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JOREK simulations of runaway electron avalanche and beam formation in DTT

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Declaration

I hereby declare that, the contents and organization of this dissertation constitute my own original work and does not compromise in any way the rights of third parties, including those relating to the security of personal data.

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2026

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Abstract

The Divertor Tokamak Test (DTT) facility plays a strategic role in the European fusion roadmap by bridging the gap between the ITER and EU-DEMO tokamaks regarding power exhaust. Achieving the necessary exhaust conditions requires operating the device at challenging parameters, such as a plasma current of $I_p = 5.5$ MA. This high-current regime introduces the risk of severe disruptions and the generation of relativistic runaway electron (RE) beams via the avalanche mechanism. If uncontrolled, these beams can inflict critical, localized damage on plasma-facing components (PFCs). This thesis provides the first predictive assessment of RE avalanche dynamics in DTT, establishing operational boundaries for the commissioning phase (Day-0, $I_p = 2$ MA) and the Full Power scenario ($I_p = 5.5$ MA).

The analysis relies on the non-linear MHD code JOREK coupled with the STARWALL resistive wall code. A fluid RE model is employed to simulate the self-consistent interaction between REs and the background plasma. Using an artificial thermal quench method, the study systematically scans the parameter space of initial RE seed currents (I_{seed}) and impurity inventories (N_{imp}) to isolate the physics of the current quench (CQ).

The results reveal a fundamental divergence in behavior between the two DTT scenarios, resulting from the exponential sensitivity of the avalanche gain to the pre-disruption plasma current. The Day-0 scenario exhibits intrinsic robustness: due to the limited poloidal flux reservoir, the avalanche gain is moderate. Simulations reveal that even under pessimistic assumptions (i.e., massive neon injection and a high-amplitude initial seed of 20 kA), RE currents peak below 1 MA. These findings confirm that in the Day-0 scenario, REs pose minimal risks to the PFCs, thereby removing the immediate need for complex RE mitigation strategies.

Conversely, the Full Power scenario is characterized by a vanishingly small safety margin against RE formation. The simulations identify a critical impurity threshold at

$N_{\text{imp}} \approx 1 \cdot 10^{21}$. Below this value, the injection yields a partially mitigated disruption, leaving the device vulnerable to severe thermal and halo current loads. Above this limit, however, the high avalanche gain ($G_{\text{av}} \sim 10^5$) amplifies even minimal RE seeds (as low as 5.5 A) into macroscopic RE beams. In this regime, the RE beam effectively replaces the thermal current in the confined plasma region, stalling the CQ. The beam termination is subsequently governed by the vertical stability of the plasma column, since the RE current is determined by a competition between avalanche generation and the physical scraping-off of the plasma volume against the first wall during the vertical displacement event.

These findings indicate that the RE safety margins of the commissioning phase do not scale to full-current operations. Consequently, the disruption mitigation system for the Full Power scenario cannot rely on pure, massive high-Z injection, as this also enhances RE formation. Instead, strategies must prioritize optimized mixtures (e.g., deuterium-rich) to balance the suppression of RE generation with the mitigation of other disruption loads.

Contents

List of Figures	x
------------------------	----------

List of Tables	xxiv
-----------------------	-------------

1 Introduction	1
1.1 Nuclear fusion: core principles	1
1.1.1 Fusion reaction and ignition criteria	1
1.1.2 Magnetic confinement and the tokamak concept	3
1.2 The European roadmap to fusion energy and the DTT facility	6
1.2.1 From ITER to EU-DEMO	6
1.2.2 The DTT facility	7
1.3 Thesis motivation and outline	12
2 Physical background	14
2.1 Magnetohydrodynamics (MHD)	14
2.1.1 Resistive MHD equations	15
2.1.2 MHD stability and the energy principle	21
2.2 Disruptions	24
2.2.1 Disruption causes and dynamics	25
2.2.2 Vertical displacement events and halo currents	27
2.2.3 Disruption mitigation systems	31

2.2.4	The DTT disruption mitigation system	34
2.3	Runaway electrons	34
2.3.1	Runaway dynamics	36
2.3.2	Primary generation	39
2.3.3	Secondary generation (RE avalanche)	42
2.3.4	Transport and losses	46
3	Numerical tools and methodology	49
3.1	The JOEREK code	50
3.1.1	Spatial discretization and numerical implementation	50
3.1.2	The reduced MHD model	52
3.1.3	Equilibrium construction	57
3.2	The RE fluid model	57
3.2.1	Governing equations for the fluid approximation	58
3.2.2	RE sources in the fluid model	59
3.2.3	Self-consistent coupling to MHD	62
3.3	Free boundary simulations	63
3.3.1	Coupling scheme with the STARWALL code	64
3.3.2	DTT structures in STARWALL	66
3.4	Simulation setup and methodology	68
3.4.1	General modeling assumptions	68
3.4.2	Parameter scanning strategy	70
3.4.3	The artificial thermal quench (ATQ)	71
4	Runaway electron generation in the Day-0 scenario	74
4.1	Scenario specifications	75
4.1.1	Equilibrium	75

4.1.2	Physical parameters	77
4.1.3	RE seed and impurity scan	78
4.2	Results	79
4.2.1	Intermediate RE seed (200 A)	79
4.2.2	High-amplitude RE seed (20 kA)	82
4.2.3	Summary of parameter scan	84
5	Runaway electron generation in the Full Power scenario	91
5.1	Scenario specifications	92
5.1.1	Equilibrium	92
5.1.2	Physical parameters	94
5.1.3	RE seed and impurity scan	94
5.2	Results	96
5.2.1	Baseline RE seed (5.5 A)	96
5.2.2	Intermediate RE seed (550 A)	99
5.2.3	High-amplitude RE seed (55 kA)	102
5.3	Discussion and general assessment of RE risks in DTT	106
5.3.1	Comparison: Day-0 vs. Full Power	107
5.3.2	The role of avalanche multiplication	108
5.3.3	Vertical stability and threat assessment	109
5.3.4	Impact of impurities	110
5.3.5	Strategies for Full Power operation	111
6	Conclusions and outlook	112
6.1	Summary of main achievements	112
6.2	Implications for DTT operations	113
6.3	Future outlook	114

Contents	ix
Appendix A Finite element method (FEM)	117
A.1 Weak formulation	117
A.2 Discretization and the Galerkin method	118
Appendix B JOREK normalization	120
Nomenclature	122
References	124

List of Figures

- 1.1 Schematic representation of a tokamak. The confining helical magnetic field line (green) is generated by the superposition of the toroidal field B_ϕ (yellow), produced by the toroidal magnetic field coils, and the poloidal field B_θ (blue), induced by the plasma current I_p (red). Adapted from: Max Planck Institute for Plasma Physics, C. Brandt. 4
- 1.2 Sketch of the diverted tokamak configuration with a single X-point (single null, SN). The magnetic topology is divided by the separatrix (or last closed flux surface, LCFS), which separates the closed flux surfaces of the main plasma (core, in red) from the open field lines of the scrape-off layer (SOL, in orange), which channel particles and heat fluxes towards the divertor plates (light blue). Adapted from [2]. 5

- 1.3 Historical progress of the fusion triple product ($nT\tau_E$) towards ignition conditions for magnetic (MCF) and inertial (ICF) confinement concepts. The red horizontal lines labeled $Q_{\text{sci}}^{\text{MCF}}$ indicate the minimum triple product required to achieve specific gain values for MCF (where $Q_{\text{sci}}^{\text{MCF}} = \infty$ corresponds to ignition). These thresholds are calculated at the central ion temperature T_{i0} that minimizes the required triple product (≈ 20 to 27 keV) and their finite widths represent a range of assumed impurity levels. The red circles illustrate the steady performance improvement over several decades of the tokamak, which is the most advanced MCF concept. The performance range for ITER (red hatched rectangle in the upper-right) is bounded by its projected peak triple product and its original stated mission goal ($Q_{\text{sci}} \geq 10$). ITER represents the crucial step required to bridge the gap between current experimental capabilities and the reactor-relevant regime in the European roadmap. Reproduced from [6]. 8
- 1.4 (a) Sectional view of the DTT facility. (b) Overview of the DTT magnet system, comprising 18 Nb₃Sn toroidal field coils, a 6-module Nb₃Sn central solenoid, 6 external poloidal field coils (two Nb₃Sn and four NbTi), and 6 in-vessel copper coils. Adapted from [12]. . . 9
- 1.5 Top-down view of the DTT heating system layout. The sectors are equipped with electron cyclotron resonance heating (ECRH) launchers, ion cyclotron resonance heating (ICRH) antennas, and a neutral beam injector (NBI). Ports reserved for the remote handling of the first wall (FW-RH) are also indicated. Adapted from [10]. . . 10
- 1.6 Poloidal cross-sections of alternative magnetic configurations feasible in DTT thanks to the flexibility of the PF coil system: single null (SN), double null, snowflake, X-divertor, and double super-X. Adapted from [12]. 11

- 2.1 Schematic evolution of plasma parameters during a disruption involving runaway electron generation. (a) The electron temperature T_e (blue line) collapses during the thermal quench phase (TQ, green shaded area), causing a rapid increase in resistivity and the induction of a large parallel electric field $E_{\parallel}$ (red line) during the current quench phase (CQ, red shaded area). (b) The total plasma current I_p (solid line) shows the characteristic initial spike in the TQ followed by a decay. If a macroscopic RE beam is present, the decaying thermal current I_{th} (dotted line) is replaced by the relativistic RE current I_{RE} (dashed line), potentially leading to a runaway plateau (yellow shaded area). Adapted from [16]. 26
- 2.2 Schematic representation of the electromechanical loads during an upward vertical displacement event (VDE). The halo current (I_h , green arrows) flows along the open magnetic field lines in the scrape-off layer, intercepts the first wall, and closes its circuit poloidally through the conductive vacuum vessel. The interaction between the poloidal component of this current and the toroidal magnetic field (B_{ϕ} , light blue arrow) generates a net vertical Lorentz force (F_z , gray arrow) that acts on the machine structure. Reproduced from [28]. © ITER Organization. 30
- 2.3 Visual evidence of localized damage resulting from the impact of a collimated RE beam on the beryllium limiter of the JET tokamak. The concentrated energy deposition led to significant surface melting of the tiles. Reproduced from [61]. 35

