REVIEW

The reversed field pinch

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The reversed field pinch

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Abstract

This paper reviews the research on the reversed field pinch (RFP) in the last three decades. Substantial experimental and theoretical progress and transformational changes have been achieved since the last review (Bodin 1990 Nucl. Fusion 30 1717–37). The experiments have been performed in devices with different sizes and capabilities. The largest are RFX-mod in Padova (Italy) and MST in Madison (USA). The experimental community includes also EXTRAP-T2R in Sweden, RELAX in Japan and KTX in China. Impressive improvements in the performance are the result of exploration of two lines: the high current operation (up to 2 MA) with the spontaneous occurrence of helical equilibria with good magnetic flux surfaces and the active control of the current profile. A crucial ingredient for the advancements obtained in the experiments has been the development of state-of-art active feedback control systems allowing the control of MHD instabilities in presence of a thin shell. The balance between achievements and still open issues leads us to the conclusion that the RFP can be a valuable and diverse contributor in the quest for fusion electricity.

Keywords: magnetic confinement, reversed field pinch, MHD

(Some figures may appear in colour only in the online journal)

1. Introduction

This paper reviews progress in understanding and improving the performance of the reversed field pinch (RFP) toroidal magnetic configuration as a fusion energy concept, covering several decades of research since the last reviews [1, 2]. The development of the RFP has progressed in parallel with other magnetic configurations, and in many respects, it is the third most developed toroidal configuration next to the tokamak and stellarator, benefiting from well-diagnosed experiments and comprehensive theory and modelling. Significant advances have accrued over the past several decades, and the perspective on the RFP is now very different. This review summarizes these advances with selected emphasis on those that have been most impactful in guiding future research directions on the RFP.

The advances in fusion research over the last few decades have been tremendous, and the grand challenge to achieve a burning plasma in ITER dominates the world’s magnetic fusion effort. Carbon-free energy sources are essential to limit global warming and avert climate change, while at the same time it is necessary to increase the abundance of affordable
energy for the benefit of all people in all nations. Nuclear fusion can and should play a central role in achieving a future based on clean, sustainable energy. While the tokamak configuration is most developed and provides the basis for ITER, fusion energy is simply too important to be narrowly constrained to one possible approach. Research on multiple approaches not only increases the possibility for fusion energy to become a reality but also exposes the underlying science more completely and stimulates innovation. The models governing fusion plasmas are largely agnostic to magnetic configuration. Experiments, theory, and modelling that cover a broad range of configurations allow the development of robust, predictive fusion science. In the end, any specific approach to fusion must have a tractable path forward, and at this point in time there are no known showstoppers exposed for the RFP. Some challenges that might have seemed insurmountable 30–40 years ago now have solutions, while others are much better understood and must be resolved with new experiments and advanced modelling. Only few key routes remain mostly unexplored for the RFP, in particular control of the plasma boundary at the material interface and power exhaust.

It has long been appreciated that the RFP offers a high beta, low field approach to fusion with all of the potential benefits that entails. However, the RFP’s unique advantage is the opportunity to achieve an ohmically ignited and inductively sustained toroidal fusion plasma, including the possibility for steady-state operation using ac magnetic helicity injection. The RFP is magnetized almost entirely by plasma current, which increases ohmic heating and reduces the demands on magnets. The poloidal current that generates toroidal magnetic field is mostly in the plasma, not in a winding. The elimination of auxiliary heating and current drive is game changing for toroidal plasma confinement. Induction of current is simple, reliable and efficient, accomplished with transformers that are located outside of the blanket. Plasma-facing antennas are absent, and perforations in the neutron shield are minimized. The applied toroidal field in an RFP is small and can even be set to zero, although the toroidal field magnet is valuable as a poloidal current transformer for optimizing inductive control and sustainment. The TITAN reactor study investigated the advantages of a compact, high beta, inductively sustained RFP plasma core, yielding one of the lowest cost-of-electricity projections to date. This study introduced novel design elements, including single-piece maintenance and an integrated blanket-magnet for the toroidal field. New scenarios and requirements have emerged in the intervening years, and updated reactor studies are starting again to affirm the RFP’s potential as a low-cost, reliable approach to fusion energy.

As a toroidal plasma, the RFP shares all of the fundamental requirements for fusion related to equilibrium, stability, energy and particle confinement, impurity control, and plasma–wall interactions (PWI) (but sans auxiliary heating). This review describes our current knowledge on these fundamental areas of fusion science for the RFP.

The first main change for the RFP in the last thirty years is that poor confinement from stochastic transport is no longer a barrier. Multiple tearing instabilities can appear (multiple helicity regime, MH), causing the magnetic field to become stochastic, the historic challenge for energy confinement in RFP plasmas. Two approaches show that stochastic transport can be circumvented, one through self-organization towards a single helicity (SH) regime that has a narrow mode spectrum, and one through control of the inductive electric field and the plasma current profile, which directly targets the free energy for instability. Energy confinement as good as for tokamak plasmas of comparable size, magnetic field strength, and heating power has now been demonstrated in RFP plasmas. The scaling of confinement has been better established from work on the largest RFP experiments, RFX-mod and MST; microturbulence from drift waves takes over when stochastic transport is reduced, likely to be the ultimate limit for thermal particle and energy confinement as it is for tokamak and stellarator plasmas. The scaling of energy confinement for larger, hotter, higher current plasmas remains an open question that can only be resolved with new experiments. The magnetic field strength in present-day RFP experiments is still five-times smaller than required for a reactor, bearing in mind $B \sim I_p/\alpha$, where $I_p$ is the toroidal plasma current, and $\alpha$ is the minor radius.

The RFP plasma also suffers instability to multiple ideal kink modes as a consequence of its low safety factor equilibrium, $q < 1$, even in the limit of vanishing thermal pressure. It is therefore essential that the plasma be surrounded by a conducting wall to prevent growth of global modes on an Alfvénic timescale. The finite conductivity of the stabilizing wall translates the instability to resistive wall modes (RWM). Multiple RWMs can be simultaneously unstable, which might have seemed an insurmountable hurdle. A second game changing development for the RFP is active mode control using saddle coils surrounding the wall and plasma, which controls all unstable modes and maintains a precise plasma boundary for pulse durations limited only by power supplies. Moreover, the control encompasses simultaneous magnetic field error correction as well. This pioneering development of active mode control for RFP plasmas has helped improve the prospects for steady-state advanced tokamak plasmas that suffer analogous high-beta kink modes, and it also creates connections to 3D physics shared by stellarator plasmas.

Advanced nonlinear, visco-resistive, 3D MHD models and simulations have been critical to understanding RFP plasmas. The development of major MHD codes now used extensively in magnetic fusion research has roots in the goal to understand the RFP’s magnetic self-organization. The knowledge gained from this work has inspired approaches to improve confinement in the RFP. Current profile control was invented in MHD modelling before being tested experimentally. Similarly, the bifurcation towards the single-helicity regime was initially discovered in modelling of high viscosity RFP plasmas and then quasi-helical states emerged in experiments. The latter were better reproduced in simulations by including a helical boundary condition.

Positive trends of temperature and confinement with plasma current have been observed in present experiments, and the upgrade of the RFX-mod device, presently underway, is
expected to improve the confidence on such scalings, giving the basis to conceive a future reactor based on the RFP concept.

Advances in understanding the RFP’s operational limits for density and beta have also been made. Interestingly, RFP plasmas obey the same empirical density limit as for tokamak plasmas, but the density is large since the plasma current density is large. The ideal beta limit for the RFP is also very large (order 40%), and the limiting behaviour is soft, most likely governed by non-ideal processes. Beta and density-driven fast disruptions are uncommon in RFP plasmas.

Investigation of energetic ions is a new area of research for the RFP in the last several years. Energetic ion confinement is found to be classical, better than the thermal particle confinement, even in the MH regime with a stochastic magnetic field. The energetic ion beta in the core exceeds the projected alpha particle $\beta$ for a reactor plasma when energetic-particle-driven activity is observed to onset.

An important possibility for the RFP is steady-state current sustainment using inductive current drive. The oscillating field current drive (OFCD) concept was proposed in conjunction with discussions of the conservation of magnetic helicity in relaxed-state plasmas. In OFCD, ac loop voltages inject cycle-average dc magnetic helicity to sustain the plasma current without accumulation of magnetizing flux. Fractional current drive by OFCD of at least 10% has been demonstrated, and nonlinear MHD modelling shows that the underlying dynamics are similar to conventional steady induction. Demonstrating full OFCD sustainment with small ac ripple is one of the key opportunities for a larger, higher current experiment. The development of high temperature superconducting magnet technology could be strongly enabling for OFCD, symbiotic with recent efforts to establish a compact, high-field tokamak demonstration reactor.

While this review is focussed on the fusion context for the RFP, the magnetic self-organization aspects of the RFP that are interesting from a basic science point of view with connections to astrophysics are implicit. Tearing instability relates to magnetic reconnection, and the nonlinear saturation process that regulates tearing relates to the magnetic dynamo process. Non-collisional ion heating and energization are particularly strong in RFP plasmas, likely associated with a turbulent cascade, and make excellent connections to similar processes in solar and astrophysical settings. These topics are embedded within the fusion context, with many of the key references provided in the bibliography.

1.1. Chapter contents

Chapter 2 gives an insight into the basic experimental properties of the RFP. The axisymmetric magnetic equilibrium model for the RFP is recalled. The dynamics of the configuration in the two different regimes of MH and quasi single helicity (QSH), characterised by different spectra of tearing modes, are described, and the helical equilibria in QSH states are introduced as well. The helical states with two magnetic axes are QSH double axis (QSH-DAx), while the helical state with a single helical axis is QSH-SHAx.

Chapter 3 summarizes the current understanding of the RFP self-organisation. Taylor’s theory of field reversal is recalled. The need of a dynamo process for the RFP sustainment is discussed, together with experimental measurements. The transition from MH to QSH states as the result of a bifurcation process is introduced. A second bifurcation is then described, from a QSH state characterised by two magnetic axes, the unperturbed axi-symmetric one and the axis related to the island of the dominant mode (QSH-DAx state), to a magnetic equilibrium featuring the helical axis only as the result of the expulsion of the magnetic separatrix (QSH-SHAx). The effect of a helical boundary in the simulation of helical states, allowing a better reproduction of the experiments, is discussed. The progress in the understanding of the helical equilibria allows a new vision of the dynamo process (electrostatic dynamo), common for multiple and QSH states.

The complex magnetic dynamics have a crucial impact on the confinement and transport properties in the RFP, which are addressed in chapter 4. In the MH state the overlap of the magnetic islands due to the core resonant dynamo modes causes a broad region of stochastic magnetic field, where the temperature profile is observed to be flat. The experimental results are compared to the Rochester and Rosenbluth (RR) theory of transport in a stochastic field [3]. In QSH states the temperature profile mirrors the magnetic topology: a small hot island is observed when the magnetic topology is characterized by the presence of two magnetic axes in QSH-DAx states. A wider region with high temperature, involving most of the plasma core and encompassed by strong gradients, is instead observed in QSH-SHAx states, which are characterized by improved transport and confinement properties. Studies about the impact of microturbulence in regimes of reduced stochastic transport, performed by gyrokinetic simulations, are also reported. A large number of experiments and simulations have been dedicated to edge transport and turbulence, and to fast ion and electron confinement as well. This is also reviewed in chapter 4.

The approach to improved confinement in the RFP based on current profile control is described in chapter 5. MHD modelling demonstrated that tearing fluctuations responsible for magnetic field stochasticization could be reduced or even eliminated with a suitable modification of the current profile. Several approaches to inductive current profile control have been studied in experiments and simulations. Temporal variations in the toroidal and/or poloidal magnetic field are used to control the inductive electric field, called pulsed poloidal (or parallel) current drive (PPCD). The dynamo electric field during PPCD is observed to become very small, so that the resistive term in the Ohm’s law is balanced by the applied electric field. This coincides with a large reduction of magnetic fluctuations and a ten-fold improvement in energy and thermal particle confinement time, with substantial enhancement of the electron and ion temperatures.

Chapter 6 deals with the limits to plasma pressure. RFPs share with tokamaks and stellarators a high density limit, traditionally described by the Greenwald empirical limit (or by the Sudor limit in the case of stellarators). This chapter describes
the experimental and modelling efforts to explain such a limit, thus contributing to the understanding of a still open issue for all magnetic configurations. The $\beta$ limit is different in the RFP with respect to the tokamak: the ideal $\beta$ limit is much higher in the RFP. In MST values of total $\beta$ as high as 26% have been measured, consistent with the Connor–Taylor scaling [4] and exceeding the Mercier criterion governing the interchange stability [5]. High beta values are also associated to neoclassical bootstrap current effects.

During the last 15 years RFP devices have become very attractive laboratories to study state-of-the-art active feedback control of MHD stability. To these studies is dedicated chapter 7. Theoretical studies on the stabilization of RWMs are described together with the experimental demonstration of their control. The focus is on EXTRAP-T2R and RFX-mod experiments, which are devices equipped with sophisticated active control systems. The chapter also deals with tearing modes, which in the RFP are ruled by nonlinear MHD (in contrast to RWM). The constructive interference due to phase-locking of the $m = 1$ and $m = 0$ tearing modes produces a toroidally localized helical deformation of the plasma column (slinky or locked mode, LM). Several control algorithms have been developed and tested to mitigate the local deformation of the magnetic field, mainly in RFX-mod, ultimately leading to a very effective reduction and slow rotation of LM.

As in other magnetic configurations, plasma-wall interactions (PWI) play a crucial role in determining the plasma behaviour and performance, in particular through the control of plasma density and impurity contamination. In the RFP PWI is related to the magnetic topology through the three-dimensional edge magnetic fields. This is discussed in chapter 8.

Chapter 9 summarizes the experimental and modelling studies made on a quasi-steady-state method of inductive current sustainment, called OFCD, based on the operation of oscillatory toroidal and poloidal loop voltages, resulting in the generation of a dc current via the nonlinear relaxation processes. Fractional current drive of at least 10% has been demonstrated, and nonlinear MHD modelling reveals the dynamics of relaxation. The scaling for the concomitant ac ripple in the equilibrium magnetic field is reviewed, which can be used to project the experimental requirements for full sustainment by OFCD.

Chapter 10 summarizes system studies for an RFP based reactor and neutron source, and concluding remarks are given in chapter 11.

1.2. The RFP devices

Since the 1990s, several RFP devices have been in operation, listed in table 1. The largest one is RFX-mod, designed to reach 2 MA maximum current, (formerly RFX, [6]) at Consorzio RFX in Padova, with major radius $R = 2$ m and minor radius $a = 0.459$ m. It is equipped with a thin copper shell with $\tau_w = 100$ ms ($\tau_w = \mu_0 a \rho \delta$, where $a$ is the minor radius, $\delta$ is the wall thickness and $\rho$ is the wall conductivity) and a network of actively controlled saddle coils for thin shell operations [7]. The MST device, located at the University of Wisconsin Madison, USA [8] ($R = 1.5$ m, $a = 0.50$ m) can operate up to 0.6 MA plasma current. The MST’s 5 cm thick aluminium shell serves as both the vacuum chamber as well as the toroidal and poloidal equilibrium field magnet. The shell has a minor radius of 0.52 m with graphite plasma limiters. The inductive current profile control technique was originally developed on MST. The facility includes a 1 MW, 25 keV, 20 ms neutral beam injector to study energetic ion confinement and stability.

EXTRAP-T2R [9], in operation at the Royal Institute of Technology, Stockholm, Sweden ($R = 1.24$ m, $a = 0.18$ m, maximum plasma current = 0.3 MA), is equipped with a copper shell with shell time 12 ms and focuses its experimental activity on the development of advanced tools for the feedback control of plasma MHD instabilities.

The TPE-RX device [10, 11] ($R = 1.72$ m, $a = 0.45$ m) had a multiple layer construction composed of a close fitting thin shell (1 mm copper) ($b/a = 1.08$) and a thick shell (5 cm aluminium). It operated from 1998 up to 2007 [12].

RELAX is a low aspect ratio RFP [13] ($R = 0.51$ m, $a = 0.25$ m, maximum plasma current = 0.12 MA) operated at the Kyoto Institute of Technology, Japan and is mainly dedicated to the study of single helicity plasmas. RELAX is exploring MHD stability and the possibility of generating bootstrap current at a low aspect ratio [14].

KTX [15–17] started its operation in 2015 at the University of Science and Technology in Hefei, China, ($R = 1.4$ m, $a = 0.4$ m). It has a magnetic boundary similar to that of RFX-mod and is foreseen to produce a plasma current of 0.5 MA in phase I and 1 MA in phase II, with an MHD active control system. KTX is somewhat smaller in size than MST, but it utilizes an air core, like RFX, in order to achieve higher plasma current levels. The unique double-C design of the KTX vacuum vessel allows access to the interior of the machine for rapid first-wall (FW) modifications and investigations of power and particle handling. The KTX device provides a platform to develop diagnostic tools [18] and to investigate the fluctuation induced transport in high temperature plasmas.

An upgrade of RFX-mod machine assembly has been designed (dubbed RFX-mod2) and is presently being implemented [19–21]. The primary goal of the upgrade, discussed in chapters 4 and 7, is to modify the magnetic front-end in order to reduce the amplitude of tearing modes and thus mitigate the magnetic chaos and the localized PWI due to wall-locking).

2. Experimental properties of the RFP configurations

2.1. The axisymmetric RFP equilibrium

The RFP equilibrium is a paramagnetic configuration: the plasma currents arrange in a way to ensure MHD equilibrium and generate most of the toroidal flux. The amount of toroidal flux converted from the poloidal flux by the dynamo process (explained in section 3.3) is proportional to the plasma current: the inverse of the proportionality constant is the figure of merit that the RFP configuration shares with all pinch configurations,
Table 1. Main parameters of present RFP devices and the large TPE-RX device that operated until 2007. Cross sections of devices are shown in figure 1.

<table>
<thead>
<tr>
<th>RELAX</th>
<th>EXTRAP-T2R</th>
<th>TPE-RX</th>
<th>MST</th>
<th>KTX phase 1 (phase 2) (projected)</th>
<th>RFX-mod</th>
</tr>
</thead>
<tbody>
<tr>
<td>R/a (m)</td>
<td>0.5/0.25</td>
<td>1.24/0.18</td>
<td>1.72/0.45</td>
<td>1.5/0.5</td>
<td>1.4/0.4</td>
</tr>
<tr>
<td>Shell time τ_w (ms)</td>
<td>5</td>
<td>12</td>
<td>300</td>
<td>1100</td>
<td>40</td>
</tr>
<tr>
<td>Max I_p (MA)</td>
<td>0.12</td>
<td>0.3</td>
<td>0.45</td>
<td>0.6</td>
<td>0.5 (1)</td>
</tr>
<tr>
<td>Pulse length (ms)</td>
<td>3.5</td>
<td>100</td>
<td>80</td>
<td>100</td>
<td>2</td>
</tr>
<tr>
<td>Max T_e (keV)</td>
<td>0.3</td>
<td>0.3</td>
<td>1</td>
<td>2.0</td>
<td>0.3 (0.8)</td>
</tr>
<tr>
<td>n_e range 10^{19} m^{-3}</td>
<td>0.2–3</td>
<td>0.5–2</td>
<td>0.5–1</td>
<td>1–4</td>
<td>1–10</td>
</tr>
</tbody>
</table>

hence the name pinch parameter \( \Theta \):

\[
\Theta = \frac{B_\theta(a)}{\langle B_\phi \rangle}
\]  

(2.1)

where \( \langle B_\phi \rangle \) is the cross section averaged toroidal magnetic field and \( a \) is the plasma minor radius. The distinguishing feature of the RFP is the reversal of the toroidal field at the edge: its magnitude, normalized to the toroidal flux is defined as the reversal parameter \( F \):

\[
F = \frac{B_\phi(a)}{B_\theta}
\]  

(2.2)

These two externally measurable parameters are used to characterize the RFP operations and allow comparing experiments of different sizes. The magnetic safety factor at \( a \) is proportional to the ratio of \( F \) and \( \Theta \), \( q(a) = (a/R)(F/\Theta) \). Experimentally, \( F \) and \( \Theta \) are correlated and follow some trajectory in the \((F, \Theta)\) plane [22]. At the beginning of the setting-up phase of a plasma pulse \( F = 1 \) and \( \Theta = 0 \); as plasma current is initiated and increased due to the induced loop voltage, the plasma begins to generate toroidal flux, so that \( F \) decreases as \( \Theta \) increases; in particular, above a \( \Theta \) threshold \( F \) becomes negative.

The axisymmetric part of the parallel current profile \( \mu = \mu_0 J \cdot B / B^2 \) measured in experiments (see e.g. in figure 1b of [1], or figure 42) is often approximated in a cylindrical geometry by the following family of functions

\[
\mu(r) = \mu(0) \left(1 - \frac{r}{a} \right)^\alpha
\]  

(2.3)

characterized by a maximum current at the magnetic axis decreasing to zero at the edge [23, 24]. The exponent \( \alpha \) allows to parametrize how much the current profile is peaked. This model is used to derive the internal magnetic profiles and to perform linear stability analysis. A toroidal version of this model has been developed [25] and has been benchmarked with the axisymmetric version of VMEC adapted for the RFP [26]. A toroidal equilibrium reconstruction code MSTFIT that solves the Grad-Shafranov equation (analogous to EFIT widely used for tokamak research) has been also developed for MST and adapted for other RFP experiments [27].

The toroidal and poloidal components of the RFP magnetic field are of comparable amplitude, in contrast with the tokamak where a stabilizing toroidal magnetic field larger than the poloidal one is needed. This is shown in figure 2, where the radial profiles of the poloidal and toroidal components of the magnetic field (normalized to the on-axis magnetic field) are reported for cylindrical RFP and tokamak configurations.

When including the flow term in the axisymmetric magnetic equilibrium, the plasma toroidal velocity in RFX-mod plasmas changes sign towards the edge [28], consistent with experimental observations [29–31].

The RFP cylindrical safety factor \( q(r) = rB_\phi / RB_\theta \) is strongly sheared, with \( q(0) \approx a/2R \) at the magnetic axis and decreasing with minor radius such that \( q(a) < 0 \), where \( a \)
and $R$ are the minor and major radius respectively (figure 3). The toroidal field at the plasma surface is therefore directed opposite to the core, giving the RFP configuration its name. This has profound consequences in terms of plasma magneto-hydrodynamic (MHD) stability. In addition to ideal kink modes, which become RWM due to the finite conductivity of the shell, there are many resonant surfaces for $m = 1$ tearing modes with $n \geq 2R/a$, where $m$ and $n$ are the poloidal and toroidal mode numbers of the usual Fourier decomposition, and the $q = 0$ surface is resonant for $m = 0$ modes.9

2.2. The dynamics of the RFP configuration: from multiple helicity to quasi single helicity

Two widely different regimes, characterized by different $m = 1$ and $m = 0$ magnetic fluctuation spectra, have been observed in RFP experiments especially as long as higher plasma currents, and consequently higher electron temperatures, have been obtained. These regimes have been defined as multiple helicity (MH) and quasi single helicity (QSH), according to the shape of the measured spectrum of the $m = 1$ and of the $m = 0$ Fourier harmonics of the magnetic field measured at the edge of the experiments. They are largely dependent on the device magnetic boundary parameters such as aspect ratio, thickness of the conducting shell and its distance from the wall. Experimentally, MH states are characterized by the presence of many $m = 1$ Fourier components of the magnetic field with comparable amplitude. In the QSH states one single $m = 1$ Fourier component grows up becoming dominant over the other ‘secondary’ modes, whose amplitude is smaller than in MH.

2.2.1. Multiple helicity regime. In both early RFP experiments [32] and 3D MHD computations [33–35], the Fourier modes measured at the edge had poloidal mode number $m = 1$ and a range of toroidal mode numbers, $n$, resonant in the plasma core given the $q$ profile (figure 3).

When in the setting-up phase of the plasma discharge the profile of the safety factor, initially flat, evolves through a series of non-monotonic profiles (until the field reversal), the

\[ q(r) = \frac{rB_0}{RB_0} \]

Figure 3. Axisymmetric safety factor profile $q(r) = rB_0/BR_0$ for a typical RFP with an aspect ratio $A = 3$. The location of the $m = 1$ and of the $m = 0$ resonant surfaces are shown.

Figure 4. Prototypical RFP sawtooth, ensemble averaged from a large number of similar events recorded on the MST RFP. Time $t = 0$ is a marker at the peak of sawtooth activity. Each panel represents the change in a key parameter during a sawtooth event. (a) Relative amplitude of magnetic fluctuations, illustrating tearing mode growth; (b) toroidal magnetic flux density, illustrating RFP dynamo activity; (c) heat flux, indicative of magnetic stochastic transport; (d) plasma flow, indicative of momentum transport; and (e) ion temperature, demonstrating ion heating.

local minimum at the edge is a source of kink instability [36]. The series of current-driven modes with $m = 1$ and increasing $n$ occurring as soon as the mode resonance appears, generate the poloidal current required to create the toroidal flux [37]. This process bears some similarity with the Kadomtsev reconnection model in the tokamak [38].

In the sustainment phase, toroidal flux generation is observed to occur either continuously or in discrete events [39, 40] referred to as sawtooth crashes [41], dynamo reconnection events or discrete reconnection events (DREs) in different contexts and papers (an example is shown in figure 4). The impulsive features of these events have been reproduced and studied in visco-resistive MHD simulations [42–45]. The frequency of such events depends on the pinch parameter $\Theta$ [41, 46] and on the magnetic boundary. Experimental observations showed such events as associated to large amplitude bursts of $m = 1$ and $m = 0$ low $n$ modes (figure 4(a)) [47–52]. The cyclic relaxation of the current profile has been measured via Faraday rotation, which shows repetitive flattening of $\mu(r)$ inside the $q = 0$ surface and peaking near the plasma edge [53]. The generation of the toroidal flux is a local, toroidally asymmetric phenomenon as a consequence of the nonlinear phase alignment [54]. Toroidally localized, magnetic field aligned current sheets have been measured in RFX-mod [55] and in MST [56, 57].

The global impact of DREs depends on the nonlinear interaction between $m = 1$ modes resonant in the core and

9 Due to the chosen conventions of the orientation of the coordinate system and of the 2D Fourier decomposition, and due to the relative orientation between the toroidal field and the plasma current, in RFX-mod and in EXTRAP-T2R internally resonant tearing modes toroidal numbers are negative, so that $n \leq -2R/a$. 
\( m = 0 \) modes resonant near the plasma surface. For ‘non reversed’ plasmas with \( F = 0 \) (vanishing toroidal field at \( r = a \)), the \( m = 0 \) resonance is removed from the plasma, and the global impact of reconnection events is muted [58]. Operation with \( F \approx 0 \) is an empirical control that helps stimulate QSH, likely a consequence of reduced coupling between \( m = 1 \) and \( m = 0 \) modes.

While visco-resistive MHD describes well the sawtooth relaxation dynamics as a consequence of tearing instability, flattening of the current profile has also been modelled as a kinetic effect related to magnetic chaos [59, 60]: in the radial domain of chaotic magnetic field, transport is fast and equilibrium almost force-free, so that \( \mathbf{J} = \mu(r)\mathbf{B} \), where \( \mu \) may be radius dependent. Setting this in \( \nabla \cdot \mathbf{J} = 0 \) implies \( \nabla \mu(r) \cdot \mathbf{B} = 0 \), which shows that \( \mu \) must be constant along the field lines.

The scaling of magnetic fluctuation amplitudes with the Lundquist number \( S \) has been studied in a number of RFP devices in the MH regime [61–64]. The Lundquist number \( S \), that in RFPs is usually \( S \gg 1 \) is defined as the ratio of the resistive diffusion time \( \tau_R = \mu_0 L^2/\eta \) (being \( \eta \propto Z_e T_e^{1/2} \) the plasma resistivity and \( L \) the scale of the system) to the Alfvén time \( \tau_A = L/V_A \), where \( V_A = B/\sqrt{\mu_0} \) is the Alfvén velocity and \( \rho \) the mass density. The magnitude of magnetic fluctuations decreases with \( S \) in RFP plasmas, so this scaling is important to characterize transport related to magnetic stochasticity. A sufficiently strong scaling is necessary so that RFP plasmas are not limited by stochastic transport.

The generation of the toroidal flux through the nonlinear interaction of the tearing modes is a local, toroidally asymmetric phenomenon [54].

During DREs, the radial profile of the toroidal plasma rotation can change rapidly and spontaneously, representing a rapid radial transport of toroidal momentum. This effect is about two orders of magnitude more rapid than what would be expected from classical viscosity and it is explained by nonlinear interactions involving triplets of tearing modes producing internal torques that redistribute momentum and therefore lead to flow profile flattening [68].

