma. Besides the non-inductively driven current, the inductively driven current is of importance for control of the island growth. The electron transport properties in the magnetic islands vicinity are however poorly known. For this reason, perturbative electron heat transport experiments in the vicinity of magnetic islands have been performed. Magnetic islands with a flat electron temperature profile yielded low electron heat conductivity coefficients. In previous experiments with peaked electron temperature profiles, high electron heat conductivity coefficients have been found. An increase in the electron heat conductivity coefficient inside the island causes a decrease of the efficiency by which an inductive current is driven in the island.

Thirdly it has been investigated whether different sawtooth instabilities could be identified, based on the two-dimensional heat transport they cause. Sawtooth crashes can be caused by different instabilities. The two main instabilities held responsible for the sawtooth crash are the 'resistive internal kink mode' and the 'quasi-interchange mode'. The control of the different instabilities requires different applications of heating and current drive. For control it is important to distinguish the different instability modes. In general the sawtooth crash happens on a time-scale shorter than one millisecond, which makes it difficult to distinguish the different instability modes. These modes can be distinguished from their characteristic two-dimensional temperature redistribution pattern during the sawtooth crash. It has been shown to be difficult to distinguish these instability modes with conventional one-dimensional temperature measurements or line-integrated measurements. Therefore the application of direct two-dimensional ECE-Imaging measurements has been evaluated. Two different temperature normalizations for ECE-Imaging data have been evaluated for their ability to distinguishing the two different mode structures. It has been shown that both normalization methods allow a clear distinction of the temperature structures due to the quasi-interchange mode and the resistive internal kink mode. The application of these normaliza-
tions to two-dimensional ECE-Imaging data from sawtooth-crashes in TEXTOR clearly showed a two-dimensional temperature evolution that is well described by a resistive internal kink mode.
Een veel belovende toekomstige energiebron is nucleaire fusie. Het bereiken van energieproduktie met fusiereacties op aarde is echter een grote uitdaging. De fusiebrandstof, een plasma van waterstof isotopen, moet een hoge temperatuur en een hoge dichtheid bereiken, om netto energie te produceren.

Fusie machines van het type Tokamak hebben, door gebruik te maken van magnetische opsluiting van de fusiebrandstof, de vereiste fusie condities al benaderd. Door de vorming van magnetische instabiliteiten in het plasma, zoals magnetische eilanden en zaagtanden, wordt de fusie energieproductie negatief beïnvloed.

Om het fusie proces te controleren, is het van belang de groei van magnetische instabiliteiten te kunnen controleren. De groei van magnetische instabiliteiten is afhankelijk van het stroomprofiel in het plasma. Met de depositie van hoog vermogen millimetergolven, kan men het stroomprofiel beïnvloeden.

Met millimetergolven kan direct en/of inductief stroom gedreven worden. Inductieve stroomdrijving vindt plaats door een lokale modificatie van het elektronen temperatuurprofiel. De efficiëntie waarmee inductief stroom wordt gedreven, is afhankelijk van de elektronen transport eigenschappen van het plasma. Verwacht wordt dat op toekomstige tokamak machines zoals ITER, de directe stroomdrijving groter is dan de inductieve stroomdrijving.

Het effect van inductieve stroomdrijving op magnetische instabiliteiten is afhankelijk van de transport eigenschappen in de nabijheid van de instabiliteiten. De transport eigenschappen nabij instabiliteiten zoals magnetische eilanden zijn echter niet goed gekarakteriseerd.

Met het doel om deze eigenschappen te karakteriseren zijn transport experimenten uitgevoerd op de TEXTOR tokamak. TEXTOR is uitgerust met een set experimentele gereedschappen die een nader onderzoek van de transport eigenschappen van het eiland mogelijk maken.
Als eerste is onderzocht of het elektronen temperatuurprofiel en het gere- lateerde stroomprofiel met verhitting gemodificeerd kunnen worden in plasmas zonder magnetische eilanden. De elektronen temperatuurprofielen en de elektronen dichtheidprofielen zijn gemeten met hoge resolutie Thomson verstrooiïng. Toepassing van verschillende centraal gedeponeerde verhittingsbronnen zorgden buiten het plasmacentrum voor gelijk gevormde elektronen temperatuurprofielen en stroomprofielen. In het plasmacentrum werden de profielen echter sterk beïnvloed door de zaagtand instabiliteit. Een locale modificatie van het elektronen temperatuurprofiel en het inductief stroomprofiel kon gecreëerd worden door locale electronenverhitting buiten het plasmacentrum.