- 2.4 Graphical representation of the force balance governing electron runaway. The red curve depicts the collisional friction force, which for suprathermal electrons decreases with increasing momentum. The dashed horizontal lines represent the driving force from the electric field (eE) and the critical electric field force (eE_c). A critical momentum mv_c exists at the intersection of the electric field force eE and the collisional drag force: for higher momenta (shaded red region), electrons experience a net positive acceleration and enter the runaway regime. The minimum field required to enable this process is the critical electric field E_c ; for fields below this threshold, no runaway generation occurs. Conversely, if E is high enough to exceed the slide-away value of $\sim 0.21eE_D$ (E_D being the Dreicer electric field), the whole electron population is converted into runaways. Reproduced from [65]. 37
- 2.5 Schematic representation of the electron distribution function in the early phases of a mitigated disruption. The injected impurities provide the plasma with cold electrons through ionization (gray and black curves) which slow down the hot electrons (red curve) through collisions. However, the higher-energy tail cools much more slowly and, if the induced electric field is strong enough to define a critical momentum (vertical dashed green line) that falls within this surviving hot tail, these particles are accelerated into the runaway regime. Reproduced from [16]. 41
- 2.6 Poincaré plots from a JOREK simulation showing the evolution of the RE density (n_r) and magnetic field stochasticity during a disruption. Starting from intact surfaces (a), the topology becomes highly stochastic as the disruption progresses (b-d), leading to the loss of REs. As the CQ proceeds, flux surfaces begin to re-form (e-f), effectively trapping the remaining REs in the core. Reproduced from [57]. 47
- 3.1 (a) JOREK cylindrical coordinate system (R, Z, ϕ) used to describe the tokamak geometry. (b) Single 2D Bezier element in the (R, Z) space. The element-local coordinates are also shown: $(s, t) \in [0, 1]$. Adapted from [96]. 51

- 3.2 Poloidal cross-section (R-Z plane) of the active and passive DTT conducting structures as implemented in the STARWALL code. The vacuum vessel is split into an inner shell (axisymmetric thin wall, orange line) and an outer shell (toroidal filaments, red). Active components, including the central solenoid, poloidal field coils, and in-vessel coils, are shown in black. The JOREK computational boundary (green) is defined to closely track the physical first wall (blue) while smoothing sharp corners to ensure numerical stability. The separatrix for the standard single null Day-0 scenario is included in pink for reference. 67
- 4.1 Radial profiles of electron density (n_e , top) and temperature (T , bottom) as a function of the normalized poloidal flux ψ_n . The orange curves represent the equilibrium configuration implemented in the JOREK simulations for the Day-0 scenario. The density profile used in the simulation (orange) is a fit of the original data (blue). Regarding the temperature, the adopted equilibrium employs a simplified bell-shaped profile (orange) rather than the more complex profile of the older dataset (blue). This choice allows for the utilization of the most up-to-date DTT geometry while ensuring the initial magnetic equilibrium remains MHD stable ($q > 1$). The use of a simplified initial temperature profile is physically justified by the specific focus of this study on the CQ phase, during which the temperature evolution is determined self-consistently by the balance between Ohmic heating and impurity radiation, rendering the specific details of the pre-disruption temperature profile largely irrelevant to the final RE avalanche dynamics. 76

- 4.2 Temporal evolution of key plasma parameters for the Day-0 scenario with an initial runaway electron (RE) seed of 200 A. From top to bottom: total plasma current (I_p , solid) and RE current (I_{RE} , dashed); maximum electron temperature ($T_{e,max}$); and total particle content for ions (N_{ion} , solid) and impurities (N_{imp} , dashed). Lighter shades of the same color indicate reference simulations performed without REs; these curves are largely obscured due to the negligible impact of the RE population at this RE seed level. The phases of the artificial thermal quench (ATQ) are delineated by vertical brown lines: the dotted line marks the end of the initial thermal collapse, while the dash-dotted line indicates the start of impurity injection. The sharp transient drop and subsequent recovery of $T_{e,max}$ around $t \approx 1$ ms for the highest impurity scenario (red curve) is due to this artificial setup, as Ohmic heating remains disabled until the impurity injection phase concludes. 80
- 4.3 Poloidal cross sections of the electron temperature (T_e) distribution at the beginning of the simulation ($t = 0$ ms, left) and at the end of the first step of the artificial thermal quench ($t = 0.326$ ms, right), a timestamp consistently marked by the brown dotted vertical lines in the time-evolution figures throughout this chapter. Despite the visual similarity of the relative profiles, note the difference in the color bar scales, highlighting a collapse in peak temperature from 4.28 keV down to approximately 10 eV. 81
- 4.4 Poloidal cross sections and corresponding 1D linear cuts of the current density (j) distribution. The left panels show the state at the end of the first artificial thermal quench step ($t = 0.326$ ms, marked by the brown dotted vertical lines in the time-evolution figures throughout this chapter), while the right panels show the state at the end of the second step ($t = 0.871$ ms, marked by the corresponding dash-dotted lines), immediately before impurity injection. The linear cuts clearly illustrate the flattening of the current profile inside the core during this phase. 81

- 4.5 Time traces of plasma current (I_p , solid) and runaway electron (RE) current (I_{RE} , dashed) for the high-amplitude initial seed case (20 kA). Varying impurity levels are shown by different colors. Lighter shaded lines represent reference cases without REs and the phases of the artificial thermal quench (ATQ) are delineated by vertical brown lines: the dotted line marks the end of the initial thermal collapse, while the dash-dotted line indicates the start of impurity injection. The end of the ATQ and the start of the self-consistent CQ phase is discernible by the distinct change in the I_p slope immediately following the dash-dotted line. Unlike the lower seed cases, a divergence between RE-inclusive and RE-free simulations is observed at the highest impurity level (red lines), where the growth of an RE beam notably alters the current decay dynamics. 83
- 4.6 Poloidal cross-section showing the runaway electron (RE) density distribution at the time of peak RE current ($t \approx 5.7$ ms) for the 20 kA RE seed and maximum impurity case. The white contour marks the separatrix, demonstrating that the REs are contained within the confined volume, which is being reduced by the vertical motion of the plasma. 84

- 4.7 Time evolution of the runaway electron current (top) and the vertical position of the magnetic axis (bottom) for the three initial seed regimes: baseline (2 A, dotted), intermediate (200 A, dashed), and high-amplitude (20 kA, solid). To improve readability and highlight scaling trends, the I_{RE} values for the 2 A seed (dotted lines) have been multiplied by a factor of 100. Color coding indicates the impurity injection levels, while the phases of the artificial thermal quench are delineated by vertical brown lines: the dotted line marks the end of the initial thermal collapse and the dash-dotted line indicates the start of impurity injection. The 'x' markers denote the timestamps when the plasma first contacts the plasma-facing components. The constant slopes clearly observed for high impurity levels (red and light blue) reflect exponential avalanche growth. The superposition of the scaled 2 A and 200 A traces points to a linear dependence on the initial seed magnitude, whereas the 20 kA case reveals non-linear effects. Note that the simulation for the 2 A seed at the highest impurity level (which would appear as a red dotted line) was omitted, as its negligible macroscopic impact is already bounded by the corresponding 200 A case. 86
- 4.8 Evolution of the runaway electron density profile (n_{RE} , expressed in 10^{16} m^{-3}) for the high-amplitude seed scenario ($I_{seed} = 20 \text{ kA}$) with an intermediate impurity content of $N_{imp} \approx 1 \cdot 10^{21}$ (corresponding to the light blue solid traces in Figure 4.7). The white contour denotes the separatrix, distinguishing the confined plasma from the open field line region. The sequence illustrates the volumetric collapse mechanism: as the vertical instability drives the plasma rapidly downward towards the X-point, the confined volume is progressively reduced. Consequently, although the local RE density in the core continues to rise due to the avalanche mechanism, the total RE current (I_{RE} , indicated at the top right of each panel) eventually decays after reaching its peak at $t \approx 6.3 \text{ ms}$, as the reduction in confined volume outpaces the avalanche growth prior to direct contact with the PFCs. 87

- 4.9 Current decomposition for the high-amplitude seed scenario ($I_{\text{seed}} = 20$ kA) with an impurity inventory of $N_{\text{imp}} \approx 3 \cdot 10^{21}$ (corresponding to the solid red traces in Figure 4.7). The plot displays the total plasma current (I_p , solid line), the runaway electron (RE) current (I_{RE} , dashed line), and the total current flowing within the closed flux surfaces (I_{core} , gold dotted line). The phases of the artificial thermal quench are delineated by thin vertical brown lines: the dotted line marks the end of the initial thermal collapse, the dash-dotted line indicates the onset of impurity injection, and the dashed line signals the completion of the ATQ procedure and the start of the self-consistent current quench. Notably, during the artificial current flattening phase (between the dotted and dash-dotted vertical lines), a transient I_p spike occurs alongside the development of a negative edge skin current. An additional red dashed vertical line marks the timestamp of first contact with the plasma-facing components (PFCs). The exponential growth of the RE avalanche deviates from linearity and approaches saturation as I_{RE} nears the available confined current limit (I_{core}), an effect manifesting prior to PFC contact. The subsequent convergence of the I_{core} and I_{RE} curves demonstrates that the current within the confined plasma volume is eventually carried almost entirely by the RE beam, implying that the remaining current ($I_p - I_{\text{RE}}$) consists predominantly of halo currents flowing in the scrape-off layer. 88
- 5.1 Comparison of the equilibrium profiles for the Full Power scenario as a function of the normalized poloidal flux ψ_n . From top to bottom: electron density n_e , electron temperature T_e , FF' , and safety factor q . The blue curves represent the original CREATE-NL data, while the orange curves show the fitted approximations. While the electron density profile was left unmodified from the fit, the profiles for T_e , FF' and q were modified to eliminate pre-existing core instabilities. These modified, MHD-stable profiles ($q > 1$), shown in green, are adopted as the input for the Full Power simulations. 93