An important consequence of the nonlinear coupling of tearing modes is the formation of a toroidally localized deformation called slinky (inspired by the kink shape formed when a slinky toy is twisted), or locked mode (LM) in RFP literature [92], in EXTRAP-T2R [93] and in RELAX [94]. In TPE-RX, long-lasting highly reproducible QSH states were obtained by applying a delayed reversal of the toroidal field at the edge [95], and were systematically triggered by applying a small positive pulse in the initially weakly reversed edge magnetic field [96]. The end of the QSH phase came with a distortion of the \( m = 0 \) magnetic island chain, and the occurrence of a chaotic region, expelling energy at the location of the LM [97]. In RELAX, QSH states lasting above 30% of the flat-top duration can be obtained [98].

The conditions to observe QSH states are similar in all RFP experiments: they occur at shallow reversal (i.e. \( F \) near 0) and the dominant mode typically resonates near the magnetic axis [96, 98–100]. Moreover, there is a higher probability to obtain QSH spectra with higher plasma current [65, 89, 101, 102], and their persistence (i.e. the overall fraction of the flat-top characterized by QSH states) increases with current too. In RFX-mod at plasma current above 1.5 MA the latter can exceed 90%
of the plasma current flat-top [102, 103]. The amplitude of the dominant mode and that of the secondary modes scale in opposite ways with the current and with the Lundquist number $S$, as observed both in RFX-mod and in MST and reported in figure 6 [104, 105]. Long lasting QSH states observed in RELAX, despite the low plasma current, suggest that a low aspect ratio configuration could be a favourable condition, due to the scarcer density of the central $m = 1$ tearing modes in this case [98, 99].

Two kinds of QSH states have been observed. When the amplitude of the dominant magnetic mode is up to a few percent of the central magnetic field, and several times that of the other secondary modes [90, 106] the magnetic field topology displays two magnetic axes (figure 7(a)). The first one corresponds to the unperturbed equilibrium one while the second axis is related to the island O-point [107]. Such QSH states are termed double axis (QSH-DAx) states [108].

In the highest RFX-mod plasma current regimes, where secondary modes reach the lowest amplitudes, QSH states with the expulsion of the magnetic separatrix between the helical island corresponding to the dominant mode and the axisymmetric field have been observed [107, 109]. The initially small magnetic island related to the dominant mode in QSH-DAx is replaced by a large helical structure without separatrix (see figure 7(b)). These QSH states featuring a single helical magnetic axis are dubbed QSH-SHAx states [110]. QSH-SHAx states are obtained either spontaneously, at plasma current above ≈ 1.4 MA [110], or stimulated by an external magnetic perturbation, or triggered by oscillating poloidal current drive (OPCD) [107, 111]. In [112] it was predicted that the disappearance of the magnetic separatrix in QSH states corresponds to chaos healing.

QSH-SHAx states also occur in MST, at similar Lundquist number $S$, even though lower plasma currents (about 0.5 MA) than in RFX-mod, with the dominant $m = 1, n = 5$ mode as large as 8% of the axisymmetric field [91, 105, 113].

Both QSH-SHAx and QSH-DAx periods end by sawtooth crashes, similar to the ones typical of the MH regimes, where the amplitude of secondary modes increases and the dominant mode decreases (figures 5(b) and 16(c)), but the system goes back immediately to QSH. Moreover, a small sawtooth-like activity still occurs during long lasting QSH states in RFX-mod, characterized by small crashes not leading to a back-transition to MH states [114].

2.2.3. QSH states stimulated by a helical boundary. The possibility to increase the duration of QSH phases with the application of a stationary or rotating magnetic perturbation (MP) to a prescribed slow (up to 50–100 Hz) frequency has been investigated in MST, EXTRAP-T2R and in RFX-mod. MST adopted a pre-programmed approach as in resonant magnetic perturbation (RMP) or resonant field amplification (RFA) experiments (see e.g. [115] or [116, 117]), where the currents flowing into the control coils are pre-programmed. Instead in RFX-mod the measured radial field is feedback controlled at a prescribed reference level [118]. As a result, the application of an $m = 1, n = −7$ field in plasmas with shallow reversal, strongly increases the persistence of the helical state [119], provided that the reference amplitude is above some threshold. Furthermore, QSH states can be obtained at higher densities $(n_{eG} \approx 0.5$, with $n_G \approx \frac{10^{20}}{m^3} = I_p [\text{MA}] / \pi a^2$ Greenwald density) than in the spontaneous case [120]. While first RFX-mod experiments with the application of an external 3D field resulted in a significant increase (up to 50%) of the loop voltage necessary to sustain the plasma current, such effect has then been minimized with a proper tailoring of the radial field [121]. In EXTRAP-T2R, an increase of the persistence of QSH periods up to 10% of the discharge and larger helical core structures occurred by applying a feedback controlled RMP with the corresponding helicity [122].

As predicted by numerical simulations [123], by providing a corresponding helical boundary condition in RFX-mod, it is
Figure 8. Poincaré maps from simulations at $\theta = 0$ with a 2% MP as boundary condition; (top) MP with helicity $m = 1, n = -6$, non resonant; (middle) MP with helicity $m = 1, n = -6$; (bottom) MP with helicity $m = 1, n = -8$. The $n = -6$ case shows larger areas of conserved surfaces. Reproduced courtesy of IAEA. Figure from [125]. © 2017 EURATOM.

possible to excite non internally resonant QSH $n = 6$ states [124]. When a non-resonant boundary condition is applied the area with well conserved magnetic surfaces is larger and the transport of magnetic field lines is reduced (figure 8) [125]. This has been related to the existence of Cantori, encompassing the region characterized by conserved magnetic surfaces, which act as barriers to transport of magnetic field lines [126].

The application of 3D fields to improve the properties of the configuration is not unique to the RFP configuration. The motivations for introducing the 3D magnetic fields include plasma confinement and stability optimization (stellarators), plasma control (tokamak edge-localized modes and tokamak and RFP resistive wall instabilities) and engineering/economic constraints (limited number of toroidal field coils, asymmetrical particle/energy sources and non-uniform ferrous steel structure near the plasma). A survey of common tools developed in the stellarator community can be found in [127].

2.3. The helical equilibria for the QSH-SHaX and QSH-DAx states

The internal topology of the magnetic field is reconstructed with several approaches. The first one is a perturbative approach. Marginally stable eigenfunctions [128, 129] were matched to the edge magnetic fluctuation spectrum and implemented in field line tracing codes [130, 131]. This has been done in several RFP devices [101, 132, 133]. Such codes have also been benchmarked with the NEMATO tool, which solves the magnetic field line equation on a 3D grid in a volume preserving integrator scheme [134]. In QSH-DAx Poincaré plots were displaying magnetic islands emerging out of stochastic domain closely corresponding to independently measured ‘bean’-like hot structure featuring strong electron temperature gradient at the edges (see figure 9) [87–90, 135]. In Poincaré plots of QSH-SHaX states, wider regions separated from outer stochastic domain appear into the core.

A much faster way to reconstruct internal topology is based on the SHEq code [136], where a system of helical flux coordinates is implemented by considering only the dominant mode radial profile as a perturbation on the axisymmetric equilibrium. The magnetic field reconstruction by SHEq has been compared with a direct measurement of the internal magnetic field in MST [91, 113]. In RFX-mod, the reconstructed equilibrium shows that the electron temperature and density are constant along the helical magnetic flux surfaces, see figure 10 [109].

These helical equilibria have also been reconstructed with the VMEC equilibrium code, originally developed in the stellarator community, and applied and enhanced in order to deal with the reversal of the toroidal field [26, 137]. The resolution of the equilibrium reconstruction has been then improved by coupling VMEC with the V3FIT code [138], able to include internal and external measurements as constraints. As an example, figure 11 shows the improvement of the magnetic reconstruction when the pressure profile is included as a constraint.
A fundamental property of QSH-SHAx equilibria is that the \( q \) profile is non-monotonic: it goes through a maximum located in the vicinity of the temperature gradient region, as shown in figure 12(b). Non-monotonic \( q \) profiles have also been found in numerical SH simulations by the SPECYL code [139, 140]. As the maximum value of the helical \( q \) is lower than the axisymmetric one, the innermost resonant mode of the axisymmetric safety factor \((m = 1, n = -7 \text{ in RFX-mod for example})\) is not resonant anymore according to the helical \( q \) profile (therefore, in RFX-mod, the innermost resonant mode becomes \( m = 1, n = -8 \)).

The magnetic topology of both QSH-SH Ax and QSH-D Ax states of RFX-mod can be reproduced by a one parameter fit of a minimally constrained equilibrium model using only five parameters in a two-domain generalization of Taylor’s theory (that will be summarized in section 3.1). Both states appear as a consequence of the formation of a transport barrier in the plasma core, in agreement with experimental results [141].

It is worth mentioning that helical core equilibria states are also found in tokamaks, associated to the non-monotonic \( q \) profile occurring in hybrid mode operations, as obtained in 3D equilibrium calculations [142, 143].

3. Magnetic self-organization in the RFP

The first experimental evidences of a RFP configuration as an example of self-organization dates back to the toroidal pinch research in the 1960s, with the observation of the quiet periods in the ZETA machine [144, 145]: low magnetic turbulence periods associated to the reversal of the toroidal field near the plasma edge.

The toroidal field reversal is the key ingredient for the ideal stability of internal modes, combined to a sufficiently close fitting shell for suppressing external \( m = 1 \) modes [36, 146]. Later linear ideal and resistive MHD stability analysis with more realistic profiles [23] and boundary conditions [147] confirmed and enhanced these findings.

The emergence of a quiet and ordered configuration out of the turbulent discharge setting-up was considered an example of a self-organized process, as defined by Hasegawa [148]. A weakly dissipative continuous medium in which energy and one or more other quantities are conserved in absence of dissipation, evolves towards an ordered quasi-stationary state, obtained by a constrained minimization. According to the Taylor’s conjecture in [149], magnetic helicity is the other conserved quantity.
Section 3.6 summarizes the experimental measurements of MHD dynamo and the evidences of a two-fluid contribution from the Hall dynamo, while flow and momentum transport studies are recalled in section 3.7. Finally, section 3.8 presents a new interpretation of the MHD dynamo as an electrostatic dynamo able to unify the description of both multiple and SH plasmas. Section 3.9 summarizes further theory aspects.

3.1. Taylor’s theory of the RFP

The first attempt to explain the reversal of the toroidal field was performed within the Taylor’s theory [149, 154] as being due to some unspecified MHD turbulence. Such a turbulence in a weakly resistive plasma causes reconnection of magnetic field lines so that the plasma rapidly relaxes towards a minimum energy state on ideal MHD time scales (of the order of the Alfvén time), which are faster than the classical resistive diffusion time scales.

The mathematical properties of this state are obtained by a variational argument, i.e. by minimizing the magnetic energy subject to the conjecture that, in presence of weak resistivity, the conservation of the total helicity is a constraint to be satisfied. The motivation behind Taylor’s conjecture is that the finite plasma resistivity causes reconnection of magnetic field lines and breaks the conservation of the infinity of helicity invariants of the ideal magnetohydrodynamics found by Moffatt [155]. For an isolated plasma, enclosed in a perfectly conducting wall, this minimization leads to a force free equilibrium field [156].

Despite the appeal of the Taylor conjecture, there is as yet no rigorous justification for it, although various heuristic arguments have been advanced [157]. Still, it has been successfully used as a framework for the interpretation of the RFP experimental and numerical behaviour of the RFP, especially during the DREs [158]. For example, from global magnetic measurements combined with equilibrium modelling, the magnetic energy and helicity have been inferred in the MST during DREs [24], revealing that the helicity changes relatively little while magnetic energy reduction is significant.

In Taylor’s relaxation theory the magnetic field was found to obey the force-free equilibrium equation

$$\nabla \times B = \mu B$$  \hspace{1cm} (3.1)

where $\mu$ is constant throughout the whole plasma. This implies that the plasma current density $J$ is aligned to magnetic field $B$. In cylindrical geometry, solutions of the force free equation (3.1) are Bessel functions (hence the name Bessel function model, BFM) and the $(F, \Theta)$ parameters describe a unique set of relaxed state equilibria similar to experimental observations. The theory predicts that, above a certain $\Theta$ value, an equilibrium with reversed toroidal field occurs.

On the other hand, in the light of visco-resistive MHD simulations performed from 90s, several aspects of Taylor’s theory have been criticized [159]. Among them: the states predicted by the theory do not exist, neither experimentally, nor numerically; the timescale of relaxation turns out to be intermediate between ideal and resistive; the fluctuations involved in these processes are large scale (tearing and resistive modes). Also,
3.2. The need of a dynamo for the RFP

While Taylor’s theory approximately describes the properties of the RFP discharges, it is not a self-contained theory that describes all of the detailed processes generating field reversal, i.e., the reconnection and relaxation mechanisms. Moreover, even though the axi-symmetric RFP is a steady equilibrium, its ohmic sustainment by an induction toroidal field cannot occur in axi-symmetric geometry [37], and therefore a steady-state axisymmetric solution should not exist. This is a result of the Cowling anti-dynamo theorem, which states that a symmetric magnetic field configuration cannot be maintained against resistive diffusion by a symmetric velocity field [37]. Figure 13 shows the radial profile of the parallel component of the mean applied electric field \( E_\parallel = \mathbf{E} \cdot \mathbf{B} / |\mathbf{B}| \) and compares it with the \( \eta J_\parallel \) term, clearly showing that the axisymmetric Ohm’s law is violated. In particular in the core the applied electric field is larger than required \( (E_\parallel > \eta J_\parallel) \), while in the edge it is less and even directed in the opposite direction. The dynamo process generates an additional parallel current required to balance the Ohm’s law.

At the beginning of the 90’s two possible mechanisms underlying the RFP dynamo were investigated, partially supported by experimental observations [1, 160, 165]. The first one was based on the MHD dynamo, which relied on the local mean-field single fluid Ohm’s law, in which the coupling of velocity and magnetic fluctuations generated an average electric field \( E_\parallel \) sustaining the plasma current against resistive diffusion.

A second approach was based on the current due to the free streaming of electrons along stochastic field lines, the so called kinetic dynamo theory (KDT) [59]. The measurement of fast electrons at the edge of several RFP devices was the evidence suggesting this possibility. However, further theoretical studies and experimental investigations ruled out this interpretation (see section 4.8).

3.3. The MHD dynamo

Borrowing concepts from geophysical and astrophysical turbulent dynamo theories [162], the fields and currents are separated into mean \( \langle \mathbf{B}, \mathbf{E}, \mathbf{J} \rangle \) and fluctuating components \( (\mathbf{b}, \mathbf{v}, \mathbf{E}) \), where the angle brackets represent a mean-field flux surface average, which leads to an additional mean field term \( \langle \mathbf{v} \times \mathbf{b} \rangle_\parallel \) in the parallel to \( \mathbf{B} \) components of Ohm’s law [166]:

\[
E_\parallel + \langle \mathbf{v} \times \mathbf{b} \rangle_\parallel = \eta J_\parallel
\]

In particular, this term, called in subsequent RFP literature MHD dynamo, can be non-zero at the reversal surface, where the induction electric field is zero, and can drive the required poloidal current. As the velocity and magnetic field dynamics are coupled through the Lorentz force, the determination of the MHD dynamo term is a nonlinear problem. In [166] a simplified statistical approach was adopted based on mean-field electrodynamics [162]. The MHD dynamo term was assumed to be proportional to the mean field through the free parameter \( \alpha \) (the \( \alpha \)-dynamo approach [167]) and under some assumptions the equilibrium field was found to obey the force-free condition with constant \( \mu \), as in equation (3.1).

Self-consistent investigations on the dynamics of the fluctuations were based on 3D nonlinear visco-resistive MHD computations. Sykes and Wesson showed robust field reversal in a single fluid MHD simulation using a \( 14 \times 14 \times 13 \) grid [168]. They identified the basic mechanism of field reversal as the quasi-periodic nonlinear evolution of an \( m = 1 \) global kink instability in the presence of an externally applied toroidal voltage. An intuitive explanation of this result can rely on a simple toy model, which consists in a current-carrying resistive wire initially placed on the axis of a cylindrical flux conservor, whose spontaneous kink leads to the reversal of the edge axial magnetic field, and not to a ‘disruption’ [169, 170].

This mechanism was confirmed by subsequent numerical simulations, based on the elementary cylindrical visco-resistive compressible nonlinear MHD model in the constant-pressure, constant-density approximation with a forced SH [33, 35, 171–173].

3.4. The standard paradigm: the MHD dynamo is inherently multiple helicity

The initial numerical 2D SH states were deemed too far from the experimentally observed MH plasmas. When computational physicists, taking advantage of the increase of the power of computers, began to deal with the full 3D problem, the parameters of initial numerical simulations led them into the MH state. The MHD dynamo term was inherently due to several nonlinearly interacting global tearing modes. Fully 3D MHD numerical simulations with multiple helicities were able to reproduce field reversal and the quasi-periodic RFP sawtooth dynamics (see section 2.2) [42, 160, 173–176].
An example of the radial profiles of the MHD dynamo electric field $E_\parallel = -(\mathbf{v} \times \mathbf{b}) \cdot \mathbf{B} / |\mathbf{B}|$ is shown in figure 13 along with all the terms of the Ohm’s law for a quasi-stationary solution of a 3D visco-resistive MHD simulation [177]. Figure 14 shows the terms $E_\parallel$ and $\eta J_\parallel$ derived from an experimental equilibrium reconstruction [178]. The parallel Ohm’s law becomes $E - E_\parallel = \eta J_\parallel$. Both the applied $E_\parallel$ and parallel MHD dynamo field $E_\parallel$ modify the parallel current profile $J_\parallel$ which, in turn, determines the stability and the level of dynamo fluctuations. By itself, the applied $E_\parallel$ would create a peaked $J_\parallel$ profile and could not sustain a reversed magnetic field (note that $E_\parallel$ near the edge is in a direction to suppress $J_\parallel$). The positive values of $E_\parallel$ near the axis suppress parallel (axial) current, i.e. act as an anti-dynamo, while negative values of $E_\parallel$ near the edge drive parallel (mostly poloidal) current.

The competition between these two processes has been used to interpret the common observation of the DREs, described in section 2.2. From linear theory [23, 160], the peaked $E_\parallel$ causes a gradient in $J_\parallel$ that destabilizes tearing modes. The parallel dynamo field $E_\parallel$ generated by the tearing modes and their associated flow fluctuations tend to flatten the $J_\parallel$ profile, and in so doing, sustain the reversed toroidal field. Therefore, nonlinear saturation of the tearing instability is interpreted as a balance between these two competing processes, and an equilibrium is maintained near marginal stability.

The nonlinear interaction between $m = 1$ and $m = 0$, low $n$ resonant modes was investigated in [42], by implementing a post-processing energy diagnostic to identify the path taken in Fourier space by the poloidal magnetic field energy as it is converted to axial magnetic field energy by the dynamo process. Such coupling was subsequently verified experimentally (see section 2.2.1).

Numerical simulations performed with a Lundquist number close to the experimental ones in [179] highlighted the presence of toroidally localized dynamo current sheets generated during reconnection events, which are toroidal in the core (anti-dynamo) and mainly poloidal in the edge.

3.5. The transition from multiple helicity to single helicity

Since 1990 and before QSH experimental observations, 3D numerical simulations revealed a bifurcation of the magnetic configuration from the MH state to the SH state, occurring when viscosity is increased and coming with magnetic order [179–182].

An appropriate rescaling of time and velocity in the MHD equations shows that, for given radial distributions of plasma resistivity $\eta$ and viscosity $\nu$, the dynamics is ruled by the magnetic Prandtl number $P = \nu(0)/\eta(0)$ and by the dissipation of the system $d = \eta(0)\nu(0)$, also quantified by the Hartmann number $H = d^{-1/2}$, a classical dimensionless number in conductive fluid MHD. Dissipation is the dominant parameter when the inertia term becomes negligible, which happens for a large range of simulation parameters [183]. These simulations indicate $d$ (or $H$) as the control parameter independently of $P$ and $\Theta$. Therefore, the previous high viscosity simulations turned out to be in reality high dissipation (or small $H$) simulations. As usual in nonlinear dynamics, an increase of dissipation is favourable to a laminar behaviour of the system. Depending on the initial conditions, two nearby different helicities were found to be selected by the plasma when relaxing to SH [183, 184].

Incidentally, it has to be mentioned that, while in the first simulations the plasma current and the axial magnetic flux were taken as constant, which implies the constancy of the pinch parameter $\Theta$ ([184] and references therein), from 2008 on simulations with the SpeCyl and PIXIE3D codes were performed with a constant loop voltage instead of a constant $\Theta$ [185, 186].

3.5.1. The intermediate regime: the QSH states.

As $d$ decreases, the pure helix (single helicity, SH) develops a toroidally (axially) localized bulging (slinky); this bulged helix is interrupted by longer and longer intermittent phases, in which secondary modes with other helicities show up in the Fourier spectrum. The bulged helix corresponds to a QSH state where secondary modes are not zero but smaller than the dominant one. When at least two modes have similar amplitudes, the plasma switches to the MH state. With decreasing $d$, the duration of the QSH phases decreases, as well as the percentage of time where QSH dominates, and after a transition region corresponding to $10^3 \leq H \leq 10^4$, the system reaches a non-stationary MH regime [183]. This is shown in figure 15 where the energy of the $m = 0$, $n \neq 0$ modes is used as an order parameter for the SH–QSH–MH transition; indeed these modes vanish in the SH state, since they result from the beating of at least two different helicities. It is important to note that magnetic chaos sets in progressively during the transition.

3.5.2. The role of the helical boundary conditions in MHD simulations.

Though MHD simulations were able to predict QSH states, in the first SpeCyl calculations they had a temporal behaviour different from the experimental one. The reason is that the former had a vanishing radial magnetic field at the edge, while the latter correspond to a finite one, since in experiments the time discretization of a digital feedback

![Figure 14. Results from an equilibrium reconstruction illustrating that $E_\parallel \neq \eta J_\parallel$ in standard multiple-helicity RFP operation. In particular, the existence of a radial location with zero applied electric field and finite current density makes clear the need for dynamo drive. Reprinted from [178], with the permission of AIP Publishing.](image-url)
control system prevents the achievement of zero radial field at plasma radius [187]. A good agreement with the experimental QSH dynamics was in fact obtained in simulations including a small \( m = 1, n = -7 \) boundary condition (MP) with an amplitude similar to the experimental one (about 1.5% of the central magnetic field) and with \( S = 10^7 \), consistent with experimental Lundquist number [188]. An example is shown in figure 16. A more recent work, including simulations of both spontaneous and driven experimental QSHs (see section 2.2.3), shows that the similarity with experiments bears on several features [189]: (i) the dependence of the quasi-periodic behaviour of QSH states and of its amplitude on the MP amplitude, (ii) the dependence on the MP amplitude of the position of the shear reversal and maximum \( q_0 \), (iii) the fact that the amplitude of the dominant mode is independent of dissipation, whereas the amplitude of secondary modes decreases with increasing inverse resistivity and viscosity (Lundquist and viscous Lundquist numbers). Numerically, a stationary state is reached for an MP of 10%.

Very recent extended analysis including a large set of simulations actually shows that in presence of a finite radial magnetic field at the edge, a second low dissipation region exists where SH states also develop [126].

3.5.3. Necessary condition for field reversal of the single helicity ohmic state. The SH ohmic state can be described analytically as a small helical perturbation of an axisymmetric ohmic pinch with small edge axial magnetic field and conductivity. This description uses the pinch–stellarator equation, which shows how the axial field of a cylindrical SH RFP evolves radially because of the pinch effect and of a stellarator contribution due to the helical deformation of the plasma [182, 190, 191]. Reversal is due to the edge stellarator contribution, which is more efficient when the edge pinch contribution is weak. The latter condition corresponds to a resistive edge plasma. Furthermore, simulations performed with a flat resistivity profile display a reduction of the dynamo action, which brings to marginally-reversed or even non-reversed equilibrium solutions [192–195]. This occurs because with a larger resistivity at the edge, the electrostatic field is larger too, in order to balance parallel Ohm’s law. The electrostatic drift due to such larger electrostatic field provides an enhanced dynamo action, which is sufficient to sustain the reversed configuration [192]. The edge resistivity of RFPs is high enough for the reversed axisymmetric part of the toroidal field to be rather constant in the reversal domain after a steady decrease in the non-reversed one. The pinch–stellarator equation enables the
derivation of a necessary criterion for the reversal of the edge axial field [191], which in combination with MHD simulations with a small RMP shows that a finite edge radial magnetic field of given helicity is favourable for field reversal. The criterion is found to be satisfied in RFX-mod when a finite edge radial magnetic field is applied, but only marginally in spontaneous QSH states [191].

3.6. Measurements of the MHD dynamo

The MHD dynamo has been measured both in MH and in QSH plasmas in several RFP devices. Experimental tests of the MHD dynamo theory have been performed by measuring terms of the mean field Ohm’s law equation. For conciseness we write here the two fluid version, following [39]:

$$ E + v \times B = \eta J + \frac{J \times B}{ne} - \frac{\nabla p_e}{ne} - \frac{m_e}{e} \frac{dv_e}{dt} \tag{3.3} $$

where $v_e$ is the electron fluid velocity, $v \approx v_i$ (the ion fluid velocity), $p_e$ is the electron pressure, $n$ is the plasma density, $m_e$ is the electron mass, and $e$ is the electron charge. The last three terms on the right-hand side of equation (3.3) are the Hall term, the gradient of electron pressure, and the electron inertia: they represent effects beyond standard single-fluid MHD.

By splitting the quantities in equation (3.3) into mean and fluctuating parts, averaging over an axisymmetric flux surface, and taking parallel components, the generalized parallel Ohm’s law becomes

$$ \eta J_{||} = E_{||} + \langle \tilde{v} \times \tilde{b} \rangle \parallel - \langle \tilde{j} \times \tilde{b} \rangle \parallel / ne \tag{3.4} $$

where the last term on the right-hand side is referred to as the Hall dynamo. Note that the electron inertia term is usually neglected as small, and the parallel component of the pressure gradient term vanishes upon axisymmetric poloidal and toroidal averaging. In the parallel two-fluid Ohm’s law $v_i$ and $v_e$ are combined in the usual relations for the fluid velocity $v = (m_i v_i + m_e v_e)$ and current $J = e(n_i v_i - n_e v_e)$ [196]:

$$ \eta J_{||} - E_{||} = \langle \tilde{v} \times \tilde{b} \rangle \parallel \approx \langle \tilde{v}_e \times \tilde{b} \rangle \parallel \tag{3.5} $$

The appearance of only electron velocity in the rhs is the consequence of the fact that Ohm’s law in two fluids MHD represents the force balance equation for the electron fluid.

The Hall dynamo term is a two-fluid MHD effect: it arises when the ion flow does not follow electron flow, which implies that electrons and ions experience different forces and therefore the electron pressure force can drive flows only in electrons [167]. The effect of electron pressure fluctuations can be made explicit, so that it can be related to local measurements. Extracting the perpendicular component of the electron flow velocity and inserting it into equation (3.4), two terms can be measured by insertable probes become apparent:

$$ \eta J_{||} - E_{||} = \langle \tilde{E}_e \cdot \tilde{b}_e \rangle \parallel + \langle \tilde{V}_e \cdot \tilde{b}_e \rangle \parallel \tag{3.6} $$

Figure 17. Comparison of the dynamo electric field $\langle \tilde{E} \cdot \tilde{b} \rangle$ to $\eta J_{||} - E_{||}$ from an ensemble of sawtooth relaxation events in the MST RFP. Time $t = 0$ is a reproducible time marker for aligning individual sawtooth events during ensemble averaging. Reprinted from [196], with the permission of AIP Publishing.

The first term arises from the fluctuating $E \times B$ drift and is called the MHD dynamo, and the second term is a two fluid effect called the diamagnetic dynamo, related to the electron diamagnetic drift. The diamagnetic dynamo is therefore a manifestation of the Hall dynamo. In the single-fluid limit where ions and electrons drift together, the Hall dynamo is absent.

3.6.1. Edge measurements of MHD dynamo. The first direct measurements of the dynamo were made in the REPUTE RFP edge [197], and later in the MST edge [39], by directly correlating fluctuations in the electric and magnetic fields. In REPUTE, the measured dynamo electric field was far below that required to balance resistive dissipation, suggesting that two-fluid effects could have an impact. However, the measurements in MST revealed a dynamo electric field of direction and magnitude consistent with current sustainment, i.e. balancing parallel Ohm’s law. It is interesting to highlight that MHD dynamo effects were measured in a spheromak using a similar technique [198].

The Hall dynamo term was also directly measured by correlating current density and magnetic field fluctuations using probes in MST plasmas. It was found to be small in the far edge [199], but it increases and is sufficient to balance Ohm’s law near the reversal surface [200], showing that two fluid dynamics are important in the relaxation process.