Als tweede is onderzocht of magnetische eilanden de lokale elektronen warmte transport eigenschappen beïnvloeden. Voor het controleren van de eilandgroei is naast niet-inductieve stroomdrijving ook inductieve stroomdrijving van belang. De transport eigenschappen nabij een magnetisch eiland zijn echter slecht gekaracteriseerd. Om deze reden zijn experimenten gedaan met verstoring van het elektronen warmte transport nabij magnetische eilanden. Voor magnetische eilanden met een vlak temperatuursprofiel zijn lage elektronen warmte geleidings coëfficiënten gevonden. In eerdere experimenten met gepiekte temperatuursprofielen in eilanden zijn hoge warmtegeleidingscoëfficiënten gevonden. De elektronen warmtegeleiding coëfficiënt in het eiland neemt dus toe met extra verhitting in het eiland. Een toename in de elektronen warmtegeleiding coëfficiënt in het eiland zorgt voor een afname in de efficiëntie waarmee een inductieve stroom in een magnetisch eiland wordt gedreven.

Ten derde is onderzocht of verschillende zaagtandinstabiliteiten geïdentificeerd kunnen worden op basis van het twee-dimensionaal warmtetransport dat zij veroorzaken. Zaagtandcrashes kunnen veroorzaakt worden door verschillende instabiliteiten. De twee belangrijkste instabiliteiten die verantwoordelijk worden gehouden voor de zaagtandcrash zijn de ‘resistive internal kink mode’ en de ‘quasi-interchange mode’. De controle van de verschillende instabiliteiten vereist verschillende toepassingen van hitte depositie en stroomdrijving. Voor controle is het van belang om de verschillende instabiliteiten te kunnen onderscheiden. De zaagtandcrash vindt in het algemeen plaats op een tijdsschaal welke korter is dan één milliseconde, wat het moeilijk maakt om het type instabiliteit te onderscheiden. Deze instabiliteiten kunnen onderscheiden worden door hun karakteristieke twee-dimensionale temperatuurs herverdeling gedurende de zaagtandcrash. Het is moeilijk gebleken om deze instabiliteiten op te lossen met conventionele één-dimensionale temperatuursmetingen of lijngeïntegreerde metingen. Daarom is
de toepassing van directe twee-dimensionale metingen met ECE-Imaging geëvalueerd. Twee verschillende temperatuur normalisaties voor ECE-Imaging data zijn geëvalueerd op hun vermogen om de twee instabiliteiten te onderscheiden. Het is aangetoond dat beiden duidelijk onderscheid kunnen maken tussen de "quasi-interchange mode" en de "resistive internal kink mode". Toepassing op ECE-Imaging data van zaag tand-crashes in TEXTOR liet duidelijk een twee-dimensionale temperatuursevolutie zien, overeenkomstig met de evolutie van een 'resistive internal kink mode'. 
Curriculum Vitae

I was born on December 15, 1978 in the Dutch town Holten. I attended the secondary school ‘De Waerdenborch’ in Holten and obtained my VWO diploma in 1997. I studied applied physics at the University of Twente, where I graduated in 2004. As part of my study I did a three month traineeship at the University of Surrey on determining the pulse shape of ultra-short laserpulses. In 2005 I started working as a Ph.D. student at the FOM Institute for Plasma Physics Rijnhuizen. The experimental work presented in this thesis was carried out on the TEXTOR tokamak at the Forschungszentrum Jülich.
I like to thank a number of people who have supported me during my research. Especially, I like to thank my promotors Niek Lopes Cardozo, Chris Schüller and Roger Jaspers for their supervision. I thank Tony Donné for introducing me to the world of fusion.

Successful experiments could not have been performed without the experienced TEXTOR team. The ECE-Imaging team provided me with several upgrades of the system during the past years. I am grateful to Mark van de Pol and Calvin Domier, who repaired the ECE-Imaging system several times for me.

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