- 5.2 Temporal evolution of the total plasma current (I_p , solid lines) and runaway electron current (I_{RE} , dashed lines) for the baseline initial seed case of 5.5 A. The colors distinguish the specific neon impurity inventory (N_{imp}), spanning a range from $1 \cdot 10^{20}$ to $6 \cdot 10^{21}$ particles. The phases of the artificial thermal quench (ATQ) are delineated by vertical brown lines: the dotted line marks the end of the initial thermal collapse, while the dash-dotted line indicates the start of impurity injection. The termination of the final ATQ phase (impurity injection) is discernible by the distinct change in the I_p slope immediately following the dash-dotted line; this transition marks the onset of the self-consistent current quench, which is progressively delayed at higher impurity levels because the fixed injection rate requires a longer duration to reach the target inventory. The 'x' markers denote the timestamps of first plasma contact with the plasma-facing components. A key finding is the formation of macroscopic RE beams (≈ 0.7 MA) in the high-impurity regimes ($N_{imp} \gtrsim 3 \cdot 10^{21}$, green, orange, and red traces). This massive amplification, occurring despite both PFC contact and the minimal baseline initial seed, serves as a clear demonstration of the high avalanche gain characteristic of the Full Power scenario. 97

- 5.3 Simultaneous evolution of the total plasma current (I_p , top panel) and the vertical position of the magnetic axis (Z_{axis} , bottom panel). The simulations initialized with the baseline initial seed ($I_{seed} = 5.5$ A) are represented by thick solid lines, whereas the thin dashed lines denote reference cases in which the avalanche source term was numerically suppressed. As in the preceding figure, color-coding indicates the impurity inventory, and vertical brown lines delineate the artificial thermal quench transitions. The 'x' and diamond markers denote the timestamps of first plasma contact with the plasma-facing components (PFCs) for the simulations with and without runaway electrons (REs), respectively. The bottom panel highlights that higher impurity injections drive more rapid vertical displacement events (yielding steeper slopes), thereby advancing the onset of PFC contact. Notably, a delayed dynamical impact of the RE population is observed: due to the relatively small initial seed, the avalanche requires several milliseconds to amplify the current to macroscopic levels sufficient to alter the current quench dynamics. Consequently, the inductive braking effect on the vertical motion, discernible as the divergence between the solid and dashed curves, manifests exclusively during the late stages of the disruption. 98
- 5.4 Temporal evolution of the total plasma current (I_p , solid lines) and runaway electron (RE) current (I_{RE} , dashed lines) for the intermediate initial seed scenario ($I_{seed} = 550$ A). Color-coding and the vertical brown markers delineating the artificial thermal quench phases follow the convention established in previous figures, while the 'x' markers denote the timestamps of first plasma contact with the plasma-facing components. A comparison with the baseline seed scenario highlights a substantial reduction in the impurity threshold required for beam formation, with macroscopic RE currents emerging at significantly lower impurity inventories. Furthermore, for intermediate impurity levels, the magnitude of the resulting RE beam is notably higher, culminating in a peak RE current of approximately 1.8 MA, constituting nearly 60% of the instantaneous plasma current. 100

- 5.5 Detailed decomposition of the total plasma current components for the intermediate seed case ($I_{\text{seed}} = 550$ A) with an impurity inventory of $N_{\text{imp}} \approx 4 \cdot 10^{21}$ (corresponding to the orange traces in Figure 5.4). The plot displays the total plasma current (I_p , solid line), the runaway electron current (I_{RE} , dashed line), and the total current flowing within the closed flux surfaces (I_{core} , gold dotted line). The vertical brown lines follow the artificial thermal quench conventions established in preceding figures, with an additional brown dashed vertical line at $t \approx 1.6$ ms indicating the completion of the ATQ procedure and the onset of the self-consistent current quench. An orange dashed vertical line at $t \approx 4$ ms marks the timestamp of first contact with the plasma-facing components (PFCs). Notably, during the artificial current flattening phase (between the dotted and dash-dotted vertical lines), a transient I_p spike occurs alongside the development of a negative edge skin current. A key observation is the convergence of the I_{core} and I_{RE} curves after $t \approx 6$ ms; this overlap demonstrates that the current within the confined plasma volume is carried almost entirely by the RE beam despite continuous wall-scraping against the PFCs, implying that the remaining current ($I_p - I_{\text{RE}}$) consists predominantly of halo currents flowing in the scrape-off layer. 101
- 5.6 Temporal evolution of the total plasma current (I_p , solid curves) and runaway electron (RE) current (I_{RE} , dashed curves) for the high-amplitude initial seed case of 55 kA. Color-coding and the vertical brown markers delineating the artificial thermal quench phases follow the conventions established in preceding figures. The 'x' markers denote the timestamps of first plasma contact with the plasma-facing components (PFCs). Under these high-amplitude seed conditions, the avalanche dynamics are exceptionally efficient: the RE fraction rapidly escalates to approximately 80% of the total current within a few milliseconds, even for moderate impurity inventories ($N_{\text{imp}} \approx 1 \cdot 10^{21}$). The resulting beams reach peak values of ≈ 3.2 MA, a magnitude sufficient to induce a prominent prolongation of the current quench by significantly reducing the total plasma current decay rate. 103

- 5.7 Comparison of the total plasma current (I_p , top panel) and the vertical position of the magnetic axis (Z_{axis} , bottom panel). Thick solid lines represent the high-amplitude initial seed scenario ($I_{seed} = 55$ kA), whereas thin dashed lines denote reference simulations in which the avalanche mechanism was numerically suppressed. Color-coding and the vertical brown markers delineating the artificial thermal quench phases follow the conventions established in preceding figures. The 'x' and diamond markers denote the timestamps of first plasma contact with the plasma-facing components (PFCs) for the simulations with and without runaway electrons (REs), respectively. In this high-amplitude seed scenario, the RE beam carries a dominant fraction of the current, thereby significantly mitigating both the vertical displacement velocity and the current quench decay rate compared to the RE-free reference cases. 104
- 5.8 Evolution of the runaway electron density profile (n_{RE} , expressed in 10^{16} m^{-3}) for the high-amplitude initial seed scenario ($I_{seed} = 55$ kA) with an impurity inventory of $N_{imp} \approx 3 \cdot 10^{21}$ (corresponding to the green traces in Figures 5.6 and 5.7). The white contour denotes the separatrix, distinguishing the confined plasma from the open field line region. The sequence illustrates the wall-scraping mechanism: as the vertical instability drives the plasma upward, the confined volume is progressively reduced by interaction with the first wall. The onset of this physical contact with the plasma-facing components (PFCs) corresponds precisely to the green 'x' markers in Figures 5.6 and 5.7. Although the local RE density in the core continues to rise due to the avalanche mechanism, the total RE current (I_{RE} , indicated at the top right of each panel) eventually decays due to this geometric reduction of the confined plasma volume. Notably, the absolute peak of I_{RE} for this scenario occurs at $t \approx 5.8$ ms, falling between the discrete snapshots displayed here. 105

- 5.9 Comparative analysis of runaway electron current evolution for the Day-0 scenario ($I_p = 2$ MA, dashed lines, adapted from Figure 4.5) versus the Full Power scenario ($I_p = 5.5$ MA, solid lines). Both configurations are initialized with an identical initial seed of 20 kA to isolate the effect of the plasma current on the RE multiplication factor. The green and blue traces correspond to impurity inventories of $N_{\text{imp}} \approx 1 \cdot 10^{21}$ and $3 \cdot 10^{21}$ particles, respectively. The 'x' markers denote the timestamps of first plasma contact with the plasma-facing components (PFCs). The use of a semi-logarithmic scale underscores the divergence in avalanche dynamics: the Full Power cases exhibit visibly steeper growth rates and higher peak currents compared to the Day-0 benchmarks, confirming the exponential sensitivity of the avalanche gain to the background plasma current magnitude. . . . 107

List of Tables

1.1	Comparison of the main parameters among DTT, ITER, and EU-DEMO [11].	9
4.1	Physical parameters adopted for the Day-0 simulations. The plasma resistivity (η_e) and parallel thermal diffusivity ($\chi_{\parallel}$) follow standard scalings but are regularized by constant cut-off values. The perpendicular thermal diffusivity ($\chi_{\perp}$) utilizes a radial hyperbolic tangent profile dependent on the normalized poloidal flux ψ_n , maintaining a core value and transitioning sharply to zero at $\psi_n = 0.92$ over a transition width of 0.01. The particle diffusivity (D) is isotropic for the thermal species, including impurities. In contrast, the RE fluid employs a highly anisotropic diffusion tensor. The value of the parallel component ($D_{\parallel, \text{RE}}$) was selected based on sensitivity studies of the RE fluid model in the literature, to optimally mimic fast parallel transport and minimize artificial perpendicular diffusion [112, 113]. The perpendicular component ($D_{\perp, \text{RE}}$) is maintained at a minimal level, in order to ensure numerical stability [113].	77

5.1	Summary of the physical coefficients employed for the Full Power scenario. The resistivity (η_e) and parallel thermal diffusivity ($\chi_{\parallel}$) are based on standard Spitzer and Spitzer-Härm scalings, respectively, saturated at specific temperature thresholds. A radial hyperbolic tangent profile is used for the perpendicular thermal diffusivity ($\chi_{\perp}$), which drops to zero around $\psi_n = 0.92$ over a width of 0.01. The thermal particle diffusivity (D) is constant and isotropic. For the runaway electron fluid, a highly anisotropic diffusion tensor is employed. The large value for parallel diffusivity ($D_{\parallel, \text{RE}}$) was selected based on extensive benchmarks of the RE fluid model, as it provides the optimal balance between capturing fast parallel transport and minimizing numerical errors in the perpendicular direction [112, 113]. The specific physical value corresponds to the exact same normalized input used for the Day-0 scenario, adjusted here by the higher density normalization factor. The perpendicular component ($D_{\perp, \text{RE}}$) is kept at the minimum value required for numerical stability [113].	95
B.1	Connection between SI and normalized units for various plasma parameters, adapted from [96]. Variables with the subscript ‘SI’ denote physical units, while those without are normalized. The normalization factors n_0 and ρ_0 are the central number and central mass densities, respectively, while μ_0 is the vacuum permeability. Additionally, k_B is the Boltzmann constant, γ is the adiabatic index ($\gamma = 5/3$ for plasmas), and $\mu_{\text{imp}} = m_i/m_{\text{imp}}$	121

Chapter 1

Introduction

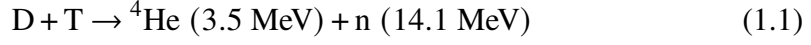
The quest for a sustainable, safe, and unlimited energy source is one of the defining challenges of the present century. As global power needs grow and the effects of climate change become increasingly severe, replacing fossil fuels is non-negotiable. Nuclear fusion offers the most promising path forward for baseload electricity generation, providing a high energy density without producing carbon emissions (during operation) or high-level, long-lived radioactive waste.