Figure 17 shows the dynamo electric field measured in the edge of the MST MH plasma ($\rho/a = 0.9$), compared to the other terms in Ohm’s law, where the resistivity $\eta$ is calculated from measured local electron temperature $T_e$ and with estimated effective charge $Z_{eff} = 2$ (the effective charge $Z_{eff}$ is defined as $Z_{eff} = \sum_i n_i Z_i^2 / \sum_i n_i Z_i$, with $i$ ranging over all the ionic species present in the plasma). The dynamo electric field increases during the sawtooth event to overcome the inductive back reaction produced by rapid changes in the plasma current profile.

Measurements performed in TPE-1RM20 edge over a wide range of electron collisionality, (defined by the ratio of electron mean free path to the plasma radius) shed light on the different results [201]. Depending on the collisionality regime, the MHD dynamo dominates in the collisionless region, and the diamagnetic dynamo dominates in the collisional regime,
indicating that diamagnetic drift controls the electron flow’s correlation with the magnetic field. This finding explains the original result from REPUTE [196]. The results from these measurements, performed in MH states, imply that the (turbulent) dynamo effect causes helicity transport and overall helicity is conserved [202].

Fast, direct and local measurements of ion velocity fluctuations $\tilde{v}$ were obtained by an insertable optical probe in MST [203]. The probe collected light from two perpendicular lines of sight limited by view dumps to produce radial localization of the measurements to 2–3 cm. The correlation analysis highlighted the contribution to the MHD dynamo of individual resonant tearing modes. The two lines of sight of the optical probe provided simultaneous measurements of the radial and toroidal components of the plasma velocity, $\tilde{v}_r$, $\tilde{v}_t$, which are the quantities arising in the parallel component of $\langle \tilde{v} \times \tilde{b} \rangle_\parallel = \langle (\tilde{v}_r \tilde{b}_t) - (\tilde{v}_t \tilde{b}_r) \rangle$. To obtain the complete dynamo product, the velocity fluctuations $\tilde{v}_r$, $\tilde{v}_t$ were correlated with the local measurements of the magnetic fluctuations $\tilde{b}_r$ and $\tilde{b}_t$, respectively. The parallel current profile $J_y$ was obtained using an insertable Rogowski coil [48] and the toroidal field profile was determined by an additional insertable magnetic pickup coil. Ensembles over repeatable plasma discharges $I_p = 210$ kA, with density $n_e = 1 \times 10^{19}$ m$^{-3}$ were considered.

The main finding of this work is that the MHD dynamo electric field balances the Ohm’s law in the outer region of the MST plasma up to the reversal surface, (qualitatively consistent with the former edge measurements), but not inside, where the Hall dynamo from two fluid dynamics becomes important [200, 204]. Correlation analysis with tearing modes also reveals that the edge structure of velocity fluctuations is mainly related to $m = 0$, low $n$ harmonics, which are correlated to tearing modes resonant at the reversal surface. A small dynamo field is measured also between crashes, confirming the continuous nature of the single-fluid MHD dynamo (at least at the edge).

3.6.2. Mode resolved core measurements of MHD dynamo. Core ion flow fluctuations have also been measured, initially by line-integrated passive spectroscopy [205] and subsequently with charge exchange recombination spectroscopy (CHERS) [206], allowing more localized observations of the flow fluctuations and of the single fluid MHD dynamo electric field. Correlation analysis with tearing modes extracted the single fluid MHD dynamo term and allowed a more detailed determination of its spatial distribution. Throughout the sawtooth cycle, the core fluctuations have been found to be mainly correlated with internally resonant modes ($m = 1, n = 5–7$ in MST). Figure 18 is an example of two components of the $\langle \tilde{v} \times \tilde{b} \rangle$ dynamo produced by the core-resonant $m = 1, n = 6$ tearing mode. This work also indicated that the velocity fluctuation eigenfunctions involved in the MHD dynamo were peaked at the rational surfaces, with widths similar to the ones expected for the islands related to tearing modes (figure 19).

3.6.3. Measurements of flow fluctuations and MHD dynamo in QSH states. The MHD dynamo in QSH states was first measured in MST [90]. Main results are: (i) the magnetic and velocity fluctuation spectrum have the same dominant wave number. (ii) The corresponding component of the velocity extends throughout the plasma volume and couples with magnetic fluctuations, producing a significant MHD dynamo electric field. (iii) The radial profile of this component is consistent with the one obtained in QSH runs of the SpeCyl MHD code [184]. A similar study performed in RFX-mod [207] showed that a helical plasma flow forms an $n = 1$ convective cell which in turn creates a localized sheared flow outside the region of high temperature gradient (figure 20). However, the experimental pattern revealed the maximum flow shear more external.

**Figure 18.** Two $m = 1, n = 6$ components of the dynamo, both of which suppress parallel current in the core of MST. The velocity and magnetic fluctuations reach peak coherence at the sawtooth event ($t = 0$) and are nearly in phase, maximizing the dynamo product. Reprinted from [205], with the permission of AIP Publishing.

**Figure 19.** Amplitude contours of the mode-resolved poloidal velocity fluctuation $\langle \sqrt{\tilde{v}^2} \rangle$. The velocity fluctuation amplitude for a particular n-mode peaks approximately at the impact parameter where that mode is expected to be resonant (estimated from the q profile). Reprinted from [205], with the permission of AIP Publishing.
displays a QSH modulation [208]. Probes have been used in particular its shear [119]. In RFX-mod, in fact, the edge flow (section 2.2.3) can modify significantly the flow profile and in flows remain peaked near their resonant surfaces.

tain a similar flow structure during QSH while all tearing mode to the magnetic spectrum. The dominate the flow spectrum over the entire radius, in contrast intervals are larger and broader than in MH, but they do not

dominate the plasma discharge. The flow is radially directed towards the helical magnetic axis in the core and changes the sign moving to the edge. The helical axis is represented by the reconstructed helical magnetic surfaces. Reproduced courtesy of IAEA. Figure from [207]. Copyright 2011 IAEA.

with respect to the predictions of 3D MHD simulations, possibly due to an ambipolar component of the helical electric field on top of the MHD one.

Localized flow measurements in MST [206] show that poloidal flows associated with the \( n = 6 \) mode during QSH intervals are larger and broader than in MH, but they do not dominate the flow spectrum over the entire radius, in contrast to the magnetic spectrum. The \( n = 7–12 \) modes, in fact, maintain a similar flow structure during QSH while all tearing mode flows remain peaked near their resonant surfaces.

Even weak external 3D fields, used to stimulate QSH states (section 2.2.3) can modify significantly the flow profile and in particular its shear [119]. In RFX-mod, in fact, the edge flow displays a QSH modulation [208]. Probes have been used in low current discharges (up to 0.45 MA) of RFX-mod, enabling the observation in stimulated QSH states of a modulation of the perpendicular components of the flow [209] concomitant with the density modulation along the toroidal angle [210, 211].

3.6.4. Mode resolved measurements of the Hall dynamo. In MST the Hall dynamo term has been measured in the plasma core using an eleven-chord laser Faraday rotation diagnostic and correlating the current and magnetic field fluctuations over an ensemble of several sawtooth cycles [212]. The Hall dynamo is found to be spatially localized in the region of the resonant surface, while an additional dynamo mechanism may be required to balance the induced electric field elsewhere. Such a picture is qualitatively consistent with two-fluids quasi-linear theory [213] which predicts a localized Hall dynamo and a more diffuse MHD dynamo away from the resonant surfaces. The Hall dynamo term is found to become large (up to 40 V m\(^{-1}\)) during the sawtooth crashes.

3.7. Plasma flow and momentum transport

Experimentally, during a DRE rapid momentum transport is found as a flattening of the radial profile of the toroidal and poloidal rotation, as mentioned in section 2.2.1 [68]. Such flattening has been explained by the generalization of Taylor’s relaxation theory to two-fluids, in which context both the electron and ion helicities are conserved separately [214, 215]. In the resulting relaxed state, the spatial constant \( \mu = J_\parallel / B \) is recovered from the standard MHD analysis and represents the dynamics of the electrons. Additionally, a new spatial constant appears, \( nV_\parallel / B \), representing the dynamics of the ions and describing the relaxed state of the parallel plasma momentum profile.

Momentum transport is tightly linked to MHD dynamo process. In fact, the mean-field parallel momentum balance equation is

\[
\rho \frac{\partial v_\parallel}{\partial t} = (\tilde{j} \times \tilde{b})_\parallel - \rho (\tilde{v} \cdot \nabla \tilde{v})_\parallel \tag{3.7}
\]

where the left-hand side represents ion inertia with \( \rho \) as the mass density. On the right-hand side are the fluctuation-induced Maxwell and Reynolds stresses, respectively. Note that both the Maxwell stress term in equation (3.7) and the Hall dynamo one in equation (3.5) contain the term \((\tilde{j} \times \tilde{b})_\parallel\), indicating that the dynamics of the parallel current and parallel flow are coupled when two fluid effects are important. In MHD simulations, the relaxation of the parallel momentum profile is driven by the same reconnection events causing relaxation of the parallel current profile [216]. Experiments in MST confirmed such a tight coupling: MST measurements using probes in the edge and Faraday rotation in the core identified large Reynolds and Maxwell stresses associated with fluctuations at the tearing mode scale. The force densities from these stresses are large compared to the plasma inertia, but they tend to oppose each other [200, 212] so that their effect is negligible.

Recent two-fluid computations [44, 45, 217] show that there is a significant contribution from the Hall dynamo and coupling to momentum relaxation. The total dynamo is a smooth superposition of MHD and Hall effects that together balance Ohm’s law, but each component varies considerably across the plasma radius, with each vanishing locally at particular locations. This is similar to the behaviour measured in MST plasmas.

3.8. A new interpretation of the MHD dynamo: the electrostatic dynamo

Both in SH and MH states, the spatially fluctuating component of the velocity may be seen as an electrostatic drift velocity related to an electrostatic field which is spatially modulated [218–220]. This is the basis for many of the probe measurements of the dynamo described above [39, 196, 197]. The electrostatic drift comes quite natural in SH. In fact, the helical displacement of magnetic surfaces produces a modulation of the parallel current density along each flux tube, which requires the build-up of a helical electrostatic potential producing the dynamo flow as an electrostatic drift (laminar MHD dynamo). This statement is supported by an elementary calculation in cylindrical geometry in [220]: in the parallel Ohm’s law (scalar product of equation (3.3) with \( B \), the \( v \times B \) term and the Hall term \( J \times B \) bring vanishing contributions. The
gradient of electron pressure vanishes too, because it is perpendicular to magnetic surfaces. Moreover, electron inertia is generally neglected in the RFP. If the magnetic field is known, the parallel Ohm’s law enables the calculation of the electrostatic potential \( \phi \) since \( E = E_0 \mathbf{e}_z - \nabla \phi \), with \( E_0 \) the inductive (loop) electric field. Finally, the full Ohm’s law equation (3.3) yields the velocity field, which turns out to be slaved to the magnetic field. In the plasma core, the contribution from the applied electric field \( E_0 \) is larger than that from the mean parallel current density, but it is smaller in the edge, as shown in figure 13. The difference is balanced by the electrostatic term, which provides an anti-dynamo contribution in the core and a dynamo contribution in the edge. The same interpretation can be applied to the MH case, where the electrostatic field is due to charge separation patterns produced by the periodic relaxation events and is again the main contributor to the fluctuating velocity field [218, 219, 221].

The pinch velocity leads to a build-up of the plasma density on the helical axis in MHD simulations where density is free to evolve [193]. This does not occur in a genuine plasma, where the central density profile is flat. This inconsistency was avoided in [222] by adding a diffusive term in the continuity equation for the density.

It is worth noting that a dynamo conversion of poloidal flux to toroidal flux, analogous to the RFP one, is invoked to explain the nonlinear saturation of the \( m = 1, n = 1 \) mode, i.e., the helical core observed in the hybrid mode of operation of the tokamak [143, 223]. This type of dynamo is also present in a flux rope configuration susceptible to the kink instability [224].

3.9. Further theoretical aspects

3.9.1. Toroidal effects. When performed in toroidal geometry, 3D nonlinear visco-resistive MHD simulations show that toroidal coupling prevents the system from reaching a pure SH state when dissipation increases, but that magnetic chaos due to toroidal coupling stays limited close to SH states for the aspect ratios of the largest present RFP’s [225]. This is confirmed by toroidal simulations with the PIXIE3D code, also showing an \( m = 0 \) island chain induced by the toroidal coupling at the \( q = 0 \) reversal surface [124]. Incompressible 3D nonlinear visco-resistive MHD simulations show that, in contrast to the cylindrical case, the toroidal one presents a double poloidal recirculation cell with a shear localized at the plasma edge [226].

3.9.2. Two fluids effects on single helicity states. Analytical calculations in [220] show that (i) the SH state is the same on using single or two-fluid Ohm’s law, which backs up a numerical result (section VA of [227]); (ii) the SH mode amplitude is insensitive to \( S \), which backs up experimental [104] and numerical results in a two-fluid context [227]. Adding gyroviscosity in the force balance equation enables to get the experimental ratio of secondary to primary modes, while it is twice too large otherwise [227]. More analytical results can be found in the course [228].

3.9.3. Role of anisotropic resistivity. QSH states were obtained in simulations performed with an anisotropic thermal conductivity and using a multiple-time-scale analysis. In these simulations, the temperature distribution indicates the existence of closed magnetic surfaces, and there is a hot confined region [229]. When resistivity increases steeply at the edge, SH states are obtained for parameters similar to the case without thermal transport [230, 231].

3.9.4. Aspect ratio scaling of single helicity. The aspect ratio scaling law for the toroidal mode number of experimental QSH states was shown to be close to the one corresponding to the optimal electromagnetic response of the toroidal shell surrounding the plasma [232]. In [233] it is shown that the observed scaling with the aspect ratio and reversal parameter for the dominant mode in the single helical states can be obtained by minimizing the distance of the relaxed state described in [234] from a state which is constructed as a two region generalization of Taylor relaxation model [235].

3.9.5. Shear stabilization of \( m = 1 \) couplings in multiple helicity. Recently, a new MHD approach was proposed to explain why high current is favourable to QSH-SHAx states, why the innermost resonant \( m = 1 \) mode is the spontaneous dominant mode and why there are crashes of this mode. It invokes the shear stabilization associated with the 3D structure of the dominant mode to interrupt the nonlinear mode–mode coupling occurring in the MH regime [236]. This kind of ideas had been first introduced in [237]. The crashes can also be interpreted as the consequence of pressure-driven resistive modes, which introduce a feedback between transport and the MHD stability of the system [238].

4. Transport and confinement in the multiple helicity and quasi single helicity RFP

Energy and particle transport in RFP plasmas are strongly influenced by the presence and magnitude of tearing instabilities. Multiple magnetic islands can overlap and cause the magnetic field to be stochastic in a large volume of the plasma. Not surprisingly, the role of stochastic transport has been thoroughly examined in RFP plasmas. The well-known test-particle transport model for parallel streaming in a stochastic field works well, apart from nuances. However, QSH states and the application of current profile control both have reduced tearing instability, allowing confinement that is not limited by stochastic transport. Microinstability associated with several branches of drift waves plays an important role when tearing is suppressed.

In QSH states, the healing of magnetic surfaces allows hot regions to form inside the helical island structure. A transport barrier forms at the boundary of the helical core, and thermal transport is locally reduced. Microinstability from ion temperature gradient (ITG) and/or microtearing (MT) is dominant in the transport barrier. Trapped electron modes (TEM) are the dominant microinstabilities with current profile control, as discussed in chapter 5.
Impulsive DREs (aka sawteeth) that appear in MH plasmas as well in QSH plasmas during back-transitions to MH impart strong effects on transport and heating. The amplitudes of magnetic fluctuations burst in these events, which cools electrons, in addition to regulating magnetic flux generation and dynamo sustainment. A significant fraction (10%–30%) of the stored magnetic energy is released abruptly in a sawtooth crash, but much of this released energy remains confined as ion thermal energy through a powerful non-collisional heating process. Ion heating occurs simultaneously with electron cooling, and the sustained ion temperature is a large fraction of the electron temperature, even when collisional coupling of the species is weak.

Neutral beam injection (NBI) experiments have allowed assessment of energetic (or fast) ion confinement and stability. Importantly, energetic ion confinement in the RFP is classical, even in the MH regime with strong magnetic stochasticity. This is related to the drift orbits of fast ions. The energetic ion confinement is also good in QSH states, provided secondary modes remain sufficiently small. The fast ion beta in the core of MST plasmas reaches 8%, by which point energetic particle and Alfvén eigenmodes are observed that limit the fast ion population. The relatively weak but strongly sheared magnetic field in the RFP generates a distinct Alfvén continuum and provides a complementary environment to tokamak and stellarator plasmas for the study of fast ion driven instabilities.

In the following, core heat and particle transport properties are described in sections 4.1 and 4.2, both in MH and QSH states. An overview of the studies of the effect of microinstabilities on transport in RFPs is also given in section 4.3; ion heating and the effect on transport during DREs (sawtooth cycle) are discussed in section 4.4. The main features of transport in the edge region are reported in section 4.5. Section 4.6 gives a short summary of the isotopic effect in the RFP. Finally, sections 4.7 and 4.8 summarise the observations related to fast particle confinement in the RFP.

4.1. Core heat transport and confinement

In standard MH plasmas, strong temperature and density gradients are observed nearby the \( q = 0 \) toroidal field reversal surface [239–241], which naturally distinguishes two domains in describing transport: a core region inside the reversal surface, and an edge region from the reversal surface to the wall/limiter.

In the edge, transport is strongly influenced by a chain of \( m = 0 \) magnetic islands developing at the reversal surface [242, 243], and forming in fact a transport barrier corresponding to the \( q = 0 \) surface. The local pressure gradient is large, and the magnetic stochasticity is small compared to the core. Microinstabilities also play a role in the edge.

In the core, stochastic transport is dominant when magnetic islands created by multiple tearing modes are overlapped. The degree of magnetic stochasticity is quantified by the Chirikov stochastic instability parameter \( s \), defined as [244]

\[
s = \frac{1}{2} \frac{|W_{mn} + W_{n'm'}|}{|r_{mn} - r_{n'm'}|} \tag{4.1}
\]

where \( m \) and \( m' \) (\( n \) and \( n' \)) are the poloidal (toroidal) mode numbers, \( W_{mn} \) is the width of the magnetic island related to the \((n, m)\) mode, and \( r_{mn} \) the radial location of the resonant surface. Overlap of neighbouring magnetic islands occurs when \( s \gg 1 \), and the field becomes fully stochastic when \( s \gg 1 \).

In a domain where \( s \gg 1 \), transport caused by parallel streaming in a stochastic magnetic field can be large, as proposed by Rechester and Rosenbluth (RR) [3]. According to RR, the thermal diffusivity \( \chi_{RR} \) depends on the magnetic fluctuations as:

\[
\chi_{RR} = v_0b_mD_m = v_0\left(\frac{b_t}{B}\right)^2L_{ac} \tag{4.2}
\]

where \( D_m \) is the stochastic magnetic field diffusion coefficient, \( b_t \) is the magnetic fluctuation amplitude radial to unperturbed magnetic surfaces, \( B \) is the equilibrium field, \( v_0 \) is the particle thermal speed, and \( L_{ac} \) is an auto correlation length characterizing the random diffusion of magnetic field lines. Stochastic transport is not inherently ambipolar, and the particle flux is given by [245]

\[
\Gamma = -D\nabla n + v \cdot n \tag{4.3}
\]

where \( D \) is the stochastic particle diffusion coefficient, and \( v \) is the convective velocity. Assuming thermal energy loss is dominated by electrons due to their high mobility (\( T_e \approx T_i \)), the ambipolarity constraint leads to \( \chi_{RR}/D = (m_i/m_e)^{1/2} \).

Since stochastic heat transport scales as \( \chi_{RR} \sim \left(\frac{b_t}{B}\right)^2L_{ac} \), it is small if \( b_t \) is small, or if the field is weakly disordered (small effective \( L_{ac} \)). The region near the \( q = 0 \) surface has large magnetic shear and densely packed \( m = 1 \) resonant surfaces. This encourages island overlap, but if the fluctuations are small, stochastic diffusion is small. As an example, the minimum heat diffusivity in high current pulsed poloidal current drive (PPCD) plasmas \( \chi_e \sim 2 \text{ m}^2/\text{s} \).
The measured equation (4.2) (crosses). The shaded region represents uncertainty in line tracing (diamonds) and with the analytic RR formula in MST (solid line), compared with stochastic diffusion from field from [239], Copyright 2003 by the American Physical Society. Adapted figure with permission.

The broad spectrum of $m = 1$ and $m = 0$ tearing modes in the MH regime allows for the appearance of many magnetic islands on resonant surfaces that span the minor radius of the plasma. As an example, the top panel of figure 21 shows the safety factor profile, $q(r)$, for a typical MH plasma in MST, with the widths of the largest $m = 1$ magnetic islands indicated by horizontal lines, each centred on its resonant surface, $r_{\text{res}}$. The bottom panel shows the electron temperature profile measured at a time between discrete relaxation events, which is flat in the core. The island widths are estimated using the amplitudes of the radial magnetic fluctuation, $b_{\text{rms}}(r_{\text{res}})$, calculated by the nonlinear 3D MHD code DEBS [174] but scaled to match the fluctuation amplitudes measured at the plasma surface. The amplitudes in the simulations are larger than in the experiment by a factor $\sim 1.7$ [239].

The electron heat diffusivity, $\chi_e(r)$, in the plasmas represented by figure 21 was measured using global power balance [239], shown as the solid line in figure 22(a). Two comparisons with the predicted stochastic heat diffusivity, $\chi_{\text{RR}}(r)$, are overlaid. Analytic values of $\chi_{\text{RR}}$ using equation (4.2) based on $b_{\text{rms}}(r_{\text{res}})$ and fixed correlation length, $L_{ac}$, are shown as crosses at the resonant surfaces, $r_{\text{res}}$, for the largest $m = 1$ modes. These analytic values of $\chi_{\text{RR}}$ agree reasonably well with $\chi_e$ near midradius, $r/a \gtrsim 0.6$, but they significantly overestimate the heat diffusivity near $r = 0$. The points plotted as diamonds are for a more native calculation of $\chi_{\text{RR}} = v_e, a D_m$ based on the magnetic diffusion, $D_m = \Delta r^2/2 \Delta L$, determined from an ensemble of direct field line tracings in the 3D field provided by the DEBS simulations noted above. Magnetic diffusivity, $D_m$, measures the average radial excursion, $\Delta r$, of field line trajectories over length, $\Delta L$, along the field. This direct evaluation of $\chi_{\text{RR}}$ agrees well with the measured $\chi_e(r)$. The analytic RR formula assumes the magnetic field is fully stochastic. This requires large values for the Chirikov parameter, $s$, which are shown in figure 22(b) for the data corresponding to figure 21. Evidently, $s > 5$ is required for an accurate estimate using equation (4.2), which is why the analytic values fail for $r/a \lesssim 0.6$ where $s \approx 1$. These comparisons highlight that it is local island overlap that regulate magnetic diffusion. As another example of this, since the magnetic fluctuation amplitudes of modes resonant in the edge region of MST plasmas are small, stochastic transport vanishes in the outer region near the $q = 0$ surface [249], despite a size-able MP from modes resonant deep in the core that can advect field lines but do not cause them to become stochastic.

The radial profile of the electron thermal diffusivity $\chi_e$ has been analysed in various experiments by expressing $\chi_e \propto (b/B)^{\alpha}$, as inspired by the RR scaling: while a dependence on magnetic fluctuations has been confirmed in all devices, the $\alpha$ exponent has not always been found to be 2 as in equation (4.2). In RFX and RFX-mod, where electron and ion temperatures are close [85], steady state analyses of the electron thermal transport adopted a single fluid approach. An effective thermal diffusivity $\chi_{\text{eff}}$ has been calculated, whose dependence on the magnetic fluctuations has been found to be $\chi_{\text{eff}} \propto (b/B)^{1.4}$ [63, 64]. This result has been interpreted by a model taking into account the effect of mode spectrum on field line stochasticity [250].

A perturbative approach was applied to evaluate the thermal diffusivity in the TPE-RX device, based on the analysis of the radial propagation of a cold pulse during pellet injection [251]. In the scaling with magnetic fluctuations $\alpha \sim 1$.

Nonlinear MHD computations of the electron energy transport in a stochastic field were also applied to characterize energy confinement in MH RFPs. When using the classical parallel electron thermal diffusivity to estimate the parallel heat diffusion, the experimental energy confinement scaling was well reproduced [252–254].

Analysis of ion energy transport is complicated by the fact that ion heating from magnetic reconnection and turbulence is very powerful in RFP plasmas. In both MST and RFX plasmas, the ion temperature is a large fraction of the electron temperature. In MST the ion temperature can even exceed the electron temperature during sawtooth relaxation events, which is not expected for collisional relaxation of an ohmically heated plasma [255]. Non-collisional ion heating is discussed further in section 4.4.
4.1.2. Quasi single helicity states. To highlight the effect of the magnetic equilibrium, figure 23 compares the electron temperature profiles measured by the Thomson scattering diagnostic in MH, QSH-DAx and QSH-SHArx in an RFX-mod plasma. Data refer to three times during the same plasma discharge. As described in section 2.3, the QSH-DAx state (black squares) corresponds to the development of a hot island involving a small fraction of plasma radius. When the magnetic separatrix is expelled and the transition to a QSH-SHArx state occurs, the hot island becomes wider and encompassed by strong electron temperature gradients (red diamonds), establishing an electron thermal transport barrier (eITB) [102, 108, 109, 256, 257]. In agreement with the theoretical description [112], a threshold has been found for the occurrence of the transition QSH-DAx/QSH-SHArx, which is observed when the dominant mode amplitude, normalized to the edge magnetic field, is larger than about 3%–4%. The amplitude of secondary modes is another important parameter that determines the appearance of wide helical structures. In particular, the widest thermal structures are found at the lowest amplitudes of the two innermost secondary modes resonant according to the helical safety factor: i.e. the $m = 1, n = -9$ in RFX-mod, as can be seen in figure 24 [258].

In RFX-mod QSH-DAx states the heat diffusivity in the temperature gradient region ranges between 6 and 35 m$^2$ s$^{-1}$, approaching the tokamak range. Despite the small size of the island, the global electron confinement time $\tau_E$ is higher than in MH states (by about a factor 2), about 1.3 ms [259]. Simulations of test ion and electron transport show the average diffusion coefficients inside the helical core to be about one order of magnitude lower than those found in MH plasmas [260]. When QSH-DAx states are stimulated by a helical edge magnetic field (see section 2.2.3), the magnetic chaos decreases too, involving the region outside the island, and allowing for a global enhancement of confinement by about 50% [261]. Indications of improved electron energy confinement have been found in TPE-RX [262].

The thermal diffusivity inside a magnetic island $m = 1, n = 6$ close to the axis (QSH-DAx) has been measured in MST, showing flat temperature profiles with rapid parallel heat conduction along helical magnetic field lines in-between the reconnection events. Instead, just after a DRE, a temperature gradient develops inside the $m = 1, n = 5$ island, when the $m = 1, n = 5$ mode may briefly come into resonance near the magnetic axis. This suggests local heating and relatively good confinement within the island [263].

Regarding the confinement in QSH-SHArx states, in MST the global energy confinement improves by about 50%. A three-fold improvement in confinement is obtained by forcing a slow decay of the plasma current [105]. In RFX-mod, where extensive experiments and analysis involving QSH-SHArx states have been carried out, the eITBs encompassing the wide central island are characterized by an electronic thermal diffusivity lower than 20 m$^2$ s$^{-1}$, with minimum values as low as 2 m$^2$ s$^{-1}$. Such a value of $\chi_e$, though still higher than the classical one (under 1 m$^2$ s$^{-1}$), is much lower than outside the barrier or in MH states, where it ranges around 100 m$^2$ s$^{-1}$ [264–266]. Both thermal diffusivity and electron temperature gradient length scale with the total amplitude of the secondary $m = 1$ modes [257, 267]. So does the inverse of the maximum electron temperature gradient [103, 267]. These scalings indicate a strong link between the level of magnetic chaos and the strength of the barrier.

The DKES/PENTA codes [268, 269], developed for stellarators, have been applied to the RFP for local neoclassical computations of the thermal diffusion coefficients with electron density and temperature profiles typical of an RFX-mod QSH-SHArx plasma. A comparison with power balance estimates shows that residual chaos due to secondary tearing modes and small-scale turbulence still contributes to drive anomalous transport in the barrier region [137].
It has to be mentioned, though, that in helical states the electron temperature gradients show a time behaviour not necessarily related to the dynamics of secondary modes [265]. The highest gradients (above keV m$^{-1}$) occur in the rising phase of the QSH state, and start decreasing (0.5–1 keV m$^{-1}$) before the destabilization of secondary modes [103]. Despite this intermittent behaviour of eITBs, there are examples of long-lasting eITBs, up to 18 ms in deuterium plasmas, in particular with QSH stimulated by 3D MPs externally applied [121].