This work focuses on one of the main operational challenges in magnetic confinement fusion: the generation and mitigation of runaway electrons (REs) during plasma disruptions. Specifically, it presents a predictive analysis for the Divertor Tokamak Test (DTT) facility, evaluating the risks associated with its initial commissioning phase (Day-0) and its reactor-relevant (Full Power) scenario using the non-linear magnetohydrodynamic (MHD) code JOREK.

1.1 Nuclear fusion: core principles

1.1.1 Fusion reaction and ignition criteria

Nuclear fusion is the process that powers stars, where light atomic nuclei combine to form heavier ones, releasing massive amounts of energy. To replicate this process on Earth, the most viable reaction utilizes two isotopes of hydrogen: deuterium (D) and tritium (T).



To achieve a sufficiently high reaction probability, temperatures higher than 10 keV (approximately 100 million Kelvin) are required. At these extreme thermal energies, electrons are stripped from their atomic nuclei, creating a fully ionized gas: the *plasma*. Often described as the fourth state of matter, plasma consists of a quasi-neutral collective of positive ions and free electrons. This state is characterized by a fundamental length scale known as the *Debye length* (λ_D): on scales smaller than λ_D , individual particle interactions dominate, whereas on spatial scales larger than λ_D , the plasma maintains quasi-neutrality and exhibits collective behavior. It is precisely this collective electromagnetic (EM) nature that distinguishes plasma from a neutral gas, rendering it electrically conductive and highly responsive to EM fields, thereby enabling its magnetic confinement.

D-T reactions release 17.6 MeV per event, the majority of which is carried by high-speed neutrons that are captured to generate electricity. In contrast, the alpha particles (helium nuclei) deposit their energy back into the plasma. For this self-heating to drive *ignition* (a self-sustained reaction), the plasma must satisfy the *Lawson criterion*: the *triple product* of particle density (n), temperature (T), and energy confinement time (τ_E) must exceed a critical threshold [1]:

$$nT\tau_E \geq 5 \cdot 10^{21} \text{ m}^{-3} \text{ s keV} \quad (1.2)$$

Due to the difficulty of attaining ignition, external energy is needed to sustain the reactions. The proximity to ignition is quantified by the *scientific energy gain* Q_{sci} , which is defined as the generated fusion power divided by the external heating power delivered to the plasma. This parameter goes to infinity when ignition is reached, and *scientific breakeven* is defined as $Q_{\text{sci}} = 1$. However, achieving scientific breakeven does not guarantee net electricity generation due to the power required for operation, as well as energy conversion inefficiencies and system losses. When referring to the electric power ratio (electric output power divided by electric power consumed), we refer to the *engineering energy gain* Q_{eng} , which is strictly lower than Q_{sci} . Reaching $Q_{\text{eng}} > 1$ is a necessary condition to obtain an economically profitable nuclear fusion power plant: in that case, the system would function as an electrical energy amplifier.

1.1.2 Magnetic confinement and the tokamak concept

To satisfy the Lawson criterion, the plasma must be effectively confined. Two primary approaches exist to achieve this: inertial confinement fusion (ICF), which relies on the inertia of the fuel mass to resist expansion during rapid compression, and magnetic confinement fusion (MCF), which employs strong magnetic fields to confine the electrically charged plasma.

For continuous energy production, MCF is currently the most developed approach. The fundamental principle relies on the Lorentz force, which compels charged particles to gyrate around magnetic field lines, effectively restricting their motion perpendicular to the field. To prevent the end losses associated with linear configurations, the magnetic field lines are closed into a toroidal (doughnut-shaped) geometry. However, a purely toroidal magnetic field is insufficient for confinement: the curvature and gradient of the magnetic field cause ions and electrons to drift in opposite vertical directions. This charge separation generates a vertical electric field, which, coupled with the magnetic field ($\mathbf{E} \times \mathbf{B}$ drift), would rapidly expel the plasma outward. To counteract these drifts, the magnetic field lines must be twisted helically as they traverse the torus, connecting regions of favorable and unfavorable curvature, averaging out the drift effects, and ensuring equilibrium.

Although numerous magnetic configurations have been explored, the *tokamak* (a Russian acronym for “toroidal machine with magnetic coils”) has emerged as the most scientifically mature concept, representing the primary focus of this work. As illustrated in Figure 1.1, the tokamak is an axisymmetric device based on a toroidal geometry, characterized by a major radius R (distance from the machine axis to the plasma center) and a minor radius a (radius of the plasma cross-section). In this topology, two principal directions are defined: the toroidal direction, which follows the major circumference of the torus (yellow lines in Figure 1.1), and the orthogonal poloidal direction, which follows the minor circumference of the plasma cross-section (blue lines in Figure 1.1). The fundamental helical magnetic structure required for confinement is generated through the superposition of two distinct magnetic field components, defined along these directions:

1. The toroidal field B_ϕ : This dominant component is produced by a set of external coils wound around the torus: the toroidal field (TF) coils. Physically,

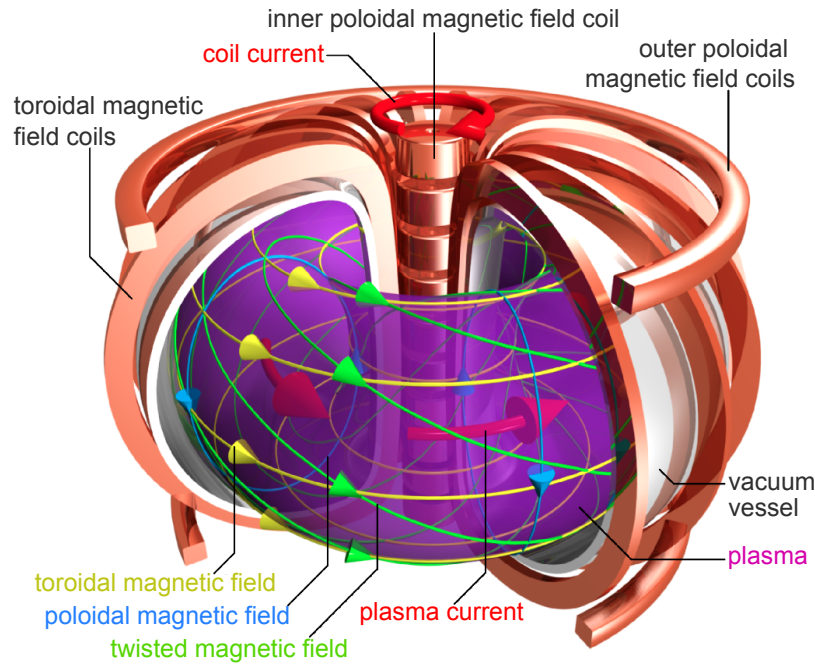


Fig. 1.1 Schematic representation of a tokamak. The confining helical magnetic field line (green) is generated by the superposition of the toroidal field B_ϕ (yellow), produced by the toroidal magnetic field coils, and the poloidal field B_θ (blue), induced by the plasma current I_p (red). Adapted from: Max Planck Institute for Plasma Physics, C. Brandt.

these coils approximate a solenoid bent into a circle, where electric current flows in the poloidal direction to generate a field in the toroidal direction.

2. The poloidal field B_θ : This field is generated primarily by a strong toroidal electric current flowing within the plasma itself (I_p). This current is driven via the transformer principle: the plasma ring acts as the secondary winding, while a central solenoid (CS) serves as the primary. By ramping the current in the central solenoid, a time-varying magnetic flux is induced, generating the electromotive force required to drive I_p . A significant intrinsic limitation of this inductive method is that the magnetic flux variation can only be sustained for a finite time. Consequently, without the application of non-inductive current drive techniques, the tokamak is inherently a pulsed device, unable to sustain a continuous plasma current (and thus confinement) indefinitely.

Beyond these two fields, external poloidal field (PF) coils are employed to generate a vertical magnetic field component, which is essential for maintaining radial equilibrium and for shaping the poloidal cross-section of the plasma.