In RFX-mod, in QSH-SHAx states the highest $\tau_{Ee}$'s values are obtained. This is shown in figure 25, where a clear relation between the electron energy confinement time and the secondary mode amplitude appears [103, 104, 270–272]. This is also found in a three parameter fit of the energy confinement time with plasma current, average density and amplitude of secondary modes $b_r$, which yields a scaling $\tau_{Ee} \approx I^{0.81 \pm 0.02} \langle n_e \rangle^{0.28 \pm 0.02} b_r^{-0.88 \pm 0.01}$ (figure 26) [273]. This result strongly motivates the upgrade of the RFX-mod magnetic front-end presently under implementation [20], whose main aim is to obtain a decrease of $b_r$ by moving the plasma edge closer to the thin shell.

As yet, spontaneous helical states are observed at relatively low plasma densities, $n/n_G \leq 0.35$ [274]. Such a limit has been related to the effect of an anomalous viscosity due to micro-turbulence which can act outside the reversal region [211, 271, 275].

To interpret the flat temperature profile of the hot region encompassed by the temperature gradients in QSH-SHAx states, transport related to self-consistently generated vortical drift motions due to electrostatic turbulence was proposed in [276].

### 4.2. Particle transport and confinement

In the RFX device, an interpretative analysis of the electron density profiles measured by the 13 chord two colour interferometer [277] in MH plasmas yields a particle diffusion coefficient of $\approx 10^{-20}$ m$^2$ s$^{-1}$ in the core decreasing to about $5$ m$^2$ s$^{-1}$ towards the edge. The same analysis also indicates the presence of a particle outward pinch, necessary to reconstruct the hollow or flat experimental profiles [278]. The ratio of thermal to particle diffusivities turns out to be consistent with RR theory in the core. A subsequent analysis carried out in RFX-mod both in MH and QSH plasmas also yields $\chi/D = (m_i/m_e)^{1/2} n/n_G^{1/2} b_r^{-0.88 \pm 0.01}$ and a dependence $D^* \propto (b/B)^{1.5}$ (with $D^*$ core diffusivity normalized to the thermal velocity) [279]. This implies that, though residual magnetic chaos remains the main drive for particle transport also in QSH plasmas, the diffusion coefficient is lower, as expected, due to its dependence on magnetic fluctuations: figure 27 shows that the same power law is able to describe the behaviour of $D^*$ in both MH and SHAx states, with $\alpha = 1.5$ in equation (4.2).

The fact that in RFX and RFX-mod the scaling of particle diffusivity is weaker than in RR theory has been related to a sub-diffusive nature of transport. In [280, 281] the ORBIT code [130] is applied, showing that the velocity in equation (4.3) can be interpreted as a sub-diffusive correction to the diffusive evaluation of fluxes. Indeed, in RFX-mod the magnetic field is chaotic, but not much above the stochastic
threshold, so that hidden structures are anyway present and the RR assumption of random phase approximation and fully isotropic stochasticity might be questionable.

Simulations in [282] predict a strong improvement of particle confinement in QSH with respect to MH. This has been experimentally confirmed in several devices: in TPE-RX, the transition from MH to QSH came with an enhancement of the particle confinement time $\tau_p$ by about 30% [283]. The improvement was confirmed in RFX-mod, in particular by using pellet injection [26]. A pellet entering the central hot region in a QSH-SHAx plasma induces a higher density with a possibly peaked profile, which triggers both an enhancement of $\tau_{\text{Be}}$ by a factor 2 to 3, and an increase of $\tau_p$, with a maximum value of 12 ms obtained at 1.5 MA [26].

Single particle trajectory numerical studies applied to RFX-mod plasmas show that, while passing particles are confined for a very long time, particle trapping dominates transport across the helical structure at the rather low collisionality of QSH regimes. The neoclassical diffusion coefficient is one order of magnitude larger than the classical one [260, 284, 285]. When trapped particles drift out of the helical core, at $r/a \approx 0.6$, they become almost passing without being lost, because the helical distortion decreases when going outwards. This results in a very low population of superbanana particles at the experimental levels of helical magnetic field (about 10% of the total) [286]. Therefore, transport is proportional to the collision frequency, in contrast with unoptimized stellarators where the lower bound of transport scales inversely with this frequency [286]. The diffusion coefficients computed by this approach were confirmed by local neoclassical transport calculations including the radial electric field.

The computation of single particle trajectories exhibits a different influence of helical states on main gas and impurity transport [260]. Several studies of impurity transport have been carried out, both in MST and in RFX and RFX-mod (see for example [287–290]). In MH regimes the experimental impurity diffusion coefficient is around 10–20 m$^2$ s$^{-1}$, and the convective term in equation (4.3) is outward, in both experiments. An example for RFX-mod is given in figure 28. These values are larger than those predicted by the stochastic model (also shown in the figure). The analysis of impurity behaviour in RFX-mod QSH states shows that inside the core region the confinement of impurities does not increase, which makes unlikely their central accumulation [272]. A core impurity diffusivity one order of magnitude higher than at the edge has been estimated [291], with an outwardly directed pinch velocity over the whole plasma featuring a strong maximum around the core/edge transition point of the diffusivity, much wider than in MH.

4.3. Effect of microinstabilities

Gyrokinetic numerical simulations and analytical analyses have been applied to study the role of microturbulence as a driving mechanism of transport in the RFP. In axisymmetric configurations, the ITG mode threshold is larger than in the Tokamak by a factor $R/a$, due to a strong Landau damping related to short field connection lengths of plasmas acting in low-$q$ configurations [292–294]. In addition, non-linear calculations show that in axisymmetric plasmas the impact of residual zonal flow, beneficial in controlling ITG, can be strong [295]. The threshold for ITG stability decreases in helical states, because of the enhancement of temperature gradients in the outer part of the helical deformation where magnetic surfaces are close-packed [295]. Therefore, ITG turbulence might be an important contributor to the total heat transport in helical states, also in combination with a decreased impact of the stochastic transport related to magnetic chaos. The stability of ITG modes might be further decreased in the presence of strongly outwardly peaked impurity profiles [296]. However, quasilinear and nonlinear three-species simulations of ITG turbulence prove that the inward impurity flux corresponding to a strong ITG would not be compatible with the measured positive impurity peaking [297].

Linear gyrokinetic simulations of RFX-mod QSH-SHAx states show that microtearing modes are unstable at the eITB, as shown in figure 29. For the strongest eITB’s, quasi-linear estimates of the associated transport show MT-driven thermal transport to be comparable to the experimental one [298]. It has to be noted that MT modes have actually been experimentally observed in RFX-mod helical states, with amplitude well correlated to the electron temperature gradient [299].
Microturbulence has been modelled for MST plasmas with a focus primarily on plasmas with inductive current profile control (PPCD) [300–302]. The ITG and MT channels have been analysed, and the role of finite $\beta$, already well known in the tokamak [303], has been elucidated for the RFP, with ITG unstable at low $\beta$ (vanishing around $\beta = 6\%–10\%$) and MT unstable at high $\beta$. For PPCD plasmas, density-gradient-driven trapped-electron modes (TEM) are the dominant microinstability, shown to be in good agreement with experimental measurements. This is discussed in chapter 6.

The occurrence of collisionless MT instabilities has been studied in RFX-mod [304] and MST [301], indicating that the high magnetic drifts in the RFP could favour such microinstabilities.

4.4. Non-collisional ion heating and effect of discrete reconnection events

As described in chapter 2, DREs occur in both MH and in QSH plasmas, and in the latter they correspond to a back transition to the MH state. In MST plasmas, the stored magnetic energy released during a DRE is 10%–30%, with an instantaneous power that can exceed 10 MW. Recently, the same kind of evaluation has been made for RFX-mod back transitions QSH/MH, where it is estimated that 30%–40% of the stored magnetic energy is released during reconnection and does not contribute to ohmic plasma heating [114]. A large fraction of the magnetic energy released during reconnection events is converted into ion thermal energy, heating the ions to high temperature 1–3 keV [255, 305–307]. An example of the time behaviour of ion temperature during a DRE in MST is given in figure 30. Detailed studies of ion heating in MST plasmas reveal that it is charge and mass dependent [308, 309] anisotropic in space with $T_{i,\perp} > T_{i,\parallel}$ [310] and an energetic ion tail forms spontaneously with a power-law distribution [310].

Several theories have been proposed to explain the ion heating and acceleration mechanisms. The leading candidates are ion cyclotron damping in a turbulent cascade [311, 312], and stochastic heating [308]. Both mechanisms tend to heat anisotropically with $T_{i,\perp} > T_{i,\parallel}$, a key heating characteristic. The turbulent cascade, driven by large-scale tearing, exhibits non-collisional dissipative features suggestive of a kinetic process consistent with the observed ion heating [313, 314]. While the drive is MHD tearing, the cascade exhibits drift waves and is not purely Alfvénic, either through nonlinear coupling or independent drift instability [315]. The RFP’s turbulent cascade, coupled to strong ion heating, represents a grand challenge for multiscale modelling that spans self-organization on the global scale to non-collisional dissipation at the microscale. Viscous heating has also been evaluated, but it is more difficult to do so in a consistent manner [316]. The energetic ion tail is formed, in part, due to runaway ion acceleration [317, 318]. A large inductive electric field accompanies the flux change driven by reconnection, which exceeds the Dreicer electric field for ions.

Temperature profiles as measured by multi-chord SXR tomography in MST tend to peak before and flatten after the crash [41, 319]. In the crash phase an outward propagating heat pulse is detected, allowing to quantify the presence of stochasticity due to the multiple $m = 1$ modes [320]. In order to investigate the effect of DREs on electron thermal transport, a time dependent model has been applied to RFX-mod and EXTRAP-T2R MH plasmas [321, 322]. The $\chi_e$ radial profile is divided into three regions: (a) the core, where $\chi_e = k_1(\hat{b}/B)^{1.4}$ is assumed; (b) the reversal region, with $\chi_e = k_0$; (c) the edge region, where a much higher diffusivity is assumed. Figure 31 shows an example of the $\chi_e$ profile reconstructed before and after a DRE [321], with enhanced core thermal diffusivity just after the DRE. The same model also gives good agreement with the experiments when applied to the EXTRAP-T2R MH plasmas [321]. Experimentally, as...
reported in [51], enhanced non axisymmetric particle losses are observed in RFX-mod during DRE localized at the toroidal angle where MHD tearing modes lock in phase and to the wall. The analysis by Biewer described in section 4.1.1 was performed for a time point midway between reconnection events. Later the analysis was applied to the full evolution of the sawtooth cycle [323], summarised in figure 32. This included a large ensemble of DEBS simulations that matched experimental parameters as closely as possible. The best agreement between experiment and modelling includes a neoclassical enhancement (∼2×) of the parallel electrical resistivity associated with trapped electrons. The inboard-outboard mirror in a toroidal RFP equilibrium is comparable to that for tokamak plasmas of similar aspect ratio [62], but drifts are mostly in-surface for the RFP, and perpendicular neoclassical transport enhancement is negligible [324, 325]. Since trapped electrons are not free to travel long distance along field lines, the stochastic heat conductivity is reduced, $\chi_e = f_c v_{th,e} D_m$, where $f_c$ is the fraction of passing particles, and $D_m$ is the previous stochastic magnetic diffusivity. This reduction in heat transport is likely important in astrophysical plasmas that contain turbulent magnetic field.

The thermal diffusivity inside a magnetic island close to the axis has been measured in MST, showing flat temperature profiles with rapid parallel heat conduction along helical magnetic field lines in-between the DREs. Instead, just after a DRE, a temperature gradient develops inside the island, suggesting reduced thermal transport [263].

Magnetic fluctuation induced particle transport and density relaxation in the core of MST plasmas was measured using advanced interferometry-polarimetry [326]; the particle flux is non-ambipolar in the DRE, which implies the formation of zonal flow [327, 328].

4.5. Edge transport and turbulence

The $m = 0$ island chain which develops at the field reversal surface strongly influences the edge transport, both in the MH and QSH state. Tearing modes, which are certainly detrimental for confinement when producing intense magnetic chaos in the $q > 0$ region, have also a secondary positive effect, allowing for the formation of good magnetic surfaces close to $m = 0$ islands, which indeed act as a transport barrier and have a direct impact on plasma pressure profile [242, 281]. An example is given in figure 33.

The change in the pressure profile, whose gradient increases when the magnetic shift is outward [329], influences the edge turbulence. In particular in RFX-mod it is found that the radial correlation length of the fluctuations increases with the characteristic radial pressure length, indicating that the latter drives the radial dimension of turbulent structures [329].

In fact outside the reversal surface transport is mainly driven by electrostatic turbulence and, similar to tokamak and stellarator, is characterized by a highly sheared $E \times B$ flow [30, 31, 249, 330–332]. The heat transport driven by electrostatic fluctuations was measured in the REPUTE experiment in the region $a/2 < r < a$ [333] and in the edge region of MST [334]. In both devices, it was found that the energy transport driven by electrostatic fluctuations was relatively small with respect to the total heat loss, though in MST the particle flux were found to be comparable to the total particle loss. In particular, in [333] an evaluation of the particle and energy flux driven by stochasticity was given, which was subsequently interpreted by a model calculating the ion and electron fluxes under the ambipolar constraint [335]. Experiments on the EXTRAP-T2R device showed that Reynolds stress, driving the $E \times B$ flow against anomalous resistivity, features a high gradient in the edge region where the flow shear is large, and is mainly related to the electrostatic component, despite the high magnetic fluctuations in this region [336, 337]. Coherent structures emerge as bursts in the turbulent background, with rotation direction determined by the local velocity shear: in RFX...
and EXTRAP-T2R they have been identified as dipolar and monopolar vortices, contributing by about 50% to the radial diffusivity [338]. The statistical properties of such turbulence share common features with other magnetic configurations [339].

The edge electromagnetic turbulence measured in RFX-mod when operated as an RFP and as a tokamak has been compared [340, 341]. Electrostatic fluctuations in both configurations reveal a strong degree of intermittency with very similar scaling of the shape of the PDF of fluctuation increment, while magnetic fluctuations behave differently, with a lower degree of intermittency in the tokamak case, consistent with the lower $\beta$. In particular, the study of turbulence in the scrape off layer of RFX-mod when operated as a tokamak allowed the validation of the GIBS code [342, 343].

The edge electrostatic properties are strongly related to the underlying magnetic topology. When 3D non-axisymmetric MPs are applied in RFX-mod to stimulate QSH states (section 2.2.3), though they have a main $m = 1$ periodicity, at the edge the flux induced by electrostatic turbulence results to be modulated by an $m = 0$ magnetic island, with an enhancement close to the O-point and a reduction at the X-point; transport is due to fluctuations propagating in the electron diamagnetic drift direction [344]. Also, the connection length to the wall of the edge magnetic field $L_{cw}$ is the result of the interaction between the applied dominant $m = 1$ and the $m = 0$ modes and has a complex structure, responsible for an impure $m = 1$ behaviour observed for floating potential, electron pressure, perpendicular flow, high frequency magnetic turbulence [210]. The edge flow pattern in presence of an MP has been reconstructed and shows a convective cell structure, with the toroidal flow moving to an outer region when the deformation related to the MP ($\Delta_{\text{MP}}$) is positive (outwards) and to a more internal one when $\Delta_{\text{MP}} < 0$ [345].

As in the core, at the edge the reduction of tearing mode amplitude is related to an improved confinement: the temperature gradient, which scales almost linearly with plasma current, increases at low secondary mode amplitude, and so does the thermal diffusivity, as shown in figure 34 [243].

Recent research in RFX-mod [346] points out the crucial role of the secondary modes in influencing the plasma performance in QSH (electron temperature and confinement time, particle influx and PWI, radiated power). In QSH plasmas with a well-formed helical state, despite secondary mode amplitude is kept low by the feedback system, a deformation related to a non-stationary locked mode (LM) remains present, similar to MH plasmas [78], though with a strongly reduced size of the displacement ($<1$ cm). In figure 35 such an LM is observed as an interference pattern (panel (b), black curve), which corresponds to two dips of the magnetic field connection length to the wall, $L_{cw}$ (see panel (a)). These ‘holes’ can be seen as homoclinic lobes (‘fingers’) at two toroidal positions ($\phi \approx 110^\circ$ and $\phi \approx 133^\circ$) [346]. The analysis of RFX-mod data suggests a threshold of $\approx 0.3$ in terms of the ratio between secondary locking size and displacement of the helical structure, $S_e = \Delta_{\text{sec}}/\Delta_{1,7}$ to completely avoid the detrimental effect of the secondary mode locking, which is important in perspective of the RFX-mod upgrade, where a reduction by a factor $\sim 2$ in secondary mode amplitude is expected [347].

### 4.6. Isotopic effect

The comparison of plasmas with hydrogen and deuterium as filling gas carried out in RFX-mod has shown a positive isotopic effect on the confinement of SHAx states similar to the one in tokamaks. For plasma currents around 1.5 MA, the energy confinement time scales as $m_0^{0.3}$; the particle influx is significantly reduced and the particle confinement time scales as $m_0^{0.45}$. The effect on confinement is mainly due to
4.7 Energetic ion confinement and stability

The confinement of energetic ions is fundamental to fusion energy. A tangential NBI system on MST has allowed detailed investigation of energetic (aka fast) ion confinement and stability in both MH and QSH plasmas. The NBI sources 25 keV, H or D neutrals (varied mixture) with an effective current ≤40 A and power ≤1 MW. The NBI pulse is ≤20 ms. The energetic ion confinement is measured using the ‘beam blip’ technique, in which a population of fast ions are created in a short NBI pulse, and the decay in d–d fusion neutron emission is observed using scintillation detectors [349]. An advanced neutral particle analyser measures the fast deuterium and hydrogen distributions simultaneously, which reveals dynamics associated with energetic particle driven instabilities [350].

The energetic ion confinement time is consistent with classical slowing down (~20 ms) and an order of magnitude larger than the thermal particle confinement time (~1 ms) [349, 351]. Drifts of the fast ion guiding centres carry them off the magnetic field such that they do not experience the magnetic field’s stochasticity, even in the MH regime. The drift is quantified using a generalization of the safety factor based on the fast ion guiding centre velocity [349]. The fast ion safety factor is larger than the magnetic safety factor, thus avoiding resonance with the dominant core tearing mode. Interestingly, the fast ion safety factor helps explain stochastic diffusion of fast ions in tokamak plasmas when an m = 1, n = 2 neo-classical tearing mode is present [352, 353].

For longer duration NBI, classical heating [354] and current drive [355] are observed until the fast ion population reaches a critical threshold at which point both energetic particle modes (EPM) [356] and magnetic-island-induced Alfvén eigenmodes (MIAE) [357, 358] are triggered by saturation of a wave-particle resonance condition [359]. Example behaviour is illustrated in figure 37: the NBI current is increased in steps and onset of EPM activity occurs when the NBI current is >20 A (0.5 MW). The fast ion distribution is modelled using the NUBEAM module in the TRANSP code [354, 360]. The curve labelled ‘TRANSP’ in figure 37(b) is the expected d–d neutron emission for maximum NBI without loss of fast ions. Density and magnetic field fluctuations measured by an interferometer-polarimeter show that the energetic particle driven modes are nonlinearly coupled and lead to flattening of the fast ion profile during ‘fishbone’ type bursts, resembling predator–prey dynamics [361]; toroidal Alfvén eigenmodes are possible in an RFP plasma, but the beam energy for MST is too low to access the bandgap.

The fast ion profile has been measured using a columnated neutron detector [362]. The corresponding fast ion beta profile is shown in figure 38, with βₐ(0) ~ 8%. The onset of EPM activity occurs for a critical gradient in the fast ion beta. The profile in figure 38 is the quasi-steady state profile at marginal stability and full beam power. The limit to the fast ion beta is believed to result from the onset of the EPM and MIAE activity, which redistributes the fast ions outward in radius where they are subject to stochastic transport [351, 362]. Central βₐ ~ 8% is comparable to the alpha beta expected in a reactor plasma. Also, the 25 keV fast ions in MST (B < 0.5 T) have a normalized gyroradius, 𝑗β/𝑎 ~ 0.1, similar to alpha particles in a reactor plasma with B ~ 5 T, so the dynamics observed in MST plasmas are likely representative of expectations for alpha particles.

The fast ions have a stabilizing influence on the tearing mode resonant closest to the magnetic axis where the fast ion density is large [357, 363]. This delays the transition from MH...
Figure 38. The quasi-stationary fast ion beta profile at marginal stability, measured at $t = 30$ ms for conditions as in figure 26 with full beam power (40 A) in 300 kA MST plasmas. Reproduced courtesy of IAEA. Figure from [362]. Copyright 2019 IAEA.

Figure 39. Central electron temperature as a function of plasma current in RFX-mod. Red points refer to deuterium plasmas. Reprinted from [376], Copyright 2018, with permission from Elsevier.

to QSH, and the fast ion confinement is degraded if the secondary mode amplitudes are not reduced [364]. With QSH, the fast ions are born over a larger region in the core as a result of the helical distortion, making them more susceptible to stochastic transport in the mid-radius region. However, if the secondary modes are reduced, which is a key feature of the best performing QSH states, then the fast ion confinement recovers its classical behaviour [351, 365].

4.8. Fast electrons and the kinetic dynamo theory (KDT)

In many RFP experiments, fast electrons have been measured at the plasma edge with parallel energies characteristic of the central temperature [366–370]. In some cases, these electrons carry most of the edge equilibrium current.

The KDT [59] predicts that these electrons can provide the current required to sustain the field reversal if a pre-determined stochastic magnetic field is assumed and transport is given by the RR expression. Self-consistent numerical simulations have shown that kinetic dynamo could sustain a stationary RFP configuration [371], with the fast electrons providing a substantial contribution to the poloidal current in the edge [372].

The presence of fast electrons at the edge is not a unique feature of KDT. Indeed, due to the relatively low edge temperature, the electric field generated by the MHD dynamo can produce significant deviations of the electron distribution from the classical Spitzer–Harm distribution, as measured in the edge of RFX, significantly affecting Langmuir probe measurements [373, 374]. The main difference is that the KDT predicts the suprathermal tails to be much more robust at the edge and less so in the core compared to the MHD dynamo. In RFX the trajectory deflection of frozen hydrogen pellets has been used as a diagnostic of suprathermal electrons in the plasma, showing that the pellet trajectories are better reproduced assuming the distorted electron distribution function consistent with an MHD dynamo electric field than with the fast electron population consistent with KDT modelling [371].

Runaway electrons (i.e. electrons with energies greater than 100 keV) are observed during QSH states in MST, due to their emission of bremsstrahlung hard x-ray photons: the magnetic island provides a region of reduced stochasticity where high-energy electrons are generated and well confined [375].

As described in previous sections, both in MH and QSH states the quality of the confinement in an RFP is strictly related to the amplitude of magnetic fluctuations. Moreover, a clear dependence of temperature on current has been documented as in figure 39 [102, 376], though the experiments at the highest current values in RFX-mod have not been carried out in optimized conditions. Based on these facts an upgrade of the RFX-mod device, RFX-mod2, is presently underway, aiming at the experimental study of an RFP plasma at 2 MA with an optimized magnetic front-end (section 7.3), i.e. in condition of reduced effect of tearing modes [114]. Experimentally QSH states, though more persistent at high current, have never been obtained as a stationary equilibrium. Moreover, QSH are more likely to occur and produce transport barriers at low density: still open question is if this is related to a poor control of particle recycling, related to a strong PWI at high current. Due to a milder magnetic boundary, experiments on RFX-mod2 are expected to allow more stationary helical states at higher density and to clarify the effect of microinstabilities in conditions of reduced magnetic chaos.

5. Current profile control

Typically, RFP plasmas are formed and sustained using toroidal induction by the transformer linking the torus. The parallel electric field is therefore peaked in the core, and this destabilizes tearing modes. Nonlinear saturation of the tearing instability creates a dynamo electric field that keeps the current profile close to marginal stability in either the QSH or MH regimes. Since tearing instability results from the peaked electric field, adding current drive in the outer region allows the possibility to attain a tearing-stable current profile. This is referred to as current profile control and involves the addition of one of several possible current drive sources: (1) external field coils, (2) insertable electrostatic current sources, and (3) edge-localized radio frequency (RF) current drive.
The inductive approach using external field coils as transformers has been most successful, leading to a strong reduction in magnetic stochasticity and large improvement of energy confinement in RFP plasmas. Microturbulence emerges when tearing instability is suppressed, suggesting RFP confinement might ultimately be limited by microinstability as it is in tokamak and stellarator plasmas. These advances have been very important for the development of the RFP and for the advancement of toroidal fusion science, e.g., by extending the reach of gyrokinetic modelling.

The theoretical and experimental foundations for current profile control are described in section 5.1. In section 5.2 we describe four inductive techniques, and in section 5.3 we describe measurements of the modified current profile and fluctuation reduction. Improvements in RFP fusion performance and work towards fusion parameter scaling are described in sections 5.4 and 5.5. The emerging role of microturbulence is described in section 5.6. Results with insertable current sources and RF antennas are described in section 5.7. Section 5.8 contains other related routes to improved confinement, and section 5.9 discusses questions and future directions.

5.1. Foundations for current profile control

Comparison between experiment and nonlinear MHD established that the multiple core-resonant modes are current-gradient-driven tearing fluctuations (or resistive kinks), well described by resistive MHD [37, 165], which can stochasticize the magnetic field in the core.

MHD computations demonstrated that these tearing fluctuations could be reduced or even eliminated with a suitable modification of the current profile [158, 377–379]. The essential idea is captured in figure 40 showing results from the DEBS code. On the left are profiles of the plasma current flowing parallel to B. The standard profile corresponds to steady toroidal induction, while the controlled one results from the addition of current in the plasma edge, implemented in the simulations as a localized ad-hoc electron force profile. The additional current drive leads to dramatic changes in the core-resonant tearing mode amplitudes, such that stochasticity, is substantially reduced, see figures 40(b) and (c). The fact that a change to the edge current profile can affect core-resonant tearing modes stems from the global character of these modes, such that they are sensitive to the shape of the current profile everywhere in the plasma. The added current can either be stabilizing or destabilizing, depending on its magnitude and profile (stabilizing in the case of figure 40). The added current drive can be thought of as replacing the dynamo electric field, allowing fluctuations to be reduced or eliminated.

5.2. Inductive techniques

Four approaches to inductive current profile control have been studied experimentally and, in most cases, computationally. All entail temporal variations in the toroidal and/or poloidal magnetic field. Two of them are transient techniques: pulsed poloidal (or parallel) current drive (PPCD) primarily entails a ramp down of the toroidal magnetic field, and self-similar current decay (SSCD) entails a ramp down of both fields.

The other two techniques are quasi-steady-state. OPCD entails oscillation of the toroidal magnetic field, and OFCD entails oscillation of both fields. The most thoroughly developed of these techniques is PPCD. Hence, PPCD is the primary focus of this chapter. Though less developed than PPCD, OFCD has potentially very interesting implications for an RFP based fusion reactor; it will be separately described in chapter 9.

The PPCD technique was first studied experimentally on MST [380] and computationally in DEBS soon thereafter [378] and in later work [381, 382]. Electric and magnetic field waveforms from one of the highest-performance MST PPCD plasmas are shown in figure 41. Sustained reduction of both m = 1 and m = 0 fluctuations was found to require $E_\parallel > 0$, where $E_\parallel = E_\parallel(a) = E \cdot B/|B| = (E_\theta B_\theta + E_\phi B_\phi)/|B|$ is the surface parallel electric field, and the subscripts $\theta$ and $\phi$ refer to the poloidal and toroidal direction [383]. $E_\parallel > 0$ is maintained.

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{figure40.png}
\caption{From DEBS code: (a) $J/B$ radial profile in the standard steady toroidal induction case (blue) compared to the controlled case (red) with an auxiliary ad hoc parallel electron force (magenta). Poincaré plot of magnetic field line in controlled (b) and standard (c) cases. Adapted from [158]. © IOP Publishing Ltd. All rights reserved.}
\end{figure}

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{figure41.png}
\caption{MST 500 kA plasma with PPCD: time evolution of (a) poloidal and (b) toroidal electric fields, (c) poloidal and (d) toroidal magnetic fields and (e) parallel electric field. PPCD is applied at 8.5 ms. Reprinted from [383], with the permission of AIP Publishing.}
\end{figure}
initially through $E_\|$ via a ramp down in $B_\|$. Later in time, $E_\|$ is reversed in order to prolong the control period. The ramp down in $B_\|$ brings about a rather extreme RFP equilibrium, with the toroidal field reversal parameter, $F$, reaching about $-2$ and the pinch parameter, $\Theta$, reaching about 3.5. In addition to the $E_\|$ criterion, the production of high-quality PPCD plasmas is subject to a number of other conditions, e.g., avoiding an overly large $E_\|$ [384], providing a well-conditioned plasma-facing boundary and sufficiently low initial plasma density [383, 385, 386], avoiding sawtooth crashes at the initiation of PPCD [383], and assuring a good control of magnetic error fields and of the LM [383, 387].

The SSCD technique was first considered in two independent theoretical studies [388, 389] in which it was recognized that a particular prescribed self-similar temporal ramp-down of both $B_\|$ and $B_\|$ can yield an inductive electric field that maintains an MHD stable current profile. The self-similar solutions are space–time separable such that the equilibrium profiles are stationary in time, while the magnetic flux decay sustains the current with $E_\| = \eta J_\|$ (no dynamo). The ramp-down rate is proportional to the instantaneous value of the resistive diffusion time, close to the global $L/R$ timescale. Nonlinear, 3D simulations with SSCD programming yield essentially complete tearing stabilization [390]. The loop voltage programming required for SSCD, originally dubbed ‘catching’ [388], was first attempted on the HBTX device at Culham, but the results were never published.

The OPCD technique was first tested experimentally on RFX [391] and later examined with a 1D model simulating magnetic field profiles [392]. Entailing a sinusoidal oscillation of $B_\|$, half of each OPCD cycle is PPCD-like in the trajectory of $B_\|$, and magnetic fluctuations can be reduced in an oscillatory fashion.

5.3. Current profile modification and fluctuation reduction

In this section, some examples are provided of the experimental current profile control and of fluctuation responses to current profile control, primarily from PPCD plasmas but also from SSCD. Direct internal measurements of the current profile during PPCD were made possible on MST by a combination of laser polarimetry [393] and 2D RFP equilibrium reconstruction [27]. Figure 42 compares profiles of $\mu = \mu_0 a J_\|/B$ during PPCD ($\mu$ is the same quantity plotted in figure 40) to profiles before and after sawtooth crashes, each of which leads to flattening of the current profile in steady-induction plasmas [53, 394]. Measured with approximately 0.08 m spatial resolution, $\mu$ during PPCD falls between the pre-crash and post-crash profile for $r > 0.3$ m, while the PPCD profile is more peaked than either profile in the centre of the plasma. Such peaking was also observed in MHD modelling of PPCD plasmas in RFX with the Specyl code [395], showing that during PPCD a significant reduction of MP is accompanied by a shrinking of the magnetic profiles. Measurements of the current profile were also made in the outer 10% of MST plasmas by an inserted probe [396] with 0.01 m spatial resolution. In this region $\mu$ decreases when PPCD is applied. This result was interpreted as being due to measured drops in the dynamo electric field and the electrical conductivity, the latter resulting from a drop in the electron temperature. Overall, the change in the $\mu$ profile is fairly modest with PPCD, but additional measurements in MST showed that there are much larger changes to the inductive electric field profiles [397, 398]. Hence, there is a much greater change in the current drive than in the current density profile, consistent with marginal tearing stability.

Fluctuation measurements with current profile control are numerous and show a significant drop in magnetic fluctuations. Figure 43, representing data from RFX-mod OPCD plasmas, shows that the innermost resonant (i.e., $m = 1$, $n = -7$) mode gradually increases while the remaining $m = 1$ modes with higher-$n$ decrease, which also corresponds to higher electron temperatures [261]. The reduction of higher-$n$ modes is a feature of plasmas with improved global confinement [397, 399]. While most measurements of magnetic fluctuations are restricted to magnetic sensing coils at the plasma boundary, the polarimeter in MST allowed for first time the measurement of the fluctuation reduction in the core [394, 400] which has been found to be broadband and substantial.

Ohm’s law in steady-induction RFP plasmas is balanced by the dynamo electric field, see chapter 3. In plasmas with
substantial dynamo fluctuation reduction, such as with PPCD, the dynamo electric field becomes small. This was demonstrated by measurements in PPCD plasmas in MST, figure 44 [178, 398]. In this figure, the electric field and resistive current terms in parallel Ohm’s law are shown to be approximately equal over the entire plasma, implying that these plasmas are dynamo free. Consistent with the discussion earlier in this section, the $\eta_{neo}J_{\parallel}$ term does not change much between steady-induction and PPCD plasmas, but $E_{\parallel}$, along with the dynamo electric field, changes substantially [178, 398].

With the drop in the core-resonant magnetic fluctuations, changes occur in the behaviour of other fluctuating quantities. For example, in the core of MST PPCD plasmas, density fluctuations, measured with interferometry and a heavy ion beam probe, are found to drop significantly [401–403]. Potential fluctuations drop as well, both in the core [403] and in the edge [404]. This is not unexpected given the previously measured correlation between core-resonant magnetic fluctuations and edge-localized potential fluctuations in steady-induction plasmas in the EXTRAP-T1 RFP [405]. There is also a more localized tapering of the fluctuations in MST correlated with a region of strong flow shear (the flow shear is addressed further in section 5.8). Measurements in the edge plasma of MST during PPCD also reveal that higher-frequency magnetic turbulence is much less intermittent and has a statistical behaviour typical of self-similar turbulence, as expected in systems exhibiting self-organized criticality [406]. A similar conclusion was drawn for both magnetic and potential fluctuations in the edge of TPE-RX PPCD plasmas [407].

The reduction in core-resonant tearing fluctuations also brings about an important change to the core magnetic topology. Tomographic x-ray measurements in the core of MST PPCD plasmas revealed two discrete, non-overlapping islands, corresponding to the two innermost resonant $m = 1$ modes [135, 408]. This is shown in figure 45 and is a direct evidence of decreased stochasticity in the core. High-energy runaway electrons also emerge in these plasmas, as described in more detail in the next section, further indicating that stochasticity no longer dominates core transport.

5.4. Improved fusion performance

This section describes improvements in fusion performance brought about by SSCD, OPCD and PPCD. Most of the results are from plasmas with PPCD, and with only two exceptions, noted below, the results described here are from ohmically heated plasmas. The papers cited herein contain results from both hydrogen and deuterium plasmas.

The only published test of SSCD occurred in RFX-mod [409]. The simultaneous decay of $B_\theta$ and $B_\phi$ is provided primarily by actively reversing the sign of the surface toroidal voltage. This affects both the toroidal and poloidal current flowing in the plasma. As magnetic fluctuations drop, poloidal beta, $\beta_p = 2\rho_0(p)/B_\phi^2(a)$ increases by approximately 50% along with the global energy confinement time. A crude approximation of SSCD was also tested on MST without programmable power supplies, with the decay of $B_\theta$ and $B_\phi$ provided passively by the external poloidal field circuit. The energy confinement time in such plasmas is estimated to triple.

The quasi-steady OPCD technique was tested in both RFX [391, 410, 411] and RFX-mod [261]. One signature of OPCD is an oscillation in the central electron temperature, which increases in the PPCD-like half of OPCD cycle and then decreases in the other half. In plasmas with OPCD during the
PPCD-like phase, a transition to a QSH state occurs. A helical $T_e$ structure emerges and the maximum $\nabla T_e$ corresponds to an electron thermal diffusivity $\chi_e$ of about 10 $m^2$ s$^{-1}$. The global energy confinement time increases by about 60% [261].

The PPCD technique has been the most widely applied and studied of the inductive approaches, and it has brought about the most substantial improvements in RFP fusion performance. PPCD has been tested in MST, RFX, RFX-mod, EXTRAP-T2R, and TPE-RX, and it has led to (1) improvements in energetic electron and ion particle confinement and in thermal electron particle confinement, (2) large values of the normalized electron density, substantially exceeding the Greenwald limit, (3) classical confinement of impurity ions, (4) RFP-record values of electron and ion temperature, with a simultaneous reduction in ohmic heating power, (5) an RFP-record beta that exceeds the Mercier criterion, (6) an RFP-record tokamak-like energy confinement time, and (7) the emergence of two new regimes for the RFP, with the prominence of micro-instability in lower-$\beta$ plasmas and pressure-driven macro-instability in higher-$\beta$ plasmas. We expand upon these results in what follows.

One of the key PPCD results was the observation of runaway electrons in the core of MST plasmas [412]. In a stochastic magnetic topology, the rate of electron transport increases with the electron velocity. Hence, while the Dreicer criterion port; a diffusion model consistent with stochastic magnetic transport; $\alpha = 0$ allows no energy dependence to the diffusion coefficient, suggestive of diffusion driven by electrostatic fluctuations. The experimental x-ray energy spectrum from steady-induction plasmas is well fitted by an exponential curve with $\alpha = 1$, while the spectrum from PPCD plasmas is well fitted with $\alpha = 0$, indicating that stochastic transport is no longer playing an important role.

High-energy (20 keV) beam-injected ions may also be modestly better confined during PPCD [413]; however, it has to be mentioned that beam ions are already well confined in steady-induction MST plasmas [349].

Thermal electron particle confinement is also improved with PPCD. From measurements in MST, the particle flux decreases everywhere in the plasma, with a drop of almost two orders of magnitude in the core. Correspondingly, the global particle confinement time is estimated to improve eightfold [401, 402]. The rate of particle diffusion also drops by about a factor of two in RFX plasmas [395]. The combination of PPCD and frozen-deuterium pellet injection in TPE-RX brought about an estimated five-fold improvement in the global particle confinement time [414]. Absent pellet injection, the global confinement time improved in TPE-RX by an estimated factor of ten [407, 415].

Frozen pellet injection allows a substantial increase in the density attainable in PPCD plasmas, while retaining reduced fluctuations and improved confinement. Raising the density with edge-deposited fuelling techniques such as gas puffing bring about increased fluctuations and degraded confinement [246, 414, 416, 417]. Pellet injection can fuel the plasma core without substantially impacting the edge or degrading confinement. The synergy between pellet injection and PPCD was first noted in RFX [395, 418], where the density doubled, reaching about $4 \times 10^{19}$ m$^{-3}$. Pellet injection into TPE-RX brought about a similar doubling of the density [414], and injection into MST plasmas initially quadrupled the density, surpassing the Greenwald density limit [246, 416, 417]. Shown in table 2 are central densities and densities normalized to the Greenwald limit in MST. Subsequent injection of larger pellets in MST has resulted in an absolute density up to $9 \times 10^{19}$ m$^{-3}$ and a normalized density up to $\sim 2$ [100, 419].

The confinement of impurity ions also improves. In steady-induction RFX plasmas, the impurity particle diffusion coefficient profile drops from 20 $m^2$ s$^{-1}$ in the core to 1 $m^2$ s$^{-1}$ in the outer 10% of the plasma. With PPCD, the lower diffusivity region widens to encompass the outer 20% of the plasma [411]. In EXTRAP-T2R, the rate of impurity diffusion decreases by about a factor of two in the core with no change.
in the edge [420]. In MST, impurity transport drops down to the classical value [289, 324, 325]. Dominated by Coulomb collisions rather than neoclassical particle drift effects, impurity ions are convected outward, and the impurity density profile, along with the $Z_{\text{eff}}$ profile [421], become hollow, consistent with the temperature screening mechanism of classical transport.

One common feature of PPCD plasmas is an increase in the central electron temperature, e.g., [82, 422–426]. This is in spite of the fact that the ohmic heating power commonly drops, e.g., [380, 427]. The largest central $T_e \sim 2$ keV was achieved in MST, table 2 [428]. Neutral beam heating has also been applied to PPCD plasmas in MST, increasing the central electron temperature by about 100 eV [354].

In contrast to electron temperature, the evolution in the ion temperature in low-density PPCD plasmas is typically fairly small. The only localized measurements of the ion temperature in the core of PPCD plasmas were made in MST using neutral-beam-based charge-exchange-recombination spectroscopy. Initial measurements in low density plasmas showed that the central $T_i$ changes very little between steady-induction and PPCD plasmas, with a value of about 300 eV in the core of plasmas with a central $T_e > 1$ keV [383]. The ohmic heating power channelled from the electrons to the ions is quite small. A roughly four-fold increase in the ion temperature, to about 1.3 keV (table 2), was achieved by triggering large amplitude sawtooth crashes shortly before the start of PPCD [428]. The ion heat generated during the crashes is trapped in the plasma with the onset of PPCD and the ion energy confinement time is estimated to improve ten-fold. With pellet injection and substantially higher density, the ion temperature increases with time during PPCD plasmas in MST [416]. This is due both to a larger ohmic input power and a larger rate of energy transfer from the electrons to the ions.

Another common feature of PPCD plasmas is an increase in $\beta$. There are examples of this from RFX, TPE-RX, EXTRAP-T2R, and MST, e.g. [83, 424, 429, 430]. This will be discussed in detail in chapter 6.

In PPCD plasmas, an increase in the global energy confinement time is observed in all RFP devices. The confinement time in RFX and EXTRAP-T2R as much as doubled [83, 430], while in TPE-RX it increased as much as five-fold [251, 426]. The largest increase in confinement, as well as the largest absolute RFP confinement time was produced in MST. The energy confinement time increased by more than ten-fold, reaching about 12 ms [246, 428]. The global electron thermal diffusivity drops to about 5 m$^2$ s$^{-1}$, with a local minimum of about 2 m$^2$ s$^{-1}$ in the region of large $\nabla T_e$. The values of confinement times reached in MST are tokamak-like [397, 399, 428]. This is illustrated in figure 47, where MST standard and PPCD confinement data are compared to a tokamak reference that has the same plasma size, magnetic field strength, and heating power as specified by the IPB98(y,2) ELMy-H mode empirical scaling.

5.5. Confinement scaling

Considering plasmas with PPCD, for which there are results from multiple RFP devices, there is not as yet a firmly established universal scaling for energy confinement time, although basic parameters such as plasma minor radius very likely play a role. As suggested by figure 40(b), even with current profile control there may be regions of the plasma where stochasticity still plays a role in energy transport. In addition, as discussed in the next section, in lower-density (lower-$\beta$) PPCD plasmas, micro-scale fluctuations due to instabilities like the TEM may be important, while, as described in the next chapter, in higher-$\beta$ PPCD plasmas, macro-scale pressure-driven tearing may play a role.

As reported in figure 48, for low-density and high-density PPCD plasmas in MST, the improved energy confinement fell above, and in one case well above, the empirical Connor–Taylor scaling [4] which was based on the best standard-plasma confinement results from a variety of RFP devices [383, 417, 424]. Hence, PPCD plasmas represent a break from the previous trend.

The dependence on the dimensionless Lundquist number $S$, which was introduced in chapter 2 and varies as $I_p T_e^{3/2}/n_i^{1/2}$, was examined based on data from RFX [63, 64]. While magnetic fluctuations were smaller (and the energy confinement time larger) in RFX PPCD plasmas, the fluctuations exhibited the same amplitude scaling with $S$ as steady-induction plasmas.

The factor by which the energy confinement time improves with PPCD in low-density TPE-RX, RFX, and MST plasmas, including results up to 2002, was found to be well represented.
5.6. Emergence of microturbulence in PPCD plasmas

When tearing modes are sufficiently reduced it is plausible that microturbulence might ultimately govern energy and particle transport in RFP plasmas, as it does in tokamak and stellarator plasmas. The non-stochastic confinement properties of energetic electrons described above provided the initial indication this might be the case for PPCD plasmas. The role of microinstability in QSH plasmas was already discussed in section 4.3.

Density fluctuations measured by FIR forward scattering in MST’s standard RFP plasmas are strongly correlated with the magnetic fluctuations associated with tearing modes [435]. These fluctuations are also greatly reduced with PPCD, consistent with the large reduction in magnetic fluctuations. The spectral shape of the density fluctuations changes with PPCD, and a higher frequency peak emerges around 50 kHz, identified as density-gradient-driven TEM. The fluctuations have $k_{\perp} \rho_s \approx 0.1 - 0.2$ and propagate in the electron diamagnetic drift direction [436].

As shown in figure 49, the onset of the TEM density fluctuations exhibits a clear critical-gradient threshold, $R/L_n \approx 20$, where $L_n = n_e / |\nabla n_e|$ is the density scale length. This agrees with linear gyrokinetic modelling (GENE) predictions [300, 437]. The critical gradient threshold is larger than for tokamak plasmas by roughly the aspect ratio, $R/a$, owing to the difference in magnetic field scale length. Nonlinear GENE modelling shows a large Dimits-like shift from exceptionally strong zonal flows that suppress the TEM turbulence to near zero. This was initially a surprise, but it was discovered that a small magnetic perturbation, e.g., a remnant tearing mode, disrupts the zonal flows, as shown in figure 50. With weaker zonal flows, the turbulence saturates at larger amplitude, comparable to the level measured in MST [300, 302]. An intriguing prospect emerges: as tearing is further suppressed, either through further improvements in profile control or towards pure SH states, anomalous transport might be uniquely small in RFP plasmas.

A zonal flow has been identified experimentally in MST plasmas. Measurements using two probes with capacitive electrodes separated 180° toroidally that measure the plasma potential directly observe a zonal flow, most likely associated with the TEM turbulence [438]. Such direct observations of zonal flows are rare in fusion research, despite their underlying importance in the formation of transport barriers.

The impurity flux associated with the TEM turbulence has also been measured using a novel linearized correlation method of active spectroscopy [439]. The spectral method could be used to investigate turbulent transport that limits impurity accumulation in the core of high-performance, metal–wall tokamak and stellarator experiments.

5.7. Towards non-inductive (dc) current profile control

Three additional efforts were undertaken on MST with the goal of non-inductively (non-transiently) modifying the current profile and will be briefly summarized in this section: driving current via electrostatic current sources inserted into the...
plasma edge and launching lower-hybrid (LH) and electron-Bernstein waves (EBWs) from RF antennas.

Edge electrostatic current drive was tested with 16 miniature plasma sources inserted into the outer 10% of the plasma [440]. When the sources were inserted but inactive, they caused a perturbation leading to an estimated 50% drop in particle and energy confinement. When biased negatively to emit electron current, each source injects roughly 0.5 kA, and when the current is injected parallel to the background edge parallel current, magnetic fluctuations are reduced, and the particle and energy confinement are restored to approximately their unperturbed values. The restoration of confinement is linked to a time-averaged drop in magnetic fluctuation amplitude, due primarily to an increase in the time between sawtooth crashes. Although the plasma sources did not prove a viable means of achieving a net confinement improvement, they did demonstrate the principle of current profile control.

The foundation for LH wave injection stemmed in part from theoretical studies [441, 442]. The aim is to generate an RF current drive profile similar to the ad hoc electron force profile shown in figure 40. It was predicted that the location of the driven current can be controlled by choice of parallel wave number and RF frequency, and that the predicted efficiency of current drive is sufficient for tearing suppression with 1–2 MW of RF power. Experimentally, an interdigital-line antenna mounted on the inner surface of MST launched the correct $n_3$ spectrum and was well loaded by the plasma [443, 444]. With the source of RF power exceeding 200 kW, substantial x-ray emission from the plasma was observed [445]. However, robust operation at this high source power did not prove feasible, implying that a large number of antennas would be required to inject sufficient power to affect tearing mode stability.

The foundation for the injection of the EBW was laid in part by computational studies showing that there is a directional component to the absorption of the wave, which can be controlled by the poloidal angle at which the wave is launched, so that directed current drive is possible [446]. The conversion of Bernstein waves generated within the plasma to electromagnetic waves was confirmed by measurements of thermal emission at the plasma boundary [447]. By reciprocity, this implies that externally launched electromagnetic waves can couple to the Bernstein mode. Starting with very low power (<10 W), coupling of RF was studied experimentally with a waveguide antenna and found to agree with full-wave modelling [448]. Coupling efficiency exceeding 80% was observed under certain conditions, though coupling can be degraded by effects such as edge density fluctuations [449]. The available power for injection was increased to several hundred kW, and electron energization in the plasma is evident from x-ray emission [450]. However, recent modelling with a Fokker–Planck, RF-altered electron distribution coupled into a single-fluid MHD simulation indicates that much larger power (∼10 MW) would be needed to produce an observable effect on the tearing modes [451].

5.8. Other routes to improved confinement

In addition to current profile control and helical states, there are other routes to improved confinement in the RFP that do not entail active control, per se. We mention these here, given an observed partial synergy with and partial similarity to PPCD plasmas in MST.

In both TPE-1RM20 and MST, spontaneous periods of improved confinement were observed in conjunction with strong toroidal magnetic field reversal and a relatively large pinch parameter. In such regimes, dubbed ‘improved high theta mode’ in TPE-1RM20 [452] and ultimately ‘enhanced confinement’ (EC) periods in MST [453, 454], magnetic fluctuations drop, and both the electron temperature and the energy confinement increase, doubling in TPE-1RM20 and tripling in MST. Improved confinement was also observed in TPE-RX, corresponding to discharges with slow decay (towards zero) of the edge toroidal field [455].

The EC periods in MST are characterized in part by regularly occurring bursts of $m = 0$ mode activity. Each burst causes a temporary degradation of confinement [57]. Such bursts can also occur in PPCD plasmas, with a similar effect on confinement [319]. The hottest, highest confinement PPCD plasmas are entirely free of such bursts [424], with reduction of both $m = 1$ and $m = 0$ occurring almost immediately after the start of PPCD [383]. This immediate reduction occurs when an EC period is produced leading into the start of PPCD. In fact a possible explanation of the lower effect of PPCD in RFX and RFX-mod with respect to MST has been ascribed to a lower reduction of $m = 0$ modes. Furthermore, EC periods can terminate with very large amplitude sawtooth crashes, contributing to the pre-PPCD ion heating described above [428]. Initiating PPCD with (EC-like) strong toroidal magnetic field rever-
Sal also led to the best PPCD performance in EXTRAP-T2R [430].

The physics of these spontaneous periods of improved confinement, and the $m = 0$ bursts occurring in EC and PPCD plasmas in MST, is not well understood. However, based on the fact that the bursts bear some resemblance to edge-localized modes in tokamak H-mode plasmas [57, 456], it was hypothesized that $E \times B$ flow shear could play a role. In the PPCD case, this would not be overly surprizing given the emergence of tokamak-like microinstabilities. Actually strong flow shear was measured in the edge of both EC and PPCD plasmas [404, 454]. Later modelling of the impact of various flow profiles on tearing stability showed that the flow shear in EC and PPCD plasmas could play a role in governing the $m = 0$ bursts [457]. Modifications to $E \times B$ flow shear were also measured during PPCD in RFX [458, 459]. Finally, it is worth mentioning that biasing the edge of MST brought about a reduction of electrostatic fluctuations and an improvement in particle confinement [460]. Electrostatic fluctuations and the edge particle flux were also affected by biasing in RFX [461].

5.9. Open questions and future directions

There are many questions in the realm of current profile control. For example: (i) what is the optimal inductive $E(a)$ waveform for sustained fluctuation reduction and improved confinement, and can the degree and duration of improved confinement be increased? (ii) How does the improved energy confinement time scale with parameters like the plasma current? (iii) What would be the impact of an advanced, low-recycling boundary on plasmas with current profile control? (iv) How are PPCD/OPCD plasmas associated to the onset of QSH states? (v) What is the role of $m = 0$ modes? Though the comprehension of these plasmas has significantly progressed, further computational and experimental work is required to fully address these points.

While MHD computation provided the foundation for current profile control, experimental progress to date has been almost entirely empirical. This is due in part to the fact that fully nonlinear simulations of RFP plasmas require considerable raw computing power, the need for which increases, for example, with $S$. The most recent modelling of PPCD with the NIMROD MHD code was conducted with $S < 10^4$ [382], much lower than in the hottest PPCD plasmas in MST. The importance of $m = 0$ bursts in MST plasmas was described above, but such bursts are not typically observed in MHD computation, implying unrealistic input parameters or missing physics.

Gyrokinetic modelling is also likely to be important. The potentially missing physics noted above might be related to micro-instability. Hence, understanding present day experimental results and projecting to the future will likely require both MHD and gyrokinetic work.

In present-day RFP devices, current profile results have for the most part stemmed from simple, passive power supplies comprised of capacitors and inductors. The future of current profile control lies in digitally programed, feedback-controlled power supplies, which are in use on RFX-mod and are being developed on MST. In agreement with the modelling described just above, the power supply waveforms can be adjusted in real time based on measurements of the MHD dynamo, or at least the constituent fluctuations. But additional inductive and/or non-inductive control may also be needed for micro-scale fluctuations.

The scaling laws found in present-day devices need to be assessed at higher $I_p$, which might also allow for non-inductive control techniques like RF.

With respect to power and particle handling, RFP devices have to date operated with primitive boundaries that would be irrelevant in a reactor setting. While substantial progress in fusion performance has nonetheless been possible, it was noted above that a crucial requirement for high-quality PPCD plasmas is a ‘well-conditioned boundary,’ meaning minimized contamination by atmospheric impurities and low recycling. Indeed MHD modelling showed that changing the edge resistivity changes magnetic fluctuation amplitudes [462]. It is customary to think of current profile control in terms of changing the inductive electric field, but the current profile can be altered equally well by a change in resistivity. Together with the already mentioned possible effect of the $m = 0$ modes (section 5.8), also the need for low recycling in PPCD may explain in part the fact that the confinement improvements in TPE-RX and MST were greater than that in RFX and RFX-mod: with full graphite coverage of the plasma-facing boundary, RFX and RFX-mod had large recycling, in contrast to the other two devices with their metallic boundaries.

In this chapter and in the previous ones two ways of reducing stochasticity and improving the performance of RFP plasmas have been presented: the helical states and the current profile control. Still an open issue is whether current profile control could effectively act in QSH states. An extended experimental investigation has not been done yet. In RFX-mod at plasma current $I_p < 1$ MA, OPCD favoured the development of QSH leading to an enhancement in confinement by $\approx 50\%$ [261]. At high current QSHs spontaneously develop, so that their time phasing and duration with OPCD is difficult to control. New insights into the combined effect of QSH and OPC/PPCD will come from RFX-mod2, where more stationary helical states are expected to occur.

6. Density and $\beta$ limit

The high-density limit is nearly the same for the RFP and the tokamak, associated to a radiation barrier. It is traditionally described by the empirical Greenwald limit [463] for both configurations, suggesting that the same mechanism is behind such behaviour. Instead in the stellarator, where a high density limit exists as well, it is described by the Sudo scaling [464]. More recently, a power balance model accounting for neutrals and impurity radiation has been developed and successfully describes the high density limit in the stellarator, RFP and L-mode tokamak [465, 466]. While the RFP has the same empirical density limit as for the tokamak, the absolute value of density is large because the current density is large. An RFP reactor plasma would operate below the empirical limit, as illustrated in the TITAN power plant study (see chapter 10).
At variance with the density limit, the $\beta$ limit is different in tokamaks and RFPs. In the tokamak, the total $\beta$ is basically limited by ideal stability, which sets a disruptive limit around 5%, depending on current and aspect ratio. In addition, limitations of achievable $\beta$ related to neoclassical pressure driven resistive tearing modes have been observed in several devices, and compared to theory, see for example [467–469]. In the stellarator, ideal interchange modes are a major concern for high $\beta$ plasma production, leading to a non-disruptive limit similar to the tokamak one [470, 471], though violations of the Mercier criterion [5] have been observed in some experiments and interpreted as a local effect [472, 473].

Instead, plasmas characterized by high $\beta$ values are a feature of the RFP configuration, where, due to edge magnetic field being nearly poloidal, the total $\beta$ is almost equal to the poloidal one. As reported in papers as [36, 474], the ideal limit is much higher than in tokamaks, $\beta_p = 50\%$ for $m = 0$ instabilities and $\beta_p = 100\%$ for $m = 1$. However, current experiments achieve $\beta$ values that, though high compared to the tokamak and stellarator, are still well below the ideal stability limit.

### 6.1. Density limits

Since 80s, the operational range of RFP experiments has been found to be limited both at low and high electron density [50, 475]. The low-density limit was related to the acceleration of runaway electrons and to the growth of the magnetic fluctuations. In RFX-mod the longest durations of helical states featuring ITB have been obtained when the careful design of the magnetic boundary and the consequent reduction of the edge radial field by a proper tuning of the feedback MHD control system (clean mode control, see section 7.2) has allowed the extension of the density operational range towards the lower values [121].

Similar to tokamaks, the high density limit in RFPs has been associated to a radiation barrier, dependent on plasma current [476]. In RFX-mod, the high density limit was compared to the tokamak empirical Greenwald limit [463]:

$$n_e \left(10^{20} m^{-3}\right) < n_G = I_p[MA]/\pi a^2$$

(6.1)

with $n_e$ average plasma density, $I_p$ plasma current. It was found that in the RFP as in the tokamak the upper boundary of the density operating space is limited by the Greenwald density $n_G$ [477]. The same empirical limit was also observed in the MST RFP in standard plasmas [417], see figure 51. The fact that the same representation, based on averaged global plasma quantities, matches well data from devices characterized by different size, performances, magnetic topology and transport quality suggested that a common physics could underlie the phenomenology of the Greenwald limit. Indeed, the point of view of the RFP configuration has been compared to those of the tokamak and stellarator, significantly contributing to the general understanding of the high density limit [478]. A theory based model able to describe the density limit in RFPs, L-mode tokamaks and stellarators has been proposed, simply based on a radiative limit and taking into account the radiation from impurities and neutrals [465, 466]. The model leads to a constraint for the edge density, ruled by radiation losses from light impurities. For the ohmic tokamak and the RFP, the limit scales linearly with the plasma current, as the empirical Greenwald one. The auxiliary heating adds a further dependence, scaling with the power as $P^{0.4}$, in agreement with L-mode tokamak experiments. Compared to previous works with a similar approach and applied to ohmic tokamaks [479], this model is able to take into account profile-dependent terms.