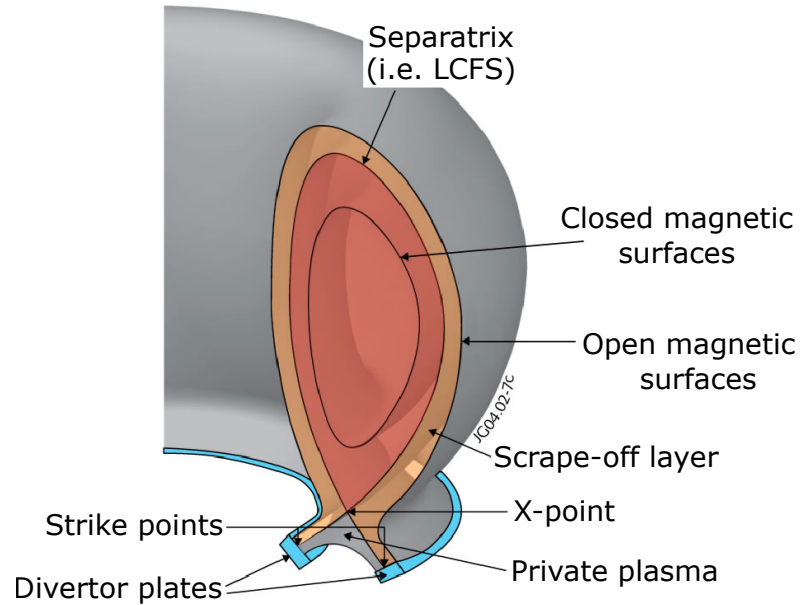


Fig. 1.2 Sketch of the diverted tokamak configuration with a single X-point (single null, SN). The magnetic topology is divided by the separatrix (or last closed flux surface, LCFS), which separates the closed flux surfaces of the main plasma (core, in red) from the open field lines of the scrape-off layer (SOL, in orange), which channel particles and heat fluxes towards the divertor plates (light blue). Adapted from [2].

The resultant magnetic field lines twist helically around the torus, tracing out nested magnetic flux surfaces. These surfaces are defined as iso-surfaces of the poloidal magnetic flux ψ , which represents the total magnetic flux passing through a surface bounded by a toroidal ring around the machine axis. Since ψ remains constant on a given magnetic surface but varies from the magnetic axis to the plasma edge, it serves as a natural radial coordinate to label the flux surfaces, which is particularly useful in shaped plasmas. In modern reactor-relevant configurations, as depicted in Figure 1.2, this magnetic topology is shaped to include a single magnetic null (*X-point*), creating what is known as the *single null* (SN) configuration. This null defines a boundary known as the *separatrix*, or *last closed flux surface* (LCFS), which divides the closed flux surfaces of the confined core plasma from the open field lines of the scrape-off layer (SOL). Particles and heat crossing the separatrix flow along these open field lines into a dedicated component called the *divertor*, where they are neutralized and pumped away.

Another fundamental parameter defining the topology of these surfaces is the safety factor q , which represents the number of toroidal turns m a field line makes for

one poloidal turn n . The radial profile of q (q -profile) plays a decisive role in the MHD stability of the plasma. Rational values of $q = m/n$ characterize resonant surfaces prone to instabilities, a mechanism that will be analyzed in detail in Chapter 2.

While the plasma current I_p is essential for confinement, it also provides initial heating via the Ohmic (Joule) effect. However, this Ohmic heating becomes increasingly inefficient as temperature rises, since the Spitzer resistivity of the plasma scales as $\eta_e \propto T^{-3/2}$ [3]. Therefore, to reach fusion-relevant temperatures, auxiliary heating systems are mandatory to overcome this limitation. These include the injection of high-energy neutral atoms (neutral beam injection, NBI) and the absorption of EM waves at specific resonant frequencies (ion cyclotron resonance heating, ICRH, and electron cyclotron resonance heating, ECRH).

With the fundamental principles of nuclear fusion and tokamaks established, the global fusion effort has shifted towards their integration into a viable energy production strategy. The transition from experimental devices to a commercial reactor requires not only achieving scientific breakeven but also solving critical engineering hurdles. Most notably, challenges arise regarding the management of extreme heat and particle fluxes at the plasma boundary: as the power crossing the separatrix increases, the structural integrity of the divertor targets becomes one of the most limiting factors for machine operation. The following section outlines the European roadmap to the realization of a commercial reactor, highlighting the role of the DTT facility as a crucial bridge to demonstration power plants.

1.2 The European roadmap to fusion energy and the DTT facility

1.2.1 From ITER to EU-DEMO

The transition from experimental physics to a commercially viable energy source is guided by the *European Research Roadmap to the Realization of Fusion Energy* [4], which outlines a stepwise approach to demonstrate the feasibility of fusion as a baseload power source. The historical progress towards this goal is visually summarized in Figure 1.3, which plots the evolution of the triple product ($nT\tau_E$) for both magnetic (MCF) and inertial (ICF) confinement concepts. As illustrated, while

current tokamak experiments (red circles) have approached the scientific breakeven threshold ($Q_{\text{sci}}^{\text{MCF}} = 1$ in the picture), a significant leap in performance is required to enter the $Q_{\text{eng}} > 1$ regime, essential for a fusion power plant [5, 6].

The central pillar of the European roadmap is the ITER project [7], currently under construction in Cadarache, France. ITER is an international effort whose main goals are to demonstrate the scientific feasibility of fusion at near power plant size by producing 500 MW of fusion power from 50 MW of input power ($Q_{\text{sci}} \geq 10$). It will also serve as a testbed for vital technologies such as on-site tritium generation, remote handling, and heating systems.

The final step before commercial deployment is the European demonstration power plant, EU-DEMO [8], which aims to prove the technological and economic feasibility of nuclear fusion by supplying net electricity to the grid and breeding its own tritium. While ITER serves as a critical milestone for EU-DEMO, significant technological gaps persist between the two devices, most notably, the problem of power exhaust. In a reactor-scale device like EU-DEMO, the power flux crossing the separatrix and flowing to the divertor (P_{sep}) will be significantly higher than in current machines. While ITER is designed to operate in an SN configuration with a tungsten divertor capable of withstanding steady-state heat fluxes of 10 MW/m² (and up to 20 MW/m² for slow transients), the conditions in EU-DEMO will be far more severe. Should the ITER solution prove insufficient (or if the specific constraints of a commercial reactor require changes), alternative solutions must be qualified, such as advanced magnetic configurations or liquid metal divertors. However, the ITER design is too rigid to accommodate major changes to the magnetic topology or divertor material composition, and testing these solutions directly in EU-DEMO might pose unacceptable operational risks. This creates a strategic need for a facility that matches the divertor conditions of ITER and EU-DEMO while retaining the flexibility to explore diverse power exhaust strategies. The Divertor Tokamak Test (DTT) facility [9, 10], currently under construction at the ENEA Frascati Research Center in Italy, is designed specifically to bridge this gap.

1.2.2 The DTT facility

The mission of DTT is to replicate the divertor conditions of ITER and EU-DEMO in a device with manageable dimensions ($R = 2.19$ m, $a = 0.70$ m), allowing

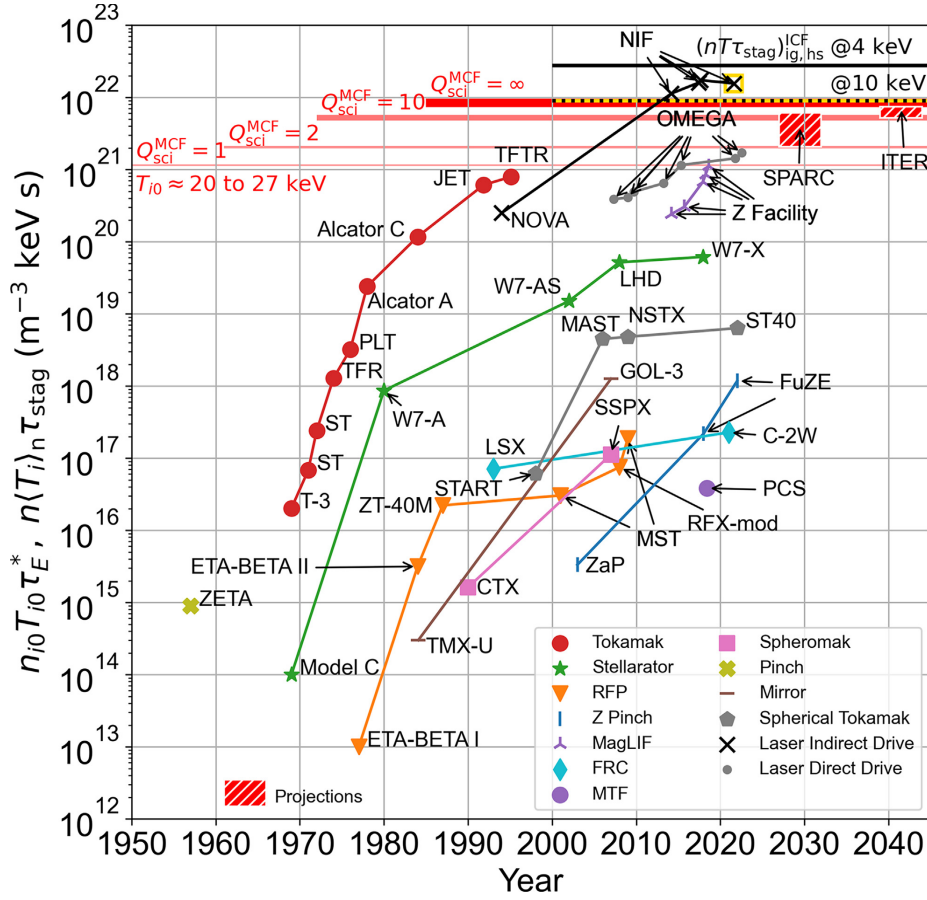


Fig. 1.3 Historical progress of the fusion triple product ($nT\tau_E$) towards ignition conditions for magnetic (MCF) and inertial (ICF) confinement concepts. The red horizontal lines labeled $Q_{\text{sci}}^{\text{MCF}}$ indicate the minimum triple product required to achieve specific gain values for MCF (where $Q_{\text{sci}}^{\text{MCF}} = \infty$ corresponds to ignition). These thresholds are calculated at the central ion temperature T_{i0} that minimizes the required triple product (≈ 20 to 27 keV) and their finite widths represent a range of assumed impurity levels. The red circles illustrate the steady performance improvement over several decades of the tokamak, which is the most advanced MCF concept. The performance range for ITER (red hatched rectangle in the upper-right) is bounded by its projected peak triple product and its original stated mission goal ($Q_{\text{sci}} \geq 10$). ITER represents the crucial step required to bridge the gap between current experimental capabilities and the reactor-relevant regime in the European roadmap. Reproduced from [6].