In RFX-mod, when the density approaches the Greenwald value, the temperature decreases, while the density increases at the edge [274]. Such edge density accumulation takes place after a back transition from the QSH to the MH regime, which occurs at $n/n_G \approx 0.35$ [208]. Due to the edge density accumulation radiation increases, in the shape of a poloidal radiating belt [477, 480], which resembles the MARFE phenomenology observed in tokamaks (with poloidal and toroidal coordinates exchanged due to the poloidal nature of the magnetic field in the RFP edge) [481] or Serpens in LHD stellarator [482].

An important difference between RFPs and tokamaks is that the density limit is not disruptive in the RFP: when the Greenwald density is approached, due to plasma cooling and resistivity increase, the toroidal flux decreases, resulting in the shrinking of the current profile with soft decrease of plasma current. Indeed, in the RFP when density increases, the input power necessary to sustain the discharge increases as well, so that the density limit can be regarded as a power limit, governed by the volt-seconds stored in the magnetizing circuit of the device [274].

Only in some circumstances, at high current and density, fast terminations have been observed in RFX, RFX-mod and MST [417, 483]. In particular, in RFX-mod, where they occur mainly at very high current and in conditions of saturated graphite wall, they have been related, to strong localized PWIs in the region of wall mode locking, where severe power loads are dissipated, leading to a sudden particle release, locally enhanced radiation and temperature drop.

Dedicated experiments in RFX-mod led to the interpretation of the density limit based on the relation between the above-mentioned localized accumulation process and the $m = 0$ mode magnetic topology. The main experimental observations are summarized in figure 52 as a function of the normalized toroidal angle $\varphi - \varphi_{\text{lock}}$ where $\varphi_{\text{lock}}$ is the angle corresponding to the maximum deformation of the $m = 1$ mode bulge due to wall mode locking (LM) [484]. The maximum $H_n$ emission corresponds to the maximum $m = 1$ mode deformation, (green and red lines in panel (d)) thus indicating at that location the main particle source. At the edge, black line in panel (a) shows the Poincaré plot of the $m = 0, n = 1$ island; the plasma flow, measured by the gas puffing imaging diagnostic [485] corresponds to the blue line and reverses at the X-point of such island. This indicates that particles are carried from the source toroidal position to a stagnation point where density accumulates (panel (c)) and, correspondingly, radiation is enhanced (panel (b)). The understanding of these observations relies on the application of the guiding-centre code ORBIT [130], upgraded to include the effect of a recycling wall [484]. The whole process is triggered by a charge separation, induced
by a differential diffusion between ions and electrons, which at high density leads to small order ($\approx 10^{14}$ m$^{-3}$) electron accumulation at the $m = 0$ island X-point. To ensure ambipolarity, a radial electric field develops, which originates a convective cell, ultimately producing the macroscopic density accumulation (order $10^{20}$ m$^{-3}$). In [486] the model has been generalized: it is shown that in any kind of device if a perturbation breaks the toroidal symmetry with a drift depending on the gyroradius (larger for ions than electrons), an ambipolar potential, with the same symmetry as the perturbation, is produced to balance the drifts.

It is well known that in tokamaks the Greenwald limit has been exceeded in several conditions, in particular including additional heating and peaking of density profile in enhanced transport regimes and the modification of the source (pellet injection, strong NBI, supersonic molecular beam), as for example reported in [463, 487]. The same is true for RFPs, where also plasmas with $n/n_G > 1$ have been produced in several conditions.

In RFX-mod, where the operation at the highest densities is limited by the power available in the plant, the Greenwald value can be often exceeded at low current levels ($< 0.3$ MA) [275]. Discharges with He as filling gas have been found to be favourable for $n/n_G > 1$ plasmas [477]. This finding, similar to what previously observed in ASDEX-Upgrade [488], is attributed to a less localized PWI in the region of LMs and, in the case of He, to a different edge flow pattern [489].

The injection of hydrogen/deuterium pellet did not allow overcoming the Greenwald limit in RFX-mod: this has been ascribed to the available pellet sizes, too large to penetrate into the plasma without destroying the helical core or too small to reach the plasma centre before the complete ablation. Pellets have rather been injected in RFX in order to study particle transport, comparing different plasma regimes, and in particular characterizing the enhanced particle confinement time in QSH states (from $\approx 3–4$ ms in MH to $\approx 12$ ms in QSH) [26] and with pulse poloidal current drive [418].

Instead, as discussed in chapter 5, in MST $n/n_G = 1.4$ ($n_e = 0.9 \times 10^{20}$ m$^{-3}$) has been observed at $I_p = 0.5$ MA in pellet-fuelled PPCD phases [100, 417]; also, in discharges with gas puffing through a high throughput valve $n/n_G = 1$ has been achieved [417].

The radiative limit is not the only high-density boundary in RFPs: a soft density limit has been observed both for the PPCD and for QSH state development. The former ([417, 430]), weakly increasing with plasma current, is related to the onset of bursty $m = 0$ activity, and in MST requires to start the current drive with a density below $0.8–1.1 \times 10^{19}$ m$^{-3}$ in the current range $0.2–0.5$ MA [383]. The latter is also linked to a threshold in $m = 0$ mode amplitude and requires $n/n_G < 0.35$ for the spontaneous development of QSH states [211].

Given the relationship between $m = 0$ modes and upper density limit, the effect of the reversal parameter $F$ has been studied too, both in RFX-mod and MST. As predicted in [179], the magnetic equilibrium plays an important role in governing the $m = 0$ modes. Their amplitude is lower when the field reversal is shallow. In fact, in RFX-mod the operation at high current is compatible with higher densities when the $F$ parameter is close to zero, which corresponds to reduced $m = 0$ mode amplitude and milder PWI [208]. Therefore, the operation at shallow $F$ is the preferred one to perform discharges at high current, since it helps in operating just below the density threshold for spontaneous QSH states. On the other hand, in such condition the distance of the reversal from the wall is small and the islands easily overlap the wall, which makes more difficult the control of the density and easier the onset of the radiative instability [274]. In MST, results reported in [100, 417] show a modestly larger density limit with $F \approx 0$.

### 6.2. High $\beta$ plasmas

Experimentally, high beta values are usually related to high density and low fluctuation (i.e. low transport) regimes. In the TPE-1RM20 RFP, in the so-called improved high theta mode the highest $\Theta = B_y(a)/(B_\phi)$ $\approx 2$ has been found to correspond to increased densities and temperatures, and therefore to the highest values of $\beta_p \approx 18\%$ [452]. This regime was related to a reduction of dynamo activity, measured as a large reduction of magnetic fluctuation power around 70 kHz. Similar $\beta_p$ values have been measured in TPE-RX when the electron density was increased by fast gas puffing [490]. Indeed, in
most RFPs increasing density in improved regimes by pellet injection or fast gas puffing has been applied as a means to improve confinement and achieve record $\beta$ values. In RFX-mod, the highest poloidal beta values of $\approx 15\%$ have been observed when $nL_\text{TC}$ approaches $1$ [491]. A statistical analysis in terms of global parameters and assuming ion temperature as a constant fraction of the electron temperature, yields $\beta_p \approx I_p^{-0.23} (nL_\text{TC})^{0.56}$, which indicates a small degradation of $\beta_p$ with plasma current and a more significant increase with $nL_\text{TC}$ [273]. Instead, 3D resistive MHD simulations by the DEBSP code, based on stochastic transport scalings [253] predict a dependence of $\beta_p$ on $I_p$ more negative than the experimental RFX-mod data ($\beta_p \approx I_p^{-0.56} (nL_\text{TC})^{0.56}$). In MST

also a modest decrease of $\beta_p$ with the current, $\beta_p \approx I_p^{-0.23}$ at high $nL_\text{TC}$, was reported in [62].

In several RFP devices the maximum achieved $\beta$ has been transiently enhanced by the combination of high-density plasmas and active control of the current profile (see chapter 5). In RFX, when an OPCD was applied, $\beta_p$ was enhanced by a factor 50% [492]. In the TPE-RX device, double $\beta_p$ (13%) and $\beta$ (11%) values have been measured during PPCD with respect to standard plasmas. In MST, the application of PPCD leads to enhanced total $\beta$ of $11\%–15\%$ (depending on plasma current), to be compared to $5\%–9\%$ in standard plasmas [399]. When during PPCD phases deuterium pellets were injected, thus increasing density and its gradient, record $\beta$ values have been obtained [246].

The TPE-RX experiment achieved increased density operation with improved confinement and $\beta_p$ up to the record value \( \approx 34\% \), by the injection of a deuterium pellet during PPCD [12].

In MST, dedicated studies on $\beta$ limit have been performed, optimizing pellet parameters in PPCD plasmas in order to produce its ablation in the core and maximize $\beta$ [100, 417]. In these experiments, as introduced in chapter 5, while density increases the electron temperature decreases, and a large fraction of ohmic power (higher at higher density) is collisionally transferred to ions. The larger stored energy leads to high beta values, $\beta \approx 26\%$, corresponding to a $\beta_p \approx 40\%$. The latter $\beta$ value coincides with the prediction by the Connor–Taylor scaling, $\beta \approx (m_e/m_i)^{1/6} (m_e$ and $m_i$ electron and ion mass respectively), which gives 25.4% for MST [417]. In such plasmas, pressure gradients exceeding the Mercier limit have been produced [246, 417], see figure 53, similar to what observed in stellarators (see for example [471, 493]).

The Mercier criterion governs the interchange stability, and can be expressed as:

$$r \frac{B^2}{2 \mu_0} \left( \frac{q'}{q} \right)^4 + 4f' (1 - q^2) > 0 \quad (6.2)$$

with the first term related to the stabilizing effect of the magnetic shear and the second one to the destabilizing effect of pressure gradient when $q < 1$.
Theoretical computations by the DEBS code have shown that in these high \( \beta \) regimes both pressure-driven interchange modes and tearing modes are linearly unstable [246]. In MST high \( \beta \) plasmas, interchange modes have not been observed (possibly due to limits of the diagnostics), so that an open issue remains whether the driving modes are interchange or global tearing modes. In [494] it is found that at low \( \beta \) the parity of pressure driven modes is the tearing parity. At \( \beta \) values around 20\% the interchange modes appear and become dominant around 40\%. Indeed in RFX-mod the observed edge magnetic fluctuations corresponding to linearly unstable pressure driven modes have the tearing parity, though they were in a first time interpreted as g-modes [495].

A disruptive limit of \( \beta \) does not occur, even when the control of the current profile with the related reduction of current driven modes and increase of \( \beta \) leads to a destabilization of pressure driven modes: at very high \( \beta \) a ‘soft’ degradation of confinement is observed, likely related to increased fluctuations at high density [417].

At variance with tokamaks, shaping effects have not been found in RFPs to enhance the \( \beta \) limit for both ideal [496] and resistive \( \beta \) modes, due to a poloidal mode coupling stronger than in the circular topology (produced by the variation of the poloidal field along the poloidal angle) and to the development of multiple trapping regions [497].

6.3. Neoclassical bootstrap current at high \( \beta \)

The bootstrap current \( J_{BS} \) depends on temperature and pressure normalized profiles and on the trapped particle fraction:

\[
\langle J_{BS} \cdot B \rangle = L_{31} \left[ \frac{p_e^{\prime}}{p_e} + \frac{1}{Z} \left( \frac{p_i^{\prime}}{p_e} + \frac{T_e}{T_e} \right) \right] + L_{32} \frac{T_e}{T_e} \tag{6.3}
\]

with \( p_e \) electron kinetic pressure \( L_{31}, L_{32}, \alpha \) functions depending on the trapped particle fraction and on the magnetic equilibrium [498, 499]. In an axisymmetric RFP, the fraction of trapped particles in the core is similar to the fraction in a tokamak with the same aspect ratio, while it becomes lower at the edge, depending on the current profile [62]. The corresponding bootstrap current is instead much lower than in tokamaks, due to the higher poloidal field. In [284], it has been calculated by two different codes (RFXTOR, based on RFX-mod magnetic measurements [25] and FLOW, solving a generalized version of Grad–Shafranov equation [500]). The result is shown in figure 54 and confirms that in present day circular RFPs only a negligible additional current can be provided by the bootstrap mechanism [28]. The total fraction of trapped particles increases in the presence of a helical structure, as in RFP QSH states, basically due to an additional field ripple introduced by the helical geometry, evaluated from 30\% in the axisymmetric case to 40\% in the helical configuration in [284]. This suggests a modest increase of the bootstrap current. The shaping effect is also related to some increase of the fraction of bootstrap current: in [497] it is evaluated to be of the order of 10\%–30\%, which maintains the total fraction very low.

![Figure 54. Total and bootstrap parallel current profiles calculated for RFX-JBS is multiplied by 10^3. Reprinted from [284], with permission from JSPF.](image)

A more relevant effect is instead expected in a low aspect ratio equilibrium, mainly due to enhanced neoclassical viscosity [501]. More specifically, in [502] it has been found that, in a reactor relevant range of parameters with an aspect ratio 2, ellipticity 1.4 and triangularity 0.4, if the pressure profile is flat and thanks to a decreased magnetic shear, an equilibrium with 94\% bootstrap current and extremely high values of \( \beta = 63\% \) stable to ideal kink and Mercier localized modes could be maintained.

6.4. Summary and possible developments

As indicated by experimental scaling laws, the confinement in the RFP scales favourably with electron density, which makes high-density regimes particularly attractive in this sense. Unfortunately, due to the low particle confinement time the density enhancement (with consequent confinement increase) induced by pellet injection is transient, so that stationary high-density regimes are reached with high rates of gas puffing. This produces edge peaked density profiles, which results in enhanced radiation, decreased temperature and enhanced request for ohmic power to sustain the current. In fact, as discussed in section 6.2, in MST the combination of PPCD with pellet injection allows the achievement of high values of confinement and \( \beta \), also exceeding the Mercier criterion, but in a transient way. On the other hand, in RFX-mod, though when pellet penetrates the helical core in QSH states the particle confinement is enhanced, the effect is anyway transient as well. Future developments could be linked to a more efficient multi-pellet injection, capable of injecting several pellets throughout the plasma discharge, coupled to a low recycling wall (see chapter 8): this could allow the development of more core-peaked profiles, avoiding the MARFE-like structures at the edge while maintaining higher \( n_e T_e \) product in the core.

7. Active control of MHD stability

7.1. Resistive wall modes

In the past, the limitation to pulse length imposed by RWM instabilities was considered a serious issue for the future of the RFP configuration. In order to characterize the dynamics
of RWMs, several RFP experiments were fitted with resistive walls, sometimes referred to as thin shells, with wall time constants shorter than the discharge pulse length. These early RFP experiments with resistive walls were commissioned in the late 1980s and included the OHTE device at General Atomic [503, 504], the Reversatron at the University of Colorado [505], the HBTX-1C device at Culham [506] and REPUTE-1 [507, 508].

A convincing experimental demonstration of the RWM instability was achieved in the HBTX-1C experiment. The wall time constant \( \tau_w = \mu_w \delta \sigma \) (where \( \delta \) is the minor radius, \( \sigma \) is the wall conductivity) for HBTX1C was 1 ms. The discharge was sustained for about 4 ms. Nonresonant ideal kinks (RWMs) were observed with growth rates about a factor of three times the wall time constant, consistent with theoretical studies assuming typical current profiles. Even though REPUTE-1 was affected by error fields, it detected a marginally nonresonant mode consistent with the stability limit of an ideal internal \( m = 1 \) kink mode [507]. The OHTE experiment did not report evidences of RWMs [504], while the EXTRAP-T2R experiment (based on OHTE but without helical coils), rebuilt in 1999–2000 with a resistive wall, actually did show clear RWMs [509].

The wall penetration time for EXTRAP-T2R was selected to be 12 ms, which was relatively long compared to the earlier resistive shell experiments in OHTE; the observed growth rates of the RWMs were on the time scale as the wall penetration time: the normalized RWM growth rates spectrum \( \dot{\gamma}^{m,n} = \tau_w \gamma^{m,n} \) for a typical RFP EXTRAP-T2R equilibrium are shown in figure 55.

Shortly after the experimental evidence for the existence of RWMs the first attempt of a simple active feedback experiment was performed in HBTX-1C [510]. Two helically wound coils (sine and cosine configuration) outside the outer shell of the vessel were incorporated in order to control the \( m = 1, n = 2 \) mode. When the mode amplitude reached a pre-set value (below the level corresponding to discharge termination), the helical coil with the correct phase was activated to suppress the growth. With this feedback algorithm, the targeted mode was suppressed below the pre-set level. Although the early demonstration of active feedback of a single RWM in HBTX-1C was noteworthy, suppression of only one RWM mode did not substantially improve the RFP performance and the extension to multiple modes would have been technically challenging.

7.1.1. Theoretical studies on RWM stabilization. Theoretical studies of RWM instability properties were carried out to determine the requirements for active control. In [511] the effect of shell resistivity and distance of the boundary on linear stability of RWMs was investigated. Analytical studies developed further the stability analysis of the RWMs, taking into account rotational stabilization and wall distance [512, 513], and supported the conclusion that active feedback was necessary for long pulse operation in the RFP.

The DEBS code was modified to include resistive wall boundary conditions and investigate nonlinear aspects of RWM stability. Nonlinear numerical simulations of an RFP with the RFX geometry performed by this code provided estimates of performance degradation due to an increase in fluctuation amplitudes [514], both due to RWMs and tearing modes, on time scales longer than the shell time. The code was further modified adding a ‘virtual’ set of helical coils in order to perform multiple RWM stabilization studies [515]. Each virtual helical coil was independently acting on a different Fourier harmonic, a feedback algorithm named mode control (MC), because each mode was independently controlled with a different gain. This algorithm has been implemented and tested experimentally, as it will be described in section 7.1.3. An extensive scan of parameters for DEBS runs showed that the external kinks are always unstable if the wall surrounding the plasma is not ideal. The instability was found independently of the initial condition, the plasma equilibrium or the plasma physical conditions.

A significant finding was that the external kinks are well described within the linear theory: their nonlinear interactions are, in fact, almost negligible during their growth phase. More generally, linear theory gives a good prediction for the \( m = 1 \) spectrum, eigenfunctions and growth rates. As far as feedback control is concerned, feedback stabilization of external kink modes is possible, if carefully chosen gains are simultaneously set on 4 or 5 unstable modes. An experimental verification of the linear RWM behaviour was obtained in EXTRAP-T2R [509], further experiments performed in RFX-mod [516] confirmed the linear behaviour of the RWMs.

Different from tokamaks [517], the stability of internally and externally non-resonant modes in RFPs is not affected by plasma rotation, except if the plasma rotation velocity becomes comparable to the Alfven velocity [513, 518]; the latter is larger than the natural rotation velocity in current RFPs. Hence, based on the above linear theory active feedback is definitely needed for a long pulse RFP. Despite these differences, studies of resistive wall instability physics for RFPs and for tokamaks are complementary, and the physical model for MHD stability for both configurations with a thin (resistive) wall has many similarities [519].

7.1.2. The intelligent shell geometry. As an alternative to helical coils, a concept for feedback stabilization of the RWM instabilities called the intelligent shell (IS) or virtual shell (VS)
was proposed by Bishop [520] at about the same time as the experiments on HBTX-1C were carried out. The IS is composed of a set of saddle coils that surround the plasma. The amount of current driven in each coil is determined by sensors located at the same position as the coil and with an identical shape. From a control theory point of view, this scheme consisted of a series of independent single input single output circuits. The aim of the IS concept is to ‘freeze’ the magnetic flux linking the powered coil, i.e. the system of coils mimics eddy currents and behaves as a perfectly conducting shell. The array linking the powered coil, i.e. the system of coils mimics eddy circuits. The aim of the IS concept is to ‘freeze’ the magnetic flux located at the same position as the coil and with an identical amount of current driven in each coil is determined by sensors for all \((m, n)\) except the \((1, 0)\) coefficient, which is instead set to 0. SVS was also used to let selected unstable \((m, n)\) modes grow uncontrolled while keeping all others at negligible levels [77, 526].

7.1.3. The mode control algorithms. Based on nonlinear feedback studies recalled in previous sections, a feedback system with a realistic geometrical description of the coils and of the sensors was studied in [527], by applying standard control theory techniques as previously adopted for tokamaks [528]. The geometry of the IS was adopted (i.e. a finite and equal number of sensors and active coils, both in poloidal and toroidal directions), but the control scheme was different. In this scheme (dubbed mode control (MC)), the feedback variables \(b^{m,n}\) are the Fourier coefficients of the flux loop sensor signals \(b^j\), while the current harmonics \(I^{m,n}\) correspond to the coefficients of the Fourier decomposition of the currents flowing into the control coils. The current flowing on each control coil is determined by an inverse Fourier transform of the spectrum of \(I^{m,n}\) current harmonics, which is proportional to the corresponding feedback variable. This mode of operation was also defined targeted-mode control [525]. Due to the linearity of the Fourier transform, the mode control feedback scheme is equivalent to the IS if the same gains are set for all modes in mode control or for all positions in IS. In mode control, a different gain for each mode can be applied, and in particular, it can be set to 0 in order to let the corresponding harmonic uncontrolled. As Fourier coefficients are complex numbers, mode control gains can be complex. Experiments have been performed with the proportional component set to complex values [76, 516]. This technique has been used to induce slow rotations of MHD modes (both TM and RWMs).

7.1.4. Influence of control coil field sidebands and their aliasing on MHD control. Due to its finite number, whenever the control coil system generates an \((m, n)\) harmonic aimed at cancelling a corresponding unstable plasma mode, it also creates sidebands with helicity \((m + kM_c, n + hN_c)\) where \(k\) and \(h\) are integer numbers. If the number of sensors is equal to the number of coils (as in the IS geometry) all of the sidebands are aliased into the \((m, n)\) coefficient of the discrete Fourier transform of the measurements \(b^{m,n}\). In other words, measured harmonics are polluted by the field generated by control coils. Equivalently, the perfect cancelation by the feedback coils of the sensor fluxes does not correspond to the cancelation of the plasma or error field harmonics. This effect was simulated for RWM control in [521] and in [527] and determined the optimal number of actuator and sensor coils. Both studies indicated that, when using a control system with the same number of sensors and actuators (i.e. with an IS geometry), the first sideband should not correspond to an unstable mode. This determines the choice of the number of coils. Given that the RWM spectrum of poloidal modes in the RFP corresponds to \(m = -1, 0\) and 1, at least \(M_c = 3\) coils in poloidal direction are necessary, but practical mechanical considerations suggest adopting \(M_c = 4\). As far as the toroidal number \(N_c\) is concerned, for the aspect ratio \(R/a = 4\), as in RFX, the minimum number of coils is 24, in order to avoid the generation of a sideband in the tearing mode part of the spectrum while stabilizing an RWM. In RFX-mod, in order to increase the stable operation window
A different approach to deal with sidebands aliasing, dubbed clean mode control (CMC) [529] and developed for RFX-mod in the context of the control of tearing modes (see next section 7.2), was successfully applied to the current driven $q(a) < 2$ tokamak RWM [530].

7.15. Experimental demonstrations of RWM stabilization. Two RFP experiments, EXTRAP-T2R in Stockholm and RFX-mod in Padua, were specifically rebuilt to provide platforms to address experimental studies of active feedback to suppress RWM instabilities. RFX-mod was equipped with a thin copper shell (with $\tau_w = 100$ ms), which replaced the thick aluminum shell of RFX, while the copper shell wall time constant of EXTRAP-T2R was $\tau_w = 12$ ms. As far as feedback was concerned, both experiments implemented the IS geometry: EXTRAP-T2R (figure 56) had $M_c = 4$, $N_c = 32$ saddle coils and corresponding saddle loop sensors, while RFX-mod was equipped with a matrix of $M_c = 4$, $N_c = 48$ (figure 57) active coils and sensors.

First experiments on simultaneous stabilization of multiple RWMs in the RFP have been performed with a limited number of active coils in EXTRAP-T2R, i.e. $N_c = 16$, $M_c = 4$ [522], i.e. half of the installed ones (figure 58). In order to reduce power supply requirements, four coils at each toroidal position were hard-wired into two pairs forming $m = 1$ active coils. The sensor coil array was $(N_s = 32, M_c = 4)$ and the four sensor coils at each toroidal position were also hard-wired forming $m = 1$ sensors [525]. With this configuration, mode coupling of $n$ and $n + 16$ modes due to the sidebands generated by the coil array could be partially resolved by the sensors coils, although not controlled by the active coils. The IS experiments clearly illustrated the effect of the violation of the stabilization criterion $N_c > \Delta n$ derived in [521]: the flux through all the sensors could be maintained at zero although two unstable modes, separated by $\Delta n = 16$, were not zero but with phases aligned to compensate for each other at actuator locations.

After the upgrade of the coil power supply system to $M_c = 4$, $N_c = 32$, complete stabilization of the unstable spectrum of RWMs was achieved in EXTRAP-T2R [531, 532] With the optimized feedback parameters, pulse lengths of 90 ms, limited by power supplies [533], were achieved corresponding to about ten resistive wall times (figure 58).

The RFX-mod device resumed operations after the substantial modifications to RFX in 2005 [7] and it also reported clear evidences of unstable RWMs when operating without feedback [77].

Since the beginning, in RFX-mod feedback operations exploited the full coverage by the saddle coil array with $N_c = 48$ and $M_c = 4$, 192 independently controlled coils. The same hardware and the same control algorithms developed for the EXTRAP-T2R device were adopted. Initial operations, based on the selective VS approach, allowed extending the RFX-mod discharge duration from 150 ms to more than 300 ms, with a significant loop voltage reduction, as reported in [85]. Different from EXTRAP-T2R, RFX-mod (as previously RFX) was affected by the wall-locking of tearing modes (as will be described in next section 7.2): therefore, the feedback system was also limiting the radial field penetration of wall locked tearing modes.

RFX-mod has also been operated as a circular ohmic tokamak with toroidal field of 0.5 T. Its particular combination of the copper shell and of the Inconel vacuum vessel geometry is able to convert the $m = 2$, $n = 1$ external ideal current-driven kink that always occurs when $q(a) < 2$ into a RWM. The number of actuator coils was not designed to satisfy the stabilization criteria of this mode: in fact, the standard mode control scheme that was successful for RFP RWMs proved insufficient for the $q(a) < 2$ RWM in the tokamak [530], no matter what
gain was used. A simple stabilization criterion for standard mode control has been derived in [534]. Once the clean mode control feedback algorithm was put in place (see section 7.2.3), the RWM was routinely kept at a negligible level for the whole duration of the discharge.

7.1.6. Advanced topics in RWM control. The demonstration of feedback stabilization of RWM instabilities was very important for the RFP concept. After RWM stabilization became routine for the RFX-mod and EXTRAP-T2R RFP experiments, the comprehensive and flexible control systems have been used for a range of additional technology and physics developments of general interest for RWM physics in fusion experiments.

Resonant field amplification (a.k.a. plasma response): the capability of resolving and targeting modes using the control system opened up the possibility to study RFA [535]. A broad spectrum of field errors is ubiquitous and field error harmonics are present in the range of the RWM spectrum. In particular, it was experimentally observed that the spectrum of RWMs, both in amplitude growth rate and in phase, was very reproducible when active control was not operating. This implied the existence of a reproducible field error spectrum, acting as a reproducible seed to the RWM dynamics. The observed growth rate spectrum was not as theoretically predicted: in particular, low $n$ modes $n = \pm 1, \pm 2$, which according to theory were marginally stable, had high growth rates. The field error spectrum contains low $n$ modes associated with shell gaps, leakage flux from the iron core or lack of symmetry in toroidal coil positioning, which should result in RFA. Because the mode growth and phase were reproducible, separate modes could be targeted using an open loop configuration. Transient external control field harmonics were produced and the response of the plasma–shell system was analysed. The growth and the damping rate of modes could be distinguished from the RFA effects and evaluated, giving good agreement with the theoretical predictions.

An important issue in all magnetic configurations is the assessment of error fields during plasma operation, which can be significantly different from error fields measured in vacuum conditions [536]. This was fulfilled by using the reference output tracking capability to produce a simulated external static field error harmonic with controlled amplitude and phase and, simultaneously, the same harmonic rotating with controlled amplitude and phase. The effect of the rotating harmonic on the static field error harmonic depends on the relative phase difference. As the phase of the rotating harmonic passes the phase of the static harmonic, growth and suppression produce an oscillation in the measured amplitude, which can be used to estimate the phase and amplitude of the field error harmonic.