for the qualification of magnetic configurations and prototypes. DTT, shown in Figure 1.4(a), is a superconducting tokamak characterized by a high toroidal magnetic field ($B_\phi = 5.85$ T) and a plasma current of up to $I_p = 5.5$ MA. Table 1.1 summarizes its main parameters, comparing them to those expected in ITER and EU-DEMO. A defining feature of the machine is its high auxiliary heating power, provided by a mix of ECRH, ICRH, and negative NBI systems, whose arrangement around the machine

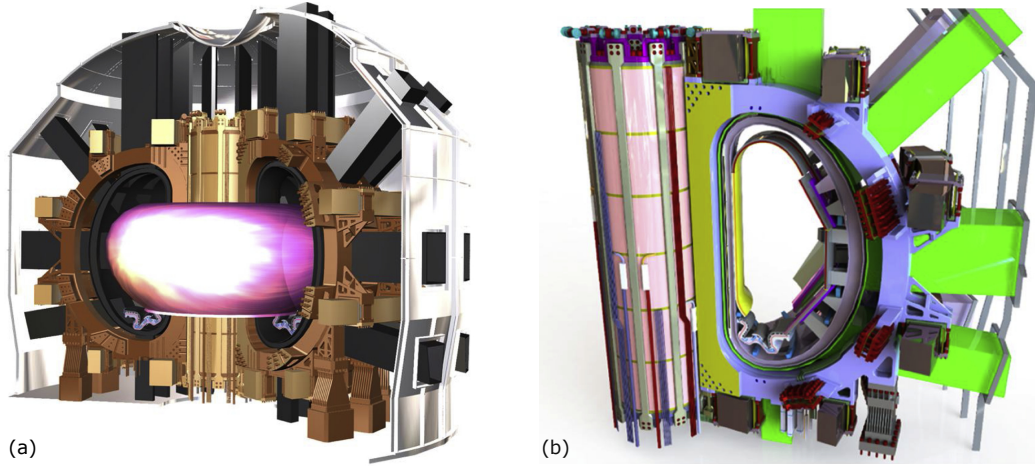


Fig. 1.4 (a) Sectional view of the DTT facility. (b) Overview of the DTT magnet system, comprising 18 Nb₃Sn toroidal field coils, a 6-module Nb₃Sn central solenoid, 6 external poloidal field coils (two Nb₃Sn and four NbTi), and 6 in-vessel copper coils. Adapted from [12].

sectors is illustrated in Figure 1.5. With a total of 45 MW, DTT can reproduce divertor heat loads comparable to those of ITER and EU-DEMO [10]. Indeed, the ratio of the power crossing the separatrix to the major radius (P_{sep}/R) is expected to reach ≈ 15 MW/m, a value similar to those of future reactor-class devices.

Figure 1.4(b) details the superconducting magnet system: the confinement is provided by 18 Nb₃Sn TF coils, a 6-module Nb₃Sn CS, 6 PF coils (two Nb₃Sn and four NbTi), and 6 in-vessel copper coils, designed for plasma control and local modifications of the magnetic configuration [12]. As shown in Figure 1.6, this magnet system allows for the realization of a wide range of magnetic topologies.

Parameter	Symbol	DTT	ITER	EU-DEMO
Major Radius	R [m]	2.19	6.2	9.1
Minor Radius	a [m]	0.70	2.0	2.93
Aspect Ratio	R/a	3.13	3.10	3.11
Toroidal Field	B_ϕ [T]	5.85	5.3	5.7
Plasma Current	I_p [MA]	5.5	15	19.6
Pulse Length	t_{pulse} [s]	100	400	7600
Heating Power	P_{aux} [MW]	45	150	460
Exhaust load	P_{sep}/R [MW/m]	15	14	17

Table 1.1 Comparison of the main parameters among DTT, ITER, and EU-DEMO [11].

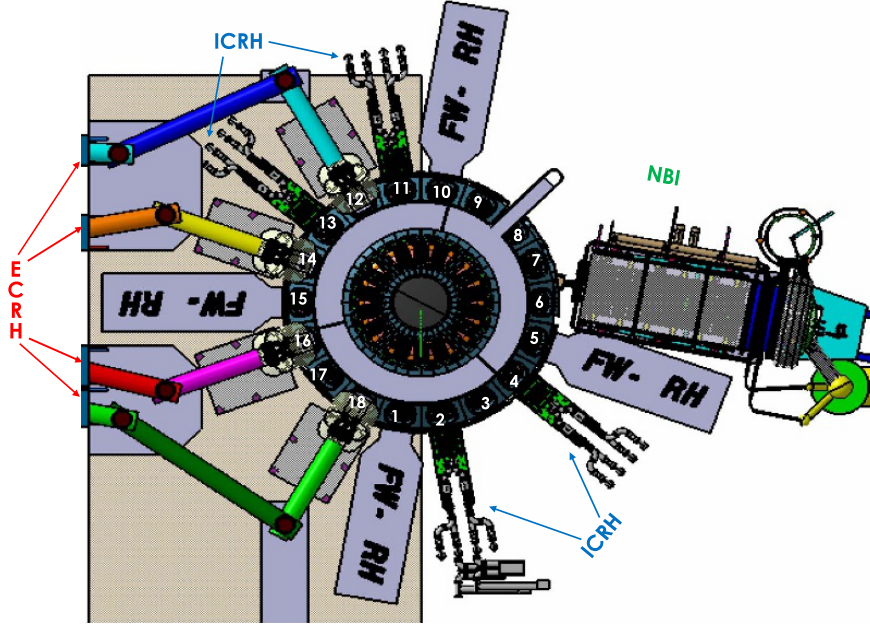


Fig. 1.5 Top-down view of the DTT heating system layout. The sectors are equipped with electron cyclotron resonance heating (ECRH) launchers, ion cyclotron resonance heating (ICRH) antennas, and a neutral beam injector (NBI). Ports reserved for the remote handling of the first wall (FW-RH) are also indicated. Adapted from [10].

These include the standard single null (SN) and alternative configurations such as the double null, the X-divertor, the snowflake, and the double super-X. Additionally, DTT is compatible with negative triangularity scenarios [12, 10]. Beyond magnetic shaping, DTT is also equipped to test alternative plasma-facing components (PFCs). While the primary divertor solution relies on solid tungsten, present plans also include testing liquid metal divertors, specifically capillary porous systems using liquid tin [10].

The experimental program of DTT will follow a staged approach, gradually increasing operational capabilities to ensure the safe commissioning of all subsystems. This roadmap is structured into three main phases [10]:

1. **Day-0** (commissioning phase): This initial phase, currently foreseen for 2029 [10], is dedicated to the first plasma operations and the commissioning of the main plant systems. It is characterized by reduced parameters, with a plasma current of $I_p = 2$ MA and a toroidal field of $B_\phi = 3$ T. In this stage, the coupled auxiliary heating power is approximately 17.4 MW, provided by the ECRH and ICRH systems.

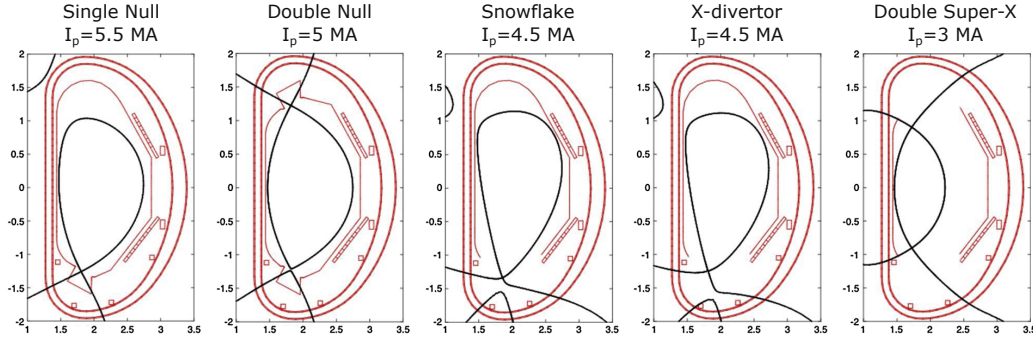


Fig. 1.6 Poloidal cross-sections of alternative magnetic configurations feasible in DTT thanks to the flexibility of the PF coil system: single null (SN), double null, snowflake, X-divertor, and double super-X. Adapted from [12].

2. **Day-1** (intermediate phase): Following the successful completion of the initial commissioning, the machine performance will be ramped up. This intermediate scenario targets a plasma current of $I_p = 4$ MA and the nominal toroidal field of 5.85 T. The heating power will be increased to ≈ 27 MW by introducing the NBI system.
3. **Full Power** (FP, nominal phase): This final stage represents the target operating point of the machine, designed to achieve reactor-relevant power exhaust conditions ($P_{\text{sep}}/R \approx 15$ MW/m). This scenario features the maximum plasma current of $I_p = 5.5$ MA, the nominal toroidal field, and the complete installation of the ≈ 45 MW auxiliary heating mix (with doubled ECRH and ICRH capacity compared to Day-0). Once this regime is established, its magnetic flexibility will be exploited to test alternative divertor configurations (Figure 1.6) and PFCs under reactor-relevant loads. Finally, the long-term roadmap envisages a subsequent phase dedicated to further upgrades to address emerging technological challenges on the path to EU-DEMO.

The predictive analysis presented in this thesis focuses on two specific milestones within this roadmap. The first is the Day-0 scenario. While thermal loads are not critical in this phase, it is of particular interest for safety analysis: it represents the starting point of machine operations, where the control systems are first tested on a live plasma, and the risk of disruptions must be carefully managed to protect PFCs. The second reference point is the FP scenario. With its plasma current of 5.5 MA, this regime corresponds to the highest stored magnetic energy and, consequently, the most severe risks in the event of RE generation and multiplication. Therefore,

assessing the safety margins of this scenario is essential for the success of the DTT mission.