RWM control with partial coverage: a general issue is the development of a control system when the number and placement of control coils is limited for technical reasons. The RFX-mod and EXTRAP-T2R experiments have full active coil coverage but the reduction of coverage can be simulated in the control algorithm software. Furthermore, both devices have well developed control models, so that extensive studies of partial or few-coils coverage can be performed numerically. In RFX-mod, control with reduced coil coverage in the RFP was studied and compared with simulations using the flight simulator with encouraging results [537, 538]. Similar studies were performed in RFX-mod tokamak discharges: as low as six active coils located in the outboard midplane were sufficient to stabilize the current driven $m = 2, n = 1$ unstable mode that occurs when $q(a)$ decreases below 2 [537]. In EXTRAP-T2R an attempt to address RWM stabilization with few coils and low coverage was made incorporating the empirically derived model-based control to simulate the control with reduced coverage [539]. Full control was not achieved in these first attempts, though the capability for studies was demonstrated.

Modelling of the control system: an example of the integration of the control system in the demanding technology of a fusion device has been performed at RFX-mod [540], by the adaptation of the CarMa code [541]. This allowed to study the influence on the control system performance of the 3D passive structures and the RFP plasma using the finite element electromagnetic code Cariddi and a 2D model for RFP plasma stability (MARS-F) [541, 542]. In this way, the MHD stability analysis includes the effects of the 3D conducting structures on the control fields. Alternatively to this model-based approach (white box), empirical data have been experimentally collected and a model created for the case without plasma (black box model) [543]. A further advancement in the capability to simulate experimental conditions was the development of a full dynamic flight simulator which includes the features of the CarMa code and a model of the RFX-mod control system [538].

Advanced system identification methods (SIM): the characterization of the RWM dynamics has been performed either theoretically by MHD modelling, or experimentally by observing the mode evolution without feedback control. An alternative approach has been actively developed in EXTRAP-T2R: the RWM dynamics is probed during the feedback-stabilized operations in the real system by simultaneously and exciting the RWM spectrum with pseudo-random spatial patterns (dithering). A minimal physics-based linear model is assumed to represent the data while all parameters (such as the RWM growth rates and the vacuum field penetration time) are extracted from the measurements [544, 545]. The model includes also a prescribed number of aliased sidebands. An important aspect of this kind of technique is that it allows constructing a model that captures the dynamics of the plasma system, including the real 3D wall effects. At the same time, it naturally gives a reasonably small state space [546] representation that can be implemented in real time. An example of the estimated RWM spectrum is shown in figure 59.

The use of closed-loop system identification techniques to enable experimental modal analysis derived solely from empirical data was extended in an experimental study in RFX-mod [547]. Contrary to EXTRAP-T2R, the stabilization diagram for RFX-mod showed unstable activity in the eigenvalue range corresponding to tearing modes. This could be attributed to
Figure 59. Empirical RWM dispersion relation for T2R. Red colour indicates high modal density (highly probable location of the plasma response eigenvalues; \( z < 1 \) indicates stability, \( z > 1 \) indicates instability. Reproduced courtesy of IAEA. Figure from [545]. Copyright 2013 IAEA.

The fact that tearing modes are typically not rotating in RFX-mod. Therefore, the corresponding radial fields (that are well described by a nonlinear dynamic model, as will be shown in next section) are detected by the sensor system and negatively affect the parameters estimate of the linear control model.

The estimated SIM model enables the synthesis of a controller that implements the multiple input multiple output capability of the system. A further refinement based on modern control theory is the development of a model predictive control (MPC) methodology which employs the SIM model for prediction and computes in real time optimal control inputs that satisfy a performance criterion [548]. Since the control model is updated in real time, MPC could have advantages for long pulse operation where discharge parameters, such as equilibrium profiles, change in time.

7.2. Control of tearing modes in the RFP

The issue of wall and phase locking of tearing modes (TM) has been intensively addressed both theoretically and experimentally. In particular, several mitigation approaches have been developed, which are described below. Once major error fields are reduced by careful design or by active coils, the ultimate cause for wall-locking of TM is the conductivity of the nearest closed shell facing the plasma. Based on the experience gained in the RFP devices and in RFX-mod in particular, a modification of the RFX-mod device (named RFX-mod2) is being implemented, as will be described at the end of this chapter.

7.2.1. Wall and phase locking of tearing modes. A basic quasi-linear phenomenological analytical model of the mechanism of wall locking of a single rotating tearing mode was described in [549] for the case of tokamak geometry. An adaptation for the \( m = 1 \) modes in the RFP has been described in [550, 551], in order to explain the different plasma current threshold for the observed wall-locking of TMs in MST and RFX.

The model assumes that the phase dynamics of a single TM depends on the balance between the electromagnetic torque, produced by the eddy currents on the first resistive conducting wall surrounding the plasma, and a phenomenological viscous torque tending to restore the natural rotation frequency of a tearing mode when it is not wall locked.

With a passive boundary, a bifurcated behaviour occurs, as shown in figure 60. The locking of the mode phase to the wall occurs above a certain mode amplitude while unlocking is found at a significantly lower threshold. The absolute value of the mode amplitude for the locking threshold depends on the resistivity of the closest continuous conducting structure surrounding the plasma. In fact, even though both RFX and MST were equipped with a thick aluminium stabilizing shell, RFX plasma was contained by a vacuum vessel made by Inconel 625 alloy [6]. The significantly higher vessel resistivity was at the origin of a lower mode amplitude threshold for wall locking (short dashed line in figure 60). The behaviour of wall locking in the TPE-RX device was intermediate between RFX and MST. In fact, despite its carefully designed magnetic front-end, with an external thick aluminium shell and an internal thin copper shell, the plasma was contained by an AISI 316L stainless steel vacuum vessel whose resistivity is half than Inconel [10].

Error fields play an important role in the tearing modes wall locking process [552]. They are non-axisymmetric magnetic fields generated, e.g., by eddy currents induced in the non-homogeneous conductors surrounding the plasma due to the gaps in the shells (necessary to allow the penetration of the vertical field) [553] or by a static misalignment of external coils. When the component of an applied error field resonant with tearing modes is above a certain threshold, wall locking is forced because tearing mode braking due to eddy currents
is more effective when error fields are present [554, 555]. In RFX-mod, EXTRAP-T2R and MST, thanks to active control coils it was possible to study the braking of plasma rotation when controlled external perturbations are applied [556]. In particular it was possible to measure the dependence of the perpendicular viscosity on the magnetic fluctuation amplitude [557] and compare it with the predicted rate of stochastic field line diffusion [558].

The locked modes phenomenon (introduced in section 2.2.1) i.e. the phase alignment among different TMs, has been modelled by a quasi-linear approach in terms of electromagnetic torques due to three wave nonlinear coupling among \( m = 1 \) and \( m = 0 \) tearing modes [552, 554, 559, 560]. The aim of the analytical model was to investigate the formation and breakup of the LM in the RFP, in order to complement 3D non-linear visco-resistive MHD simulations. The electromagnetic torque in the region near a given resonant surface is ascribed to the modification of the linear \( \delta J \) and \( \delta B \) caused by the three-wave nonlinear coupling among \((1, n), (1, n+1)\) and \((0, 1)\) modes. In the analysis, the \( m = 1 \) modes are assumed to be saturated at some level: therefore, their amplitude is considered a fixed parameter. Instead, the \( m = 0 \) modes are taken as intrinsically stable and only driven by the nonlinear interaction of the \( m = 1 \) modes. The model predicts that when the amplitude of two coupled tearing modes exceeds a critical value (locking threshold) the two modes phase-lock, while when they are below a much lower critical value (unlocking threshold) their phases are not coupled anymore. In between the two thresholds, both locked and unlocked solutions can occur, depending on the initial conditions.

A thorough analysis of RFX magnetic measurements to characterize the tearing mode spectrum along with the experimental features of the phase-locking phenomenon was documented in [75] and compared with both the Fitzpatrick analytical theory and nonlinear numerical visco-resistive simulations. A partial agreement between the analytical theory and experimental data was found: the residual discrepancy was ascribed to the \( m = 0 \) modes, which were linearly unstable in RFX due to the large shell/plasma distance \( b/a \approx 1.18 \).

A refined theory of mode locking was proposed in [554]: at high mode amplitudes the various tearing modes phase-lock together minimizing the magnitude of the electromagnetic torques exerted at the mode rational surfaces. In particular, the relative phase of \( m = 0 \) modes was assumed to remain constant in time at the value observed experimentally and that was found to correspond to the minimum absolute value of the electromagnetic torque.

### 7.2.2. Mitigation techniques of the locked mode.

In RFX [561, 562] and in MST [117], the intrinsic error fields were reduced by analog magnetic feedback systems using local saddle coils. Passive reduction of error fields, by means of an overlap of the shell edges, was implemented in TPE-RX [10] and in RFX-mod [563].

In RFX the active control of the position of the wall locked tearing modes was obtained by means of an external non axi-symmetric \( m = 0 \) magnetic field, obtained by modulating the currents in the toroidal field coils [423], thanks to the non-linear coupling between the \( m = 0 \) and the \( m = 1 \) TMs [554].

In RFX-mod, the active feedback by saddle coils allowed a significant mitigation the effects of wall locked tearing modes.

A first improvement compared to RFX and to passive RFX-mod operations [77], was obtained during initial operations with the VS algorithm [76]. Nonetheless a significant toroidally localized PWI was still occurring [79], due to the unrecognized sidebands aliasing described in section 7.1.4. Several forced rotation algorithms were therefore implemented [76], taking advantage of the direct coupling of each \((m = 1, n)\) TM with the corresponding field harmonic generated by the control coils. Slow tearing mode rotations up to 40 Hz were observed, together with a slow rotation of the PWI region. Due to the sidebands’ aliasing systematic error (section 7.1.4), the VS control algorithm only partially reduces the edge amplitudes of tearing modes, as illustrated by the Fourier spectra shown in figure 61. In a VS controlled discharge, in fact, the Fourier spectrum of radial sensor measurements is significantly low, but these harmonics do not correspond to the real amplitude of tearing modes as they are polluted by coils sidebands. Red bars show instead the clean harmonics corresponding to the real amplitudes of tearing modes, which are much higher and explain the observed PWI. The clean harmonics are obtained by a de-aliasing algorithm: the measured currents flowing into the control coils are used in a vacuum cylindrical model to obtain the time behaviour of the relevant sidebands to be subtracted from the raw harmonics (details are described in [529]).

### 7.2.3. Spontaneous rotation of tearing modes with clean mode control.

A second significant improvement of RFX-mod discharges was obtained by the implementation in real-time of the clean mode control de-aliasing algorithm [118] and allowed RFX-mod to reach the design current of 2 MA.

The correction of the aliasing error revealed that control of tearing modes exhibit a non-linear behaviour: with
clean mode control, it is possible to reduce the edge amplitude of each wall-locked tearing mode down to a minimum value, by properly increasing the corresponding proportional gain (figure 62(b)) [529]. Any further increase of gains generates a rotation of increasing frequency (figure 62(a)). On the other hand, active feedback with clean mode control was found to remove the bifurcated behaviour of the wall locking of tearing modes as described in section 7.2.1. During optimization experiments of He discharges with very low plasma current \( I_p < 150 \text{ kA} \) a regime that was never investigated in RFX the spontaneous rotation of tearing modes was observed (figure 63) [564]. On the one hand, such a low wall-locking threshold was expected, due to the high resistivity of Inconel vessel. On the other hand, a new unexpected feature of these RFX-mod experiments was the observation that the natural rotation locking threshold correspond to the unlocking one. As in RFX-mod tearing mode amplitudes scale almost linearly with plasma current \( I_p \), the reversibility of the locking threshold with plasma current reflects a reversibility with mode amplitude: this is a marked difference with respect to the hysteresis observed in all other passive RFP experiments.

7.2.4. Modelling of tearing mode phase dynamics under feedback control. The rotation of tearing modes subject to clean mode control was initially reproduced by a stationary nonlinear model [529], based on the same physics mechanisms describing the single TM wall locking in RFX as in [554]. The active control coils are modelled as a harmonic error field whose amplitude and phase are determined by a prescribed feedback law. By assuming a time-constant mode amplitude at the resonant surface, above the wall-locking threshold, in the absence of feedback the mode would be strictly wall-locked. Instead, by applying feedback, the model shows the existence of equilibrium solutions, corresponding to uniform rotations of the mode with frequencies larger than \( 1/\tau_w \) for gains above a threshold value. These rotations establish with low edge amplitude: this is the advantage of the feedback with respect to feed-forward techniques to induce LM rotations.

Subsequently, a time dependent code including nonlinear phase coupling among multiple \( m = 1 \) tearing modes, dubbed RFXLocking, was developed [565]. In this code the \( m = 0 \) low \( n \) modes are assumed to be always phase locked, consistent with the experimental RFX findings. As in previous wall-locking studies, the model evolves the mode edge amplitudes and phases under the action of the feedback coils, while the mode amplitudes at the resonant surface are estimated from the experimental measurements.

Thanks to a modified version of RFXLocking, implementing the multiple-shell structure of RFX-mod, a model-based optimization has been performed [566], in order to find an optimal gain set and reduce the edge radial magnetic field amplitude of each tearing mode to the lowest possible value compatible with the tearing mode dynamics.

RFXLocking also reproduces rather well the tearing mode dynamics as a function of the feedback gains. The effect of the proportional gain on the \((m = 1, n = -7)\) mode angular frequency, the normalized edge radial magnetic field and the coil current is shown, for example, in figure 62. The experimental data (full circles) are compared with the model prediction for the same type of discharges (red dashed line). Each point represents a time average over the discharge flattop. A reduction in the total edge radial magnetic field has been obtained with the model-based set, with respect to the empirically optimized one.

The RFX-mod CMC experiment has shown some limitations in the possibility of reducing the edge radial field even when using the radial field harmonics extrapolated at the plasma surface [118] as the feedback variable. The edge radial field cannot be reduced below a minimum non-zero value for

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**Figure 62.** Effect of a proportional gain scan on (a) the \( m = 1, n = -7 \) mode angular frequency, (b) the \((1, -7)\) normalized edge radial magnetic field and (c) the \( m = 1, n = -7 \) coil current amplitude. Each point represents a flattop average and the experimental points are from a set of 1.4 MA reproducible discharges. The error bars are 25th and the 75th percentiles of the data. They represent the variation of the above quantities due to the natural mode dynamics, not measurement errors. The red dashed curve represents the RFXLocking core prediction. Reproduced courtesy of IAEA. Figure from [566]. Copyright 2010 IAEA.
each tearing mode. This limitation is not related to the coils’ amplifiers, but is physically based, as it can be reproduced by the RFXLocking code that shows the existence of a minimum radial field in an optimum gain region [534]. More precisely, simulations with RFXLocking showed that the ultimate CMC possibilities are fixed by the plasma–shell proximity: no matter the radius where the virtual sensors are located, the actively controlled radial field is always higher than the one obtained by an ideal shell that replaces the resistive one.

7.3. Future perspectives: the design of RFX-mod2

Based on the aforementioned understanding of the interplay between passive conductive boundaries and tearing modes in an RFP, an upgrade of the RFX-mod machine assembly has been designed, dubbed RFX-mod2, and it is now being implemented [19–21, 114]. Main aim of such upgrade is to modify the resistivity and proximity of the magnetic front-end. This will allow reducing the saturated level of tearing modes, and subsequent magnetic chaos, due to increased proximity of the shell to the plasma. Moreover, it will decrease the non-axisymmetric deformation of the last closed magnetic surface (LCMS), due to the phase locked tearing modes. Finally, the plasma current threshold for the wall locking will increase, allowing a better plasma start-up phase.

In RFX-mod2, the stainless steel vacuum vessel containing the RFX-mod plasma will be removed, and the copper shell will become the first conductive surface for the plasma. The vacuum barrier will be provided by the properly modified toroidal support structure. The shell/plasma proximity will be reduced from $b/a = 1.11$ to $b/a = 1.04$. This design choice implies challenging modifications of the components of the machine close to the plasma, while maintaining the existing magnetic coil system. Innovative solutions have been conceived to fulfil vacuum and electrical requirements of the in-vessel components, in particular to ensure their reliability during the operation also in presence of abnormal events.

According to model simulations, this magnetic configuration will reduce the deformation of the LCMS by a factor about 2–3 with respect to RFX-mod, see figure 64 [347]. Moreover, the 3D nonlinear MHD Specyl, including realistic boundary conditions, predicts a reduction of the energy of unstable modes by about 30% [114].

8. Plasma–wall interactions

As in other magnetic confinement devices, in the RFP the PWI plays a crucial role in determining plasma behaviour and performance through the control of plasma density and impurity contamination. In RFPs, PWI is strictly dependent on the magnetic topology, being strongly influenced by the three-dimensional edge magnetic fields related to the wall locking of tearing modes also define locked mode (LM) that produces a localised deformation of the LCMS. When such bulge of the LCMS is wall locked, a large power is deposited onto a relatively small region of the wall, with uncontrolled release of main fuel and impurity particles.

8.1. Effect of magnetic topology on plasma–wall interaction

The LM drives large localized power fluxes onto the wall in the region intersected by the distorted magnetic surfaces. An example of the footprint on the first wall in an RFX discharge is shown in figure 65, compared with the reconstruction of the plasma surface deformation.

In RFX (and in RFX-mod afterwards), the plasma wall is fully covered by polycrystalline graphite tiles, to protect the Inconel vacuum vessel and avoid metal contamination [6]. Though local field errors introduced by the insulating poloidal gap (allowing for the penetration of the magnetic field) were reduced by a local active control system, LMs still produced a
large radial displacement of $\approx 2$ cm [567], determining power loads locally as high as $\approx 100$ MW m$^{-2}$ in the region where the bulge intercepted the wall [80]. This was associated to local enhancement of electron density and density fluctuations, suggesting a continuous ionization and loss of neutrals, and to increased transport and impurity influxes [568]. Carbon and oxygen influxes from the LM region were estimated to be of the order of 20%-40% of the total; however, tomographic measurements of total radiation showed that, due to the relatively small region of interaction, the radiation enhancement produced by LM was a fraction of the total input power of the order of $\approx 10\%$, leading to the conclusion that the main loss term is due to the local increase of transport and particle fluxes [80]. With plasma current approaching 1 MA, in proximity of LM the graphite tiles reached the carbon radiation enhanced sublimation temperature, producing carbon blooming processes, observed as sudden increase of carbon, radiation and electron density. Such strong particle release locally cooled the plasma, leading to the loss of density control and to the premature decay of the plasma current [78, 569–571]. An example of carbon behaviour when such blooming processes occurred is given in figure 66. A rather crude means to mitigate PWI and carbon blooming was found by applying a toroidally rotating $m = 0$ perturbation, thus forcing a rotation of the bulge associated to LM [423] (section 7.2.2). This technique increased the amplitude of $m = 0$ modes, with associated confinement loss, but the enhanced interaction region was distributed in time over all toroidal angles, allowing reproducible operation up to 1.1 MA. A different approach to mitigate PWI in RFX was adopted by increasing the power radiated at the edge by impurity seeding (in particular neon) [572]. This operational scenario did not affect confinement properties, while both poloidal and toroidal asymmetries of radiated power were reduced, in particular in the high-density regime.

Despite all these issues related to the strong PWI induced by the presence of the LM, in RFX quite low effective charge values were observed, $Z_{\text{eff}} \approx 1.5$ (except for very low density regimes, where the bremsstrahlung measurements become anyway difficult), independent of plasma current. This result has been interpreted as due to a screening mechanism, modelled by a Monte Carlo calculation, associated to a low ratio between impurity ionization length and Larmor radius, and to the presence of a radial electric field which points inward at the wall and outward in the edge plasma [573].

In RFX-mod, TM mitigation techniques, described in section 7.2, allowed obtaining a smoother PWI, compared to RFX. Moreover, the shape of carbon tiles in RFX-mod was optimized in order to avoid strong interactions at their edges [7].

The ratio between the maximum wall temperature increase and the total energy input times the discharge duration $\Delta T_{\text{max}} / (\Delta t_{\text{pulse}} \Delta E_{\text{pulse}})$ has been taken as an indicator of PWI intensity in [79]. This indicator has been found to be correlated with the maximum of the radial field associated to $m = 0$ modes (figure 67), thus implying a dependence of the edge properties on $m = 0$ tearing modes.

Similarly, in the EXTRAP-T2R device before the active stabilization of RWMs, spectroscopy and collector probe measurements showed large metal emissions (Mo and Cr), and arcing was identified as an important contributor to metal
Such effects were mitigated with the application of the active feedback control and allowed a strong improvement of the performance [522]. In the RFX-mod QSH states, the plasma edge is modulated by a helical ripple, thus wetting the wall over a helical region much larger than in the MH state. In such regimes, edge plasma parameters are helically modulated too [208]. In QSH, the power density deposited in the portion of the wall wetted by the helix is ≈5 MW m⁻² also at the highest current [575], and the plasma radiation is globally more symmetric than in MH [576]. On average, in QSH states the locking region accounts for ~18% of the total radiated power, which corresponds to a PWI improvement by at least a factor of three with respect to the MH states and the former RFX [346].

Even in good QSH states, in RFX-mod there is still a residual PWI associated with the phase locking of the secondary modes. Correspondingly, the magnetic connection length to the wall presents a strong decrease (‘hole’) with increased PWI (by a factor ranging between 2 and 10), which is not stationary thanks to the feedback control system. In the RFX-mod2 upgrade it has been calculated that secondary mode locking and associated PWI could be completely avoided [346].

It has to be mentioned that in RFPs the footprint of PWI also shows a poloidal asymmetry between ion and electron drifts, with electron and ion flows separately intercepted by the wall, due to the surface distortion [366, 367, 577, 578]. An example is shown in figure 68. Such asymmetry is more pronounced at low density, indicating that at low density a significant fraction of the poloidal current is due to suprathermal electrons [579]. The separate interception of ion and electron flows at the LM, is at the origin of halo currents which in RFX flowed into the Inconel vessel with intensity up to 25 kA [580], dissipating ≈4%–5% of the total power input.

8.2. Recycling control

In devices with carbon wall as RFX or RFX-mod, the value of plasma density is determined by wall desorption rather than by the fuelling rate. In fact graphite is shot by shot saturated with trapped hydrogen/deuterium, so that the plasma density after the breakdown is completely sustained by atoms released by the wall, with a recycling coefficient R which can be greater than 1 [211, 578]. The number of released particles depends both on the number of atoms stored in the graphite and on the power load on the wall, which makes density control very difficult, especially at high current. On the other hand, in RFPs, in addition to gas puffing, the only central fuelling is by pellet injection (no continuous direct core fuelling by neutral beams is applied), and the particle confinement times are relatively low (of the order of 10 ms, see for example [26]). Therefore, very low recycling regimes cannot be sustained. In summary, a good density control can be obtained with R close to 1 with a not saturated first wall, which is an equilibrium more difficult to achieve in devices with a full graphite wall: indeed, when in EXTRAP-T2R the carbon armour was replaced by molybdenum limiters, the density control improved [9].

In order to guarantee such conditions, in particular avoiding uncontrolled strong release of particles in high current discharges, effective wall conditioning treatments have been applied, including glow discharge cleaning (GDC), pulse discharge cleaning (repetition of short, low energy plasma discharges), operation with hot wall, boronization, lithization. GDC has been applied in devices with graphite wall (EXTRAP-T2, RFX, RFX-mod), mainly with helium as working gas, to remove trapped hydrogen or deuterium; only occasionally GDC in H has been used to remove impurities [581, 582]. In RFX-mod, the effect of such kind of treatment has been characterized [582], showing a toroidal non-uniformity with respect to the positions of the electrodes: the ion flux to the wall features a distribution ±30° wide around the electrodes.

In RFX, to ensure the operation with a non-saturated graphite as plasma facing material, the wall temperature has been raised up to 280 °C in some plasma sessions. A lower effective recycling characterized such pulses, but the wall response was still more unpredictable, with frequent sudden density build-ups and current quenches. This was associated with an easier overheating of graphite tiles in presence of highly localized power loads, up to the carbon sublimation temperature [583]. Such uncontrolled events were instead avoided soon after a boronization. Boronization has been applied with different techniques on several RFP devices, including EXTRAP-T2 ([584]), MST ([585–587]), RFX([588, 589]) and RFX-mod [575]. A common result in all the experiments is that boronization is effective in reducing recycling, for a number of plasma shots that is dependent on the applied technique and boron layer thickness, on plasma regimes and on carbon redeposition rates. The analysis of samples exposed to the boronization treatments showed that a crucial issue to ensure a good behaviour of a boronized wall is to obtain the deposition of a uniform layer [587, 590, 591]. In carbon wall devices where boronization is periodically performed, as RFX and RFX-mod, its lifetime is mainly limited by carbon erosion and redeposition processes, decreasing with increasing power loads [592].

Two examples of the positive effects obtained in MST and RFX are given in figures 69 and 70, respectively. Figure 69 shows how after boronization the same plasma density is sustained by a higher gas puffing [587], while figure 70 displays
a reduction of the plasma effective charge with boronization [573, 593]. In MST, it has also been found that the operation with boronized wall can be associated to discharges featuring prolonged phases (up to 20 ms) without DREs. In such phases, which are believed to be related to a lower plasma resistivity, the energy confinement time increases by a factor 3 [453].

Following the experience of several tokamaks (see for example [594–596]) and stellarators [597], conditioning by lithization has been implemented in RFX-mod by a lithium capillary–pore system [575, 598] and lithium pellet injection [117, 599]. A significant decrease of oxygen and carbon influxes and emissivities (carbon and oxygen are by far the prevailing impurities in RFX) has been found, for carbon in particular more pronounced than with boronization. The important effect was the improvement of density control, which was the main aim of these experiments and allowed operation at high current (≥1.5 MA) with higher density, up to $n_{\text{eff}} \approx 0.5$: an example is shown in figure 71. Despite a slight peaking of density profiles, confinement was not significantly improved.

Also in the case of lithization, a main issue remains the uniformity of Li deposition, though a more uniform layer has been obtained by multi-pellet injection [599]. An additional operational issue is related to the fact that, when the machine is vented and opened for maintenance, the formation of Li$_2$CO$_3$ (difficult to remove in a device lacking of a direct access to the first wall) must be prevented by a long boronization with argon flushing and careful minimization of the opening time [575].

8.3. Dust particles

Dust particles (small loose particles agglomerates, sized from tens of nanometres to millimetre) can be an issue in fusion plasma devices, as they potentially represent a safety hazard, an impurity source and a cause of surface damaging. In this respect, their behaviour has been extensively studied in tokamaks (see for example [600]), while only few papers addressing this issue and related to the EXTRAP-T2 and T2R device have been produced. Main results of such experimental and modelling studies can be summarized as follows:

(a) The preferential origin of dust formation is related again to the high power load deposited in the LM region, where small agglomerates are ejected. In EXTRAP-T2R, in discharges without active feedback more dust particles have been captured (by silica aerogel collectors) than in discharges with active feedback [601, 602];

(b) Post-mortem measurements indicate that most of dust particles are spheroids covered by deposits with granular morphology (figures 72(b) and (c)) and coated by corrugated layers containing Mo and stainless steel components, indicating that re-mobilization is an important element of intrinsic metal dust dynamics [603];

(c) Some particles elongated and aligned with plasma flow have been detected, (less than 20% of the total, see figure 72(a)) suggesting a relevant effect of ion drag [602, 603].
8.4. The problem of the power exhaust and future directions

Due to the necessity of a close-fitting conducting wall, the RFP devices presently in operation are not equipped with divertors. In the past some studies and experiments on small size devices presently in operation are not equipped with effective MHD control systems, showed sometimes some positive effects (decrease of light impurities), but also significant worsening of plasma performance, with increased fluctuations at the X-point and early termination of the discharge.

As an alternative to divertors, the vented pump limiter has been proposed for the RFP as a means for actively controlling neutral particles at the wall [589]. The concept of the vented pump limiter, as applied in the Tore Supra tokamak [609], is to provide a limiter head designed in such a way to be transparent (slots between blades) to recycled neutrals, which enter the limiter and can be extracted. A system of this kind has never been experimentally applied to RFP devices so far. The weak point was the capability of the blades in sustaining power loads that can be locally very high. Another important drawback is that the edge topology is not stationary in time neither uniform along the poloidal and toroidal angles, which indeed makes any tool designed for a tokamak of difficult application in an RFP.

A more viable solution seems to exploit the similarity of the 3D field structures of RFPs and stellarators. In fact, an interesting possibility, proposed in [610], exploits the peculiar topology of the RFP to produce a magnetic configuration resembling the island divertor of the stellarator [611]. The concept is based on the chain of magnetic islands that in QSH-SHAx states develop on the equilibrium magnetic field, but the amplitude of the ac ripple is projected to be only a few percent for a reactor plasma that has low electrical resistance. This creates an attractive, steady-state method for current sustainment that maintains the self-driven neoclassical bootstrap current and minimizes auxiliary current drive, and stellarator research aims to minimize or eliminate plasma current at the expense of toroidal symmetry.