1.3 Thesis motivation and outline

The demanding operational parameters of DTT, while necessary to replicate EU-DEMO-relevant exhaust conditions, introduce significant challenges regarding machine protection. In particular, the high plasma current and stored energy increase the potential severity of plasma *disruptions*, sudden events where the plasma thermal and magnetic energy is released on millisecond timescales. While disruptions impose severe thermal and electromechanical loads on the structural components, a specific and potentially more damaging phenomenon is the generation of runaway electrons (REs). During the rapid current quench (CQ) following a disruption, the strong induced toroidal electric field can accelerate electrons to relativistic velocities. Through the avalanche mechanism, an initial RE seed population can be amplified exponentially, potentially converting a substantial fraction of the pre-disruption plasma current into an RE beam. In high-current devices like DTT, the interaction of such high-energy beams with plasma-facing components can result in the localized deposition of tens of megajoules of energy, leading to deep melting and irreversible damage to the first wall and divertor targets.

For DTT, understanding the risks associated with REs is critical. The planned transition from the initial Day-0 scenario ($I_p = 2$ MA) to the reactor-relevant Full Power scenario ($I_p = 5.5$ MA) represents a fundamental shift in RE generation dynamics, since the avalanche gain scales exponentially with the plasma current. Thus, while in the commissioning phase the RE amplification is expected to be manageable, in the high-current regime the margin for error shrinks significantly. Therefore, a predictive assessment of RE generation is mandatory to define the safe operational boundaries in both scenarios and to inform the design of the DTT disruption mitigation system (DMS), optimizing strategies for RE avoidance and mitigation of generated beams.

This thesis aims to provide a comprehensive analysis of RE avalanche dynamics in DTT in its commissioning and full current phases (Day-0 and FP, respectively) by employing the non-linear MHD code JOREK coupled with the resistive wall code STARWALL. It is organized as follows:

- **Chapter 2** provides the theoretical background necessary to understand the physics of tokamak disruptions and REs. It begins with a description of MHD equations and stability principles, followed by a detailed analysis of disruption dynamics, including vertical displacement events (VDEs) and mitigation strategies. Finally, it covers the physics of REs, detailing their primary and secondary generation mechanisms.
- **Chapter 3** details the computational framework used in this study. It describes the MHD model implemented in JOEKE, the RE fluid model, and the coupling with the resistive wall code STARWALL. The *artificial thermal quench* (ATQ) procedure, used to initiate the disruptions in the simulations, is also explained.
- **Chapter 4** presents the results for the DTT Day-0 scenario ($I_p = 2$ MA). It analyzes the sensitivity of RE generation to initial seed currents and impurity levels, demonstrating that this early operational phase poses minimal RE risks.
- **Chapter 5** extends the analysis to the Full Power scenario ($I_p = 5.5$ MA). It highlights the dramatic increase in avalanche gain and identifies a narrow safety window, showing how even minute seed populations can lead to macroscopic RE beams in the presence of strong impurity cooling. The chapter concludes with a general assessment that synthesizes the findings from both scenarios and characterizes the risk profile for RE generation.
- **Chapter 6** summarizes the main scientific achievements and outlines future research directions.

Chapter 2

Physical background

The design and operation of a tokamak rely on the precise control of a high-temperature plasma confined by magnetic fields. To understand the macroscopic behavior of this complex system and, more specifically, the catastrophic loss of confinement known as a disruption, it is necessary to establish a theoretical framework that describes the interaction between the conducting fluid and the electromagnetic fields.

This chapter outlines the fundamental physical models used throughout this thesis. Section 2.1 introduces the magnetohydrodynamic (MHD) description of the plasma, focusing on the resistive MHD equations and the energy principle that governs stability. Section 2.2 details the phenomenology of tokamak disruptions, analyzing the sequence of events from the thermal quench (TQ) to the current quench (CQ), including the risks associated with vertical displacement events (VDEs) and the strategies employed for mitigation. Finally, Section 2.3 addresses the physics of runaway electrons (REs). It establishes the theoretical thresholds for the onset of the runaway regime, describes the primary and secondary generation mechanisms, and discusses the transport processes that determine their final impact on the plasma-facing components (PFCs).

2.1 Magnetohydrodynamics (MHD)

To describe the macroscopic behavior of a fusion plasma, one must bridge the gap between the complex kinetic behavior of individual charged particles and the collective dynamics of the system. The most successful and widely used framework

for this purpose is magnetohydrodynamics (MHD). It treats the plasma not as a collection of discrete ions and electrons, but as a single, electrically conducting fluid interacting with an electromagnetic (EM) field. Although the rigorous description of a plasma requires a kinetic approach (solving the Boltzmann equation for the distribution function of each species), such a treatment is often computationally prohibitive for global scale phenomena. However, for the low-frequency, long-wavelength dynamics characteristic of macroscopic equilibrium and stability, the MHD fluid approximation is remarkably effective.

In the context of this thesis, the MHD formalism provides the necessary tools to define the tokamak equilibrium and, crucially, to identify the instabilities that lead to disruptions. The following subsections outline the governing equations of the resistive MHD model (Section 2.1.1), which forms the basis of the JOREK code, and the fundamental principles of stability analysis (Section 2.1.2). The theoretical derivations and limits of validity presented in this section primarily follow the standard literature on MHD provided in [13, 14].

2.1.1 Resistive MHD equations

The resistive MHD model serves as a single-fluid approximation of plasma dynamics, bridging the gap between the precise but computationally demanding kinetic theory and the macroscopic behavior. The derivation of this model begins with the fundamental description of the plasma in the phase space, provided by the Boltzmann equation for the distribution function $f_\alpha(\mathbf{r}, \mathbf{v}, t)$ of each species α (ions and electrons):

$$\frac{\partial f_\alpha}{\partial t} + \mathbf{v} \cdot \nabla f_\alpha + \frac{q_\alpha}{m_\alpha} (\mathbf{E} + \mathbf{v} \times \mathbf{B}) \cdot \nabla_v f_\alpha = \sum_\beta C_{\alpha\beta} \equiv C_\alpha \equiv \left(\frac{\partial f_\alpha}{\partial t} \right)_{\text{coll}} \quad (2.1)$$

where the term on the right-hand side represents the collision operator describing short-range Coulomb interactions. The summation over β accounts for collisions with all species, including like-particle collisions ($\beta = \alpha$). To provide a complete physical description, this kinetic equation must be solved self-consistently with Maxwell's equations. The coupling arises because the EM fields determine the particle trajectories (via the Lorentz force, third term on the left-hand side), while the particles simultaneously act as source terms for the fields through their collective

charge density σ and current density $\mathbf{J}$:

$$\sigma = \sum_{\alpha} q_{\alpha} \int f_{\alpha} d\mathbf{v}, \quad \mathbf{J} = \sum_{\alpha} q_{\alpha} \int \mathbf{v} f_{\alpha} d\mathbf{v}$$

The evolution of the electric and magnetic fields is then governed by the full set of Maxwell's equations:

$$\nabla \times \mathbf{B} = \mu_0 \mathbf{J} + \frac{1}{c^2} \frac{\partial \mathbf{E}}{\partial t} \quad (2.2)$$

$$\nabla \times \mathbf{E} = -\frac{\partial \mathbf{B}}{\partial t} \quad (2.3)$$

$$\nabla \cdot \mathbf{E} = \frac{\sigma}{\epsilon_0} \quad (2.4)$$

$$\nabla \cdot \mathbf{B} = 0 \quad (2.5)$$

Solving this system of equations is computationally prohibitive because it requires evolving the distribution function in a six-dimensional phase space $(\mathbf{r}, \mathbf{v})$ over time t . To overcome this, we reduce the complexity by taking velocity moments of Equation (2.1), effectively reducing the independent variables to spatial coordinates and time $(\mathbf{r}, t)$. This process transforms the single kinetic equation for f_{α} into an infinite hierarchy of coupled fluid equations. In practice, this hierarchy is truncated by introducing closure assumptions appropriate for the physical regime of interest. While this truncation inevitably sacrifices the detailed phase-space information contained in the original kinetic model, it allows for a solvable set of macroscopic evolution equations. Explicitly, this is done by multiplying Equation (2.1) by powers of velocity $(1, m_{\alpha} \mathbf{v}, m_{\alpha} v^2/2)$ and integrating over all velocity space, yielding the multiple-fluid conservation equations, which describe the evolution of macroscopic variables for each species: the particle density n_{α} , the mean fluid velocity $\mathbf{u}_{\alpha}$, and the temperature T_{α} (or pressure p_{α}). n_{α} and $\mathbf{u}_{\alpha}$ are defined as the zeroth and first velocity moments of the distribution function, respectively:

$$n_{\alpha} = \int f_{\alpha} d\mathbf{v}, \quad \mathbf{u}_{\alpha} = \frac{1}{n_{\alpha}} \int \mathbf{v} f_{\alpha} d\mathbf{v}$$

To define higher-order moments, such as pressure and temperature, it is useful to separate the particle velocity $\mathbf{v}$ into the mean flow $\mathbf{u}_{\alpha}(\mathbf{r}, t)$ and the random thermal motion $\mathbf{w}$: $\mathbf{v} \equiv \mathbf{u}_{\alpha}(\mathbf{r}, t) + \mathbf{w}$. By definition, the average of the random velocity vanishes,

$\langle \mathbf{w} \rangle = 0$, where the bracket notation denotes the average over the distribution function:

$$\langle h \rangle = \frac{1}{n_\alpha} \int h f_\alpha d\mathbf{v}$$

We can now explicitly define the thermodynamic and transport quantities of interest as moments of the random velocity $\mathbf{w}$. These include the total pressure tensor $\underline{\mathbf{P}}_\alpha$, the scalar pressure p_α and the anisotropic stress tensor $\underline{\Pi}_\alpha$, and the temperature T_α :

$$\begin{aligned} \underline{\mathbf{P}}_\alpha &= m_\alpha n_\alpha \langle \mathbf{w} \mathbf{w} \rangle, & p_\alpha &= \frac{1}{3} m_\alpha n_\alpha \langle w^2 \rangle \\ \underline{\Pi}_\alpha &= \underline{\mathbf{P}}_\alpha - p_\alpha \underline{\mathbf{I}}, & T_\alpha &= p_\alpha / n_\alpha \end{aligned}$$

where $\underline{\mathbf{I}}$ is the identity matrix. The heat flux vector $\mathbf{q}_\alpha$ represents the transport of energy due to random thermal motion:

$$\mathbf{q}_\alpha = \frac{1}{2} m_\alpha n_\alpha \langle w^2 \mathbf{w} \rangle$$

while the collision terms describing the interaction between different species (friction force $\mathbf{R}_\alpha$ and collisional heat generation Q_α) are defined as moments of the collision operator C_α :

$$\mathbf{R}_\alpha = m_\alpha \int \mathbf{w} C_\alpha d\mathbf{w}, \quad Q_\alpha = \frac{1}{2} m_\alpha \int w^2 C_\alpha d\mathbf{w}$$