A quasi-steady-state method of inductive current sustainment, called OFCD, has been identified that creates an ohmic, steady-state sustainment scenario for the RFP. In OFCD, audio-range oscillatory toroidal and poloidal loop voltages are applied to the plasma, and the plasma’s nonlinear self-organization process generates a dc current. Accumulation of magnetizing flux is avoided, and the plasma current can be sustained indefinitely since the loop voltages have zero time-average. A remnant inductive ac ripple appears in the equilibrium magnetic field, but the amplitude of the ac ripple is projected to be only a few percent for a reactor plasma that has low electrical resistance. This creates an attractive, steady-state method for current sustainment that maintains the simplicity, robustness, and high efficiency associated with inductive current drive. The TITAN fusion power plant study employed OFCD to attain a steady-state, compact fusion power core based on the RFP (see chapter 10).

The theoretical basis for OFCD has greatly matured as part of the development of nonlinear visco-resistive MHD modelling of RFP plasmas. The physics basis and modelling specific to OFCD are reviewed in sections 9.1 and 9.2. Several
OFCD experiments have been conducted on RFP plasmas that are consistent with theoretical and modelling expectations, which are reviewed in section 9.3. Unfortunately, a demonstration of full current drive by OFCD is not expected to be possible in present day RFP devices due to the scaling for the ac ripple that sets a limit on the equilibrium excursion during an OFCD cycle. Plasmas with larger Lundquist number, i.e., larger size and temperature, are necessary to test full current drive by OFCD.

There have been a couple of experimental studies of OFCD on tokamak plasmas which are not discussed here in detail. The first was performed on the small Caltech tokamak without evidence for current drive [614]. There was also an attempt on the DIII-D tokamak using modulation of the plasma elongation to oscillate the toroidal flux, also without evidence for current drive [615]. Relaxed-state modelling shows OFCD is less effective at high safety factor [616], in addition to the concern that tokamak plasmas might not access instabilities that govern plasma relaxation.

9.1. Introduction to oscillating field current drive

The OFCD current drive method derives from the principle of a relaxed state, where (see chapter 3) the magnetic energy is minimized subject to the constraint of conserved total magnetic helicity, \( K \). The relaxed-state magnetic equilibrium satisfies \( \nabla \times B = \mu B \), with \( \mu \) = constant. Conservation of magnetic helicity motivates a helicity balance, given by \( \partial K / \partial t = 2 V_T \Phi - 2 \int \eta J \cdot B dV \), where \( V_T \) is the toroidal loop voltage at the plasma surface and \( \Phi \) is the toroidal magnetic flux embedded within the plasma. The first term is the rate of helicity injection, and the second term is ohmic dissipation of helicity resulting from the plasma’s finite resistivity, \( \eta \). Conventional induction injects magnetic helicity (and drives current) through a pulse of toroidal loop voltage generated by the transformer that links the torus.

There is a close correspondence between magnetic helicity, plasma current, and the Poynting theorem for electromagnetic energy [617]. Bevir and Gray [618] recognized that, if the toroidal and poloidal loop voltages are sinusoidal, then net cycle-average helicity injection is possible, i.e., \( \partial K / \partial t = (\bar{V}_T \bar{V}_P / \omega) \sin \delta \), where \( \omega \) is the oscillation frequency, \( \delta \) is the relative phase between the loop voltages, and the over-hat identifies the ac amplitude of the toroidal and poloidal loop voltages. Note that the poloidal loop voltage \( V_P = -d\Phi / dr \), so there is a concomitant oscillation in the applied toroidal magnetic field. If the phase, \( \delta = \pi/2 \), then the rate of helicity injection is maximized. This corresponds to driving a net dc plasma current using purely ac loop voltages, hence the terminology ‘ac helicity injection’ or OFCD. Note that phase \( \delta = -\pi/2 \) permits anti-drive, which is used experimentally to induce a current ramp-down and study the dependence on the relative phase. The frequency, \( \omega \), must be faster than the inverse resistive diffusion time, \( \tau_R^{-1} = \eta / \mu_0 a^2 \), but slower than the process that governs magnetic relaxation, i.e., for the RFP, the resistive-Alfvénic timescale \( (1/2\pi R a^2) \) associated with magnetic reconnection from tearing instability, typically in the audio frequency range. At audio frequency, the OFCD inductive current drive has Spitzer conductivity. The principal concern for OFCD may be the impact on energy confinement, since the relaxation process happens through fluctuations. The scaling of these fluctuations is important as it is for conventional pulsed induction.

The formal connection between helicity injection and current drive has been developed in detail [617, 619] and includes dc helicity injection using electrodes which link magnetic flux. This is the basis for coaxial helicity injection (CHI) used extensively for spheromak formation [152, 620] as well as non-inductive start-up of low aspect ratio tokamaks [621, 622]. It also applies to ‘localized helicity injection’ using compact plasma sources inserted into the plasma. This approach was used on MST as a means for current profile control [460, 623] and has been further developed on the PEGASUS low aspect ratio tokamak as an alternative to CHI for tokamak start-up [624]. A version of ac helicity injection termed ‘steady inductive helicity injection’ is also being investigated to sustain spheromak plasmas, which aims to minimize the ac ripple by using loop voltages sources that are not toroidally symmetric [625, 626].

9.2. Modelling and scaling predictions

While nonlinear MHD as discussed in chapters 3 and 5 provides a first-principle description for OFCD, a simplified limit is described first in which a global electromagnetic power balance is applied to a plasma that is assumed to maintain a relaxed-state instantaneously. The 3D dynamics are not resolved, but important effects on the magnetic equilibrium are revealed. This demonstrates OFCD on the basis of energy conservation and makes predictions for the evolution of a relaxed-state plasma. Because this approach resembles circuit analysis, it has also been referred to as zero-dimensional modelling [616, 627, 628].

The distinctive feature of a relaxed-state plasma is a rigid (normalized) parallel current profile, \( \mu_0 I = \mu_0 J \cdot B / B^2 \), independent of varying electrical boundary conditions. Small amplitude fluctuations are assumed to provide a turbulent emf in just the right amount at each radius to maintain the relaxed profile. Any partially relaxed state that is time-space separable as \( \mu_0 (r, t) = \lambda_0(t) \Lambda(r) \) is appropriate for this modelling, where \( \lambda_0 \) is the magnitude of \( \mu_0 \) at \( r = 0 \) and \( \Lambda(r) \) is the spatial form of the rigid profile. The tendency for fixed \( \Lambda(r) \) is the hallmark characteristic of relaxation dynamics in an RFP plasma. For example, the profiles of MST plasmas have \( \Lambda \approx 1 - (r/a)^4 \) over a range of operating conditions. The gradient \( d\Lambda / dr \) is proportional to the free energy for current-driven tearing instability, and dynamically there is a tendency to maintain marginal tearing stability at each instant in time via nonlinear feedback from the dynamo-like emf, \( \langle \nabla \times B \rangle \), in Ohm’s law.

The evolution of a relaxed-state plasma is fully determined by a global magnetic energy balance. The Poynting theorem applied to a cylindrical relaxed-state plasma with radius, \( a \), and length, \( l = 2\pi R \), can be written with dimensionless quantities
impedance is small, i.e., primarily from the plasma’s inductance since the resistive state time-average is attained. This yields a dc plasma current are specified, as well as values for the Lundquist number and frequency, amplitudes, and relative phase of the loop voltages. The boundary conditions include oscillating loop voltages. The \( R \) of RFP plasmas sustained by conventional induction, but the (DEBS code). The formalism is the same as for simulations of RFP plasmas sustained by conventional induction, but the boundary conditions include oscillating loop voltages. The evolution of the plasma current is shown in figure 74 for a case where OFCD is applied to a simulation that had reached steady conditions with conventional induction. When OFCD is initiated, the current settles into a new steady-state after a transition that lasts a couple of cycles. The simulations show that the current profile does not relax instantaneously, as represented by the cycle in \( F - \Theta \) space. Instantaneous relaxation has a nearly linear \( F - \Theta \) trajectory (see [629]). These simulations establish a sound theoretical basis for OFCD dynamics, even if limited by moderate \( S \) such that the magnitude of the ac ripple is large.

Analysis of Ohm’s law provides insight into key differences between current drive with OFCD versus conventional induction. Consider the limit of single-fluid MHD, for which Ohm’s law is \( E + V \times B = \eta J \). At low beta the current profile is governed by the component parallel to the mean-field \( B \), i.e., \( E_{||} = \langle V \times B \rangle_{||} = \eta J_{||} \). The instantaneous alignment between \( E_{0} \) and \( B_{0} \) varies over an OFCD cycle, and the cycle average of parallel Ohm’s law has non-zero \( E_{||} \), despite the loop voltages being oscillatory. An equivalent approach is to treat quantities instantaneously, in which case the OFCD mean-field current drive is represented by \( \langle E_{0} \times B_{0} \rangle = \langle E_{0} \rangle \langle B_{0} \rangle / B \), where the over-hat identifies the driven ac component. Figure 75(a) shows the profile of the OFCD mean-field drive, which peaks in the outer region. For conventional induction, \( E_{||} \) peaks at \( r = 0 \). The turbulent emf averaged over both the mean-field flux surface and OFCD cycle is shown in figure 75(b). For OFCD, the cycle-average relaxation process transports current from the outer region toward the core. This is dynamically opposite to conventional induction, for which current in the core is transported to the outer region of the plasma. Despite this opposite tendency, the shape of the parallel current profile is similar in both cases, consistent with the profile remaining close to marginal stability for current-driven tearing.

9.3. Experimental demonstrations of current drive by OFCD

The first OFCD experiments for RFP plasmas were conducted on the ZT-40M device at LANL in the 1980s [630]. High-
power oscillators based on power triodes were developed that quickly ramped to full amplitude to turn on OFCD rapidly after plasma formation. The frequency was adjustable, $\omega/2\pi = 750$ Hz to 2 kHz, and the relative phase of the oscillators was tunable so that OFCD drive and anti-drive could be tested. Since full sustainment was not possible for ZT-40M’s low $S$ plasma, the background steady induction was fixed, and the change in plasma current associated with the applied ac voltages was monitored. Experiments were conducted for varied plasma current $I_p = 50$–180 kA. The response to drive and anti-drive phasing behaved as predicted, and probe measurements showed that the current profile was relaxed throughout the ac cycle. At 50 kA a small net increase in the plasma current was observed, consistent with ZT-40M’s low $S$ plasma. There were no limiters on the Inconel wall, allowing influx of metal impurities that increased the plasma resistance. At 180 kA, the increased resistance negated net current drive, but the drive and anti-drive phase dependence held at all current levels.

The most extensive experimental tests of OFCD have been performed on the MST device. The minor radius and plasma temperature for MST plasmas are both significantly larger than for ZT-40M, which reduces the plasma resistance. High-power oscillators based on ignitron plasma-closing switches were developed that allowed turn-on at full oscillation amplitude [631]. As for ZT-40M, the background steady induction was fixed, and the change in plasma current was monitored. The oscillator phasing was adjustable to examine the phase dependence in greater detail than in the ZT-40M experiments. Most of the experiments were performed with frequency $\omega/2\pi = 280$ Hz, since the oscillators were based on resonant tank circuits that were cumbersome to change. Experiments were conducted over a range of plasma current, but detailed studies were performed at 250 kA plasma current where the fractional drive by OFCD was maximum. The amount of OFCD-driven current is most sensitive to the oscillator voltages and frequency, and the added/subtracted current did not depend strongly on the background current over MST’s range $I_p = 250$–500 kA.

The OFCD drive and anti-drive result from MST that is analogous to the ZT-40M experiment is shown in figure 76 [632]. The added current reaches 20 kA by the end of the current flattop. Note that the current rises continuously during the flattop period of the pulse, which is a consequence of the plasma’s inductive back reaction, $L/R \approx 30$ ms. If the pulse duration were longer, the anticipated OFCD-driven current would saturate at $\Delta I_p \approx 40$ kA, which is 15% of the total current.

The maximum OFCD-driven current in MST plasmas occurs for oscillator phasing $\delta = \pi/8$. The fractional increase in plasma current versus phase is shown in figure 77. There is a more structure in this phase dependence than expected for a global helicity balance. The measured amplitude of magnetic fluctuations associated with tearing instabilities is also phase dependent, as shown in figure 78. Interestingly, the cycle-average amplitude for poloidal mode $m = 0$ is somewhat smaller than for conventional flattop induction (dotted lines), with the minimum amplitude occurring for $\delta \approx \pi/8$.

The evolution of the magnetic equilibrium was studied [633] using a combination of toroidal equilibrium reconstructions [27] and relaxed-state modelling. Time-resolved evalu-
Figure 76. Time evolution of plasma current in MST discharges with OFCD (drive and antidrive) compared to the case without OFCD. Reprinted figure with permission from [632], Copyright 2006 by the American Physical Society.

Figure 77. Phase dependence of the plasma current change (cycle averaged) in MST OFCD experiments (squares) compared to a DEBS numerical simulation (green circles) and to the case without OFCD (dashed line with dotted uncertainties). Reprinted from [633], with the permission of AIP Publishing.

Figure 78. During an MST OFCD pulse: phase dependence of \( m/n = 0/1-4 \) (a) and \( m/n = 1/6-15 \) (b) tearing mode amplitude during the OFCD discharges. Error bars indicate the pulse-to-pulse variability. The OFCD off case is shown for comparison (dashed lines, with dotted lines pulse-to-pulse variability). Reprinted figure with permission from [632], Copyright 2006 by the American Physical Society.

Figure 79. In MST, phase dependence of (a) energy confinement time, (b) plasma thermal pressure, (c) thermal energy, (d) magnetic energy, (e) thermal input power, (f) power transfer from magnetic field, (g) total input power for discharge ensembles with OFCD (solid) and without (dashed, with dotted uncertainties). All quantities are cycle averages. Reprinted from [633], with the permission of AIP Publishing.

experimentations of the magnetic energy and helicity balance were also performed. Key quantities for the energy balance with varied phasing are summarized in figure 79. An important observation in the magnetic helicity balance (not shown) was a measured increase in helicity dissipation for \( \delta = \pi/2 \), negating the helicity injection added by OFCD.

Nonlinear, 3D, resistive MHD simulations with fractional OFCD current drive were performed to compare with the MST experimental results [633]. The green data points in figure 77 show that the fractional current drive in the simulations has a similar dependence on the relative phase. The lack of maximum current drive at \( \delta = \pi/2 \) is attributed to an increase in the dissipation of magnetic helicity. The experiments show an important additional effect in that the energy confinement is phase dependent, and not surprisingly maximum for \( \delta \approx \pi/8 \) (figure 79). The simulations used a fixed resistivity profile without evolution of the thermal energy, so they could not make predictions regarding changes in confinement. Nevertheless, the good agreement between the simulations and experiment provide confidence that the OFCD physics basis is sound.

Since OFCD tends to drive current in the outer region of the plasma (figure 75) while conventional induction peaks in the core, their mixture has a flatter current drive profile. This creates a form of current profile control, which has been studied for optimization in nonlinear MHD modelling [634]. The MST experiments have such a mixture, and the reduction in \( m = 0 \) tearing amplitude for \( \delta = \pi/8 \) phasing (figure 78) together with the cycle-averaged improvement in energy confinement (figure 79) are evidence for this control.

A few OFCD experiments have also been conducted on RFX-mod [635]. The ac loop voltages were small in comparison to those used on MST, and the projected current drive from zero-dimensional energy balance was only 0.4% for
800 kA plasmas, making the detection of net current challenging. Measurements of MHD behaviour suggested that the plasma relaxation process was nevertheless affected.

9.4. Future directions

In summary, the visco-resistive, 3D MHD modelling together with experiments from MST and ZT-40M have substantially improved the physics basis for OFCD. While the basic expectations hold, the dynamics associated with relaxation and confinement must be better understood to fully establish the OFCD concept. It is also necessary to test OFCD in RFP plasmas with larger size and temperature, i.e., at larger Lundquist number. Figure 73(b) illustrates the gap between MST and a typical reactor plasma. The ac ripple in the equilibrium magnetic field, indicated by the modulation in the reversal parameter, $F$, exceeds the range of normal operation for $S \lesssim 5 \times 10^2$. Hence, 100% current drive by OFCD in MST (or RFX-mod2) would have excessive equilibrium excursions. The ac ripple is not a subtle physics effect rather the expected response for an ac electric field applied to a circuit that has inductance. The scaling for the ac ripple allows a projection of the size and current for a next-step RFP experiment needed to demonstrate 100% current sustainment by OFCD.

The tendency for RFP plasmas to enter the quasi-single-helicity (QSH) regime at high Lundquist number (chapters 2 and 3) could have a substantial impact on the viability and efficacy of OFCD. This is especially true for a development path that relies on the QSH regime to achieve sufficient energy confinement. The detailed relaxation dynamics with OFCD are different than for conventional induction, and this could change access to and stationarity of QSH. Computational studies with OFCD boundary conditions at very high Lundquist number could begin to address such questions, but such simulations are already state-of-the-art with modest Lundquist and magnetic Prandtl numbers. Fractional OFCD current drive experiments in RFX-mod2 at 2 MA could probe the compatibility and impact of and on QSH. Such experiments might require significant upgrades to the power supplies.

10. RFP reactor studies

Several fusion power plant studies based on the RFP configuration were undertaken in the 1980s with a goal to identify compact, high-power-density options as the basis for fusion reactors with mass power density competitive with fission reactors [636, 637]. The TITAN reactor concept [638–641] was the most detailed of these studies for the RFP. More recently, based on the results of recent experiments and on the expected advancements by RFX-mod2, new studies on the RFP as a neutron source for a fusion–fission hybrid reactor have been undertaken [376, 642].

10.1. The TITAN system study

The high $\beta$, high density RFP plasma makes it well suited for high power density. Small size is also enabled by its relatively small field strength at the magnets, i.e., the engineering beta is very high. The possibility for ohmic ignition [612] eliminates major complications related to auxiliary heating and current drive. The TITAN study introduced several novel features in fusion reactor design, including steady-state using inductive current drive via OFCD, single-piece maintenance for high availability, high core radiation to distribute the heat load, and an integrated-blanket-coil concept that allowed the lithium blanket circuit to serve as the toroidal field magnet, since the toroidal field for the RFP is small. The primary magnets were permitted to be normal conducting, which reduced shielding requirements. Superconducting windings were only used for the vertical field equilibrium coils.

The TITAN design is for a 1 GW e class power plant, and the fusion power core is very compact. The key power plant design parameters are given in table 3. The system analysis showed a strong correlation between the neutron wall load ($I_w$) and the mass power density (MPD), i.e., the ratio of the net electric power to the mass of the fusion power core, about 800 kWe/tonne for the chosen design. The large $I_w = 18$ MW m$^{-2}$ operating point was selected, in part, to bolster the concept of single-piece maintenance for a high plant availability of 85%. Single-piece maintenance was considered possible for a neutron wall load $I_w > 10$ MW m$^{-2}$ [639].

The study’s analysis showed that the projected COE decreases with both MPD and neutron wall load with a formal minimum value at $I_w = 18$ MW/m$^2$, which was taken to be the TITAN operating point. However, the COE flattens for neutron wall load starting around 5 MW m$^{-2}$, and the design purposely targeted very high MPD. An update to TITAN’s costing was later performed by Miller [643], replacing the original costing model with those used in the ARIES tokamak studies, culminating in the ARIES-AT advanced tokamak design [644]. The update also incorporated the level of safety assurance passive-safety cost credits introduced by the ESECOM, and the blanket was updated to a modern configuration used in the ARIES-CS compact stellarator study [645]. The COE versus neutron wall load for the updated analysis is shown in figure 80, revealing a new formal minimum at $I_w = 13$ MW m$^{-2}$ and still flat COE above 5 MW m$^{-2}$. The figure details the ranges of operating points for different aspect ratios (all with circular cross section).

A neutron wall load $I_w > 5$ MW m$^{-2}$ is today considered very challenging and perhaps beyond material feasibility. The 1 GW e class ARIES studies based on tokamak and stellarator configurations all settled on values $I_w < 5$ MW m$^{-2}$. While the TITAN study shows a compact fusion power core based on the RFP is possible, it should not be viewed as defining the RFP target. A more general goal of the study was to assess the merits of high power density, and few magnetic configurations allow the possibility for extreme compactness. In a contemporary context, the TITAN study provides support for a compact RFP power core with lower net electric output that might be better suited for today’s energy market, and it also provides support for a 1 GW e reactor having $I_w < 5$ MW m$^{-2}$ without a large cost penalty by maintaining the simplicity of ohmic ignition, inductive sustainment, and minimal field strength at the magnets, while eliminating auxiliary heating and current drive.
Table 3. Operating parameters of the TITAN fusion power core. There are two design versions with differing structural materials, TITAN-I with vanadium alloy and TITAN-II with ferritic steel, each with separate cooling methods. Adapted from [639], Copyright 1993, with permission from Elsevier.

<table>
<thead>
<tr>
<th>Parameter</th>
<th>TITAN-I</th>
<th>TITAN-II</th>
</tr>
</thead>
<tbody>
<tr>
<td>Major radius (m)</td>
<td>3.9</td>
<td></td>
</tr>
<tr>
<td>Minor radius (m)</td>
<td>0.6</td>
<td></td>
</tr>
<tr>
<td>Plasma current (MA)</td>
<td>17.8</td>
<td></td>
</tr>
<tr>
<td>Toroidal field at plasma surface (T)</td>
<td>0.36</td>
<td></td>
</tr>
<tr>
<td>Poloidal beta</td>
<td>0.23</td>
<td></td>
</tr>
<tr>
<td>Plasma ohmic heating at operating point (MW)</td>
<td>30</td>
<td></td>
</tr>
<tr>
<td>Neutron wall load (MW m$^{-2}$)</td>
<td>18</td>
<td></td>
</tr>
<tr>
<td>Radiation heat flux on first wall (MW m$^{-2}$)</td>
<td>4.6</td>
<td></td>
</tr>
<tr>
<td>Fusion power (MW)</td>
<td>2301 (I), 2290 (II)</td>
<td></td>
</tr>
<tr>
<td>Total thermal power (MW)</td>
<td>2935 (I), 3027 (II)</td>
<td></td>
</tr>
<tr>
<td>Net electric power (MW)</td>
<td>970 (I), 900 (II)</td>
<td></td>
</tr>
<tr>
<td>MPD (kWe/tonne)</td>
<td>757 (I), 806 (II)</td>
<td></td>
</tr>
</tbody>
</table>

![Figure 80. 1 GWe RFP power-plant COE as a function of FW neutron wall load, $I_w$, for several plasma aspect ratios, $A = R/a$. The plant availability is assumed to be $p_f = 0.85$. [643] 2009, reprinted by permission of the publisher (Taylor & Francis Ltd, http://www.tandfonline.com.)](image)

10.2. Studies on the RFP as a neutron source

The RFP as a neutron source for the development of fusion materials and technology was also discussed in the 1980s. In just the past few years, a renewed interest in the RFP neutron source for a fusion–fission hybrid reactor has emerged [376, 642]. Most of the features that make TITAN attractive carry over immediately for the RFP as a compact neutron source. One difference in operation is the use of conventional pulsed induction with short dwell time between pulses as a way to achieve up to 90% duty cycle, motivated by the fact that the existing RFP database is based almost entirely on pulsed operation and the development path might be faster, considering also that the OFCD is very challenging. Preliminary studies suggest a compact device with $R = 4$ m and $a = 0.8$ m, which could produce 30 MW of fusion power with fusion gain $Q = 0.5$ and a neutron wall load below 1 MW m$^{-2}$ [376, 642]. A suggested three-stage approach for a pilot experiment, mainly based on the extrapolation of the RFX-mod performance to be confirmed by the upgraded device RFX-mod2 [114], could fill the gap between today’s experiments and a fusion reactor operating as a neutron source. The first phase would resolve key physics and technology issues by extending operation for plasma current up to $I_p = 9$ MA with 1.5 s pulse duration. In phase 2, superconducting coils for magnetizing and field-shaping windings would be adopted, and operation up to $I_p = 14$ MA with 8 s pulses would be developed. Phase 3 would be a fusion–fission nuclear reactor operating with D-T. The development of a neutron source would validate most of the physics basis and technology required for a pure fusion reactor.

11. Concluding remarks

Over the last three decades research on the RFP magnetic configuration has produced many very important contributions to fusion and plasma science. The potential of the RFP as a fusion power core has been increased, and several fundamental challenges for the RFP are either resolved or better understood. Research on the RFP has also contributed to the science and technology for magnetic fusion more broadly, as well as advanced the physics associated with naturally occurring plasmas, especially in the context of magnetic self-organization. These accomplishments accrued through the successful operation of well-diagnosed experiments coupled closely with comprehensive theory and modelling.

A key transformational result is the experimental demonstration that magnetic fluctuations and stochasticity that challenge confinement are not intrinsic to the RFP configuration. This happens either through inductive control of the plasma current (PPCD) or through magnetic self-organization towards the SH regime which allows RFP sustainment in a state with reduced magnetic fluctuations. Each approach leads to substantially improved confinement. The physics basis for PPCD and SH is now well established, supported by a broad set of nonlinear MHD simulations and theoretical work.
A second transformational achievement for the RFP is the demonstration of reliable control of magnetic instabilities using saddle coils in a feedback process. Active control of MHD stability was pioneered in RFP experiments, and the techniques developed are used within the entire magnetic fusion community, especially in the development of the advanced tokamak. The stabilization of multiple, simultaneously occurring RWMS for pulse duration exceeding many resistive wall times has been routinely achieved. The active control mitigates effects related to magnetic field errors and tearing modes as well, a crucial ingredient for the achievement of long-lived SH states.

RFP research has also been a leading force in the development of modelling of fusion plasmas using a variety of numerical codes and methodologies. This has profoundly influenced our comprehension of RFP physics. They include, for example, nonlinear MHD simulations, tools for the modelling of active control of stability, codes for the threedimensional reconstruction of helical plasma equilibria, codes for the study of turbulence and micro-instabilities, and codes for the description of the dynamics of fast ions and electrons.

The flexible capabilities and set of experimental and modelling tools developed for active control on RFX-mod as an RFP have also prompted the direct exploration of tokamak scenarios in RFX-mod, where \( q(a) < 2 \) plasmas have been produced, stimulating similar experiments in the DIII-D tokamak [120]. A second example concerns the observation that a dynamo conversion of poloidal flux to toroidal flux, analogous to the RFP one, can explain the helical core observed in the hybrid mode of operation of the tokamak [143, 223]. Other examples include multi-configuration modelling of the empirical density limit [466] and the pioneering development of advanced diagnostics for fusion plasmas that include high repetition rate Thomson scattering, combined interferometry-polarimetry (Faraday rotation), full-spectrum motional Stark effect, and a number of advanced x-ray diagnostics. Many of these diagnostics are appearing on new fusion facilities, including ITER.

While the progress is substantial, there remain open questions to resolve the potential of the RFP as an attractive fusion power core that can access ohmic ignition, achieve efficient inductive sustainment, and attain compact high power density. These issues can only be addressed with new experiments combined with further advances in theory and modelling. The upgrade to RFX-mod that is currently underway is a major next step to further advance the RFP at 2 MA, and new experiments with yet larger plasma current and additional capabilities will be needed to fully resolve the viability of the RFP for fusion application. Present-day experiments operate with magnetic field strength that is still at least a factor of five smaller than required in a reactor. Hence the scaling of confinement with current and size remains a key issue, with large room for growth. The improved confinement via inductive control and SH self-organization show promising trends for plasmas having temperatures in excess of 1 keV. The limit to the reduction in magnetic fluctuations has not been identified, and in the cases of highest confinement, evidence suggested microturbulence plays an increasingly important role. A key question is, what transport process limits confinement for the RFP?

The development of suitable control of the plasma–boundary interface is the least developed of the key issues for the RFP. A divertor and/or alternative approaches to plasma exhaust must be tested. There are some special constraints, and possibly opportunities, associated with the integration with the helical equilibrium. Boundary control is anticipated to be important for optimization and stationarity of the SH regime. This will be a significant element of the RFX-mod2 programme, taking advantage of the improved magnetic front-end and much smoother boundary.

None of the key issues for the RFP appear to have show-stoppers. Some of them originate from uncertainty in the physics that calls for continued fundamental theoretical and experimental work, but others are difficult to advance without new experimental capabilities. The RFP programme is at least one generation behind in development relative to the tokamak, mostly due to the smaller resources invested in RFP research.

The RFP continues to have a valuable role in the scientific and technological path towards fusion electricity, to help completing the fusion grand experiment and to assess its own fusion potential, in particular based on its unique opportunity to exploit the low magnetic field at the magnets, the absence of known limitations on aspect ratio, and the ohmic ignition. These features could be extremely valuable in achieving an economical fusion power plant. The future will tell us whether the community will be able to find efficient and smart ways to exploit the opportunity given by the RFP. Hopefully, keeping in mind the concept expressed by Enrico Fermi in 1947: ‘...the vocation of the scientist is to move forward the frontiers of our knowledge in all directions, not just along those that promise more immediate rewards or applause’.

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