Exploiting these definitions, the resulting system of equations for species α reads:

$$\left(\frac{dn_\alpha}{dt} \right)_\alpha + n_\alpha \nabla \cdot \mathbf{u}_\alpha = 0 \quad (2.6)$$

$$m_\alpha n_\alpha \left(\frac{d\mathbf{u}_\alpha}{dt} \right)_\alpha - q_\alpha n_\alpha (\mathbf{E} + \mathbf{u}_\alpha \times \mathbf{B}) + \nabla \cdot \underline{\mathbf{P}}_\alpha = \mathbf{R}_\alpha \quad (2.7)$$

$$\frac{3}{2} n_\alpha \left(\frac{dT_\alpha}{dt} \right)_\alpha + \underline{\mathbf{P}}_\alpha : \nabla \mathbf{u}_\alpha + \nabla \cdot \mathbf{q}_\alpha = Q_\alpha \quad (2.8)$$

In these equations, the time evolution is described using the convective derivative (or material derivative) associated with species α , defined as:

$$\left(\frac{d}{dt} \right)_\alpha \equiv \frac{\partial}{\partial t} + \mathbf{u}_\alpha \cdot \nabla \quad (2.9)$$

This operator represents the rate of change of a quantity moving with the local fluid velocity $\mathbf{u}_\alpha$. The EM fields continue to evolve according to Maxwell's equations introduced before (Equations (2.2)–(2.5)); however, the source terms σ and $\mathbf{J}$ are now explicitly defined as:

$$\sigma = \sum_{\alpha} q_{\alpha} n_{\alpha}, \quad \mathbf{J} = \sum_{\alpha} q_{\alpha} n_{\alpha} \mathbf{u}_{\alpha}$$

To transition from the multiple-fluid description to the single-fluid MHD model, we form linear combinations of the equations for the various species, assuming a simplified composition of electrons and singly charged ions. This reduction relies on two critical physical assumptions. First, we assume quasi-neutrality ($n_i \approx n_e \equiv n$) on length scales much larger than the Debye length and frequencies much smaller than the electron plasma frequency ω_{pe} , eliminating the need to solve Poisson's equation. Additionally, at such low frequencies ($\omega \ll \omega_{pe}, \omega_{ce}$, where ω_{ce} is the electron cyclotron frequency), the displacement current in the Ampère–Maxwell law can also be neglected, simplifying it to $\nabla \times \mathbf{B} = \mu_0 \mathbf{J}$. Finally, due to the large mass difference between electrons and ions ($m_e/m_i \ll 1$), the electron mass is neglected. This implies that, on MHD timescales, electrons have an infinitely fast response, and fluid inertia is dominated by the ions. Therefore, we define the total mass density and the single-fluid velocity as:

$$\rho \approx m_i n, \quad \mathbf{v} \equiv \frac{n_e m_e \mathbf{u}_e + n_i m_i \mathbf{u}_i}{n_e m_e + n_i m_i} \approx \mathbf{u}_i$$

The thermodynamic state of the plasma is described by the total pressure p and the single-fluid temperature T . The former is the sum of the partial pressures of ions and electrons, while the latter is the average of the species' temperatures:

$$p = p_i + p_e = 2nT, \quad T = \frac{T_i + T_e}{2}$$

The current density $\mathbf{J}$ is defined as the net flux of charge, which for singly charged ions and under the assumption of quasi-neutrality is:

$$\mathbf{J} \approx en(\mathbf{u}_i - \mathbf{u}_e)$$

This allows us to express the electron fluid velocity $\mathbf{u}_e$ and the electron friction force $\mathbf{R}_e$ and collisional heat generation Q as:

$$\mathbf{u}_e \approx \mathbf{v} - \frac{\mathbf{J}}{en}, \quad \mathbf{R}_e = -\mathbf{R}_i \approx en\eta\mathbf{J}, \quad Q = Q_i + Q_e \approx (\mathbf{u}_i - \mathbf{u}_e) \cdot \mathbf{R}_e \approx \eta J^2$$

where η is the electrical resistivity. Finally, in order to close the system, we can adopt standard approximations for the transport terms:

$$\underline{\Pi} \equiv \underline{\Pi}_i + \underline{\Pi}_e \equiv -\underline{\tau} \approx -\underline{\nu}\nabla\mathbf{v}, \quad \mathbf{q} \approx -\underline{\kappa}\nabla T \equiv -n\underline{\chi}\nabla T$$

where $\underline{\tau}$ is the viscous stress tensor, $\underline{\nu}$ is the kinematic viscosity, and $\underline{\kappa}$ represents the thermal conductivity (related to the thermal diffusivity $\underline{\chi}$ by $\underline{\kappa} = n\underline{\chi}$). In magnetized plasmas, thermal diffusion is highly anisotropic, meaning that transport along magnetic field lines is orders of magnitude faster than transport across them ($\chi_{\parallel} \gg \chi_{\perp}$). Lastly, γ is the ratio of specific heats, which for plasmas is typically equal to 5/3.

The resistive MHD model, obtained by combining electron and ion fluid equations and applying the closure approximations discussed above, is given by:

$$\frac{\partial \rho}{\partial t} + \nabla \cdot (\rho \mathbf{v}) = 0 \quad (2.10)$$

$$\rho \frac{\partial \mathbf{v}}{\partial t} + \rho(\mathbf{v} \cdot \nabla)\mathbf{v} + \nabla p - \mathbf{J} \times \mathbf{B} = \nabla \cdot (\underline{\nu}\nabla\mathbf{v}) \quad (2.11)$$

$$\frac{\partial p}{\partial t} + \mathbf{v} \cdot \nabla p + \gamma p \nabla \cdot \mathbf{v} = (\gamma - 1) \left(\eta J^2 + \nabla \cdot (n\underline{\chi}\nabla T) + \underline{\tau} : \nabla\mathbf{v} \right) \quad (2.12)$$

together with the generalized Ohm's law and the low-frequency approximation of Maxwell's equations:

$$\mathbf{E} + \mathbf{v} \times \mathbf{B} = \eta\mathbf{J} + \frac{\mathbf{J} \times \mathbf{B}}{en} - \frac{\nabla \cdot \mathbf{P}_e}{en} \quad (2.13)$$

$$\nabla \times \mathbf{B} = \mu_0 \mathbf{J} \quad (2.14)$$

$$\nabla \times \mathbf{E} = -\frac{\partial \mathbf{B}}{\partial t} \quad (2.15)$$

$$\nabla \cdot \mathbf{B} = 0 \quad (2.16)$$

The generalized Ohm's law (Equation (2.13)) retains terms describing Hall currents ($\mathbf{J} \times \mathbf{B}$) and electron diamagnetic drift ($\mathbf{P}_e$). However, for large-scale fusion experiments, where the characteristic length scales are much larger than the ion gyro radius, these two terms become negligible. Dropping them yields the standard resistive Ohm's law, $\mathbf{E} + \mathbf{v} \times \mathbf{B} = \eta \mathbf{J}$.

Equations (2.10)–(2.13) effectively show a separation between ideal MHD equations (left-hand side) and dissipative terms including viscosity, resistivity, and thermal conduction (right-hand side). In ideal MHD dynamics, the magnetic flux is *frozen* into the plasma, and the magnetic topology is preserved. Crucially, in resistive MHD the finite η relaxes the frozen-in constraint and enables *magnetic reconnection*, which is the fundamental mechanism behind tearing modes and disruptions investigated here.

It is worth noting a fundamental paradox regarding the validity of the MHD model. The single-fluid approximation assumes high plasma collisionality, which implies local thermodynamic equilibrium and ensures that ions and electrons share a common temperature. Physically, this requires the temperature equilibration time to be much shorter than the characteristic timescale of the macroscopic dynamics. In fusion-relevant plasmas, however, this condition is almost never satisfied, as the plasma is typically in a collisionless regime. Under such circumstances, ion and electron temperatures decouple, theoretically necessitating a two-fluid description or an extended MHD model with separate temperature evolutions. Despite this, the standard single-fluid MHD model remains remarkably robust in describing macroscopic stability. This is possible because the strong magnetic field in fusion experiments strictly confines particle motion perpendicular to $\mathbf{B}$. This renders the perpendicular dynamics effectively fluid-like regardless of the collision frequency, thus justifying the use of standard MHD for global stability analysis. Additionally, when the plasma can be assumed to be incompressible ($\nabla \cdot \mathbf{v} = 0$), the MHD equations suitably describe the parallel dynamics as well.

Finally, while these equations represent the baseline theoretical model, practical simulations often require extensions that incorporate, in a simplified manner, effects neglected in the standard derivation. For instance, the continuity equation is typically augmented with diffusion and particle source terms, or MHD equations are modified to account for the above-mentioned separation between ion and electron temperatures. Furthermore, the MHD model describes only the thermal bulk of the plasma and

cannot capture high-energy non-thermal populations. Since this thesis focuses on runaway electrons, this standard MHD framework serves as the background which must be extended with appropriate RE equations, as detailed in Section 2.3 and 3.2.

2.1.2 MHD stability and the energy principle

The MHD equations described in the previous section allow for the determination, in a given configuration, of a plasma equilibrium, defined as a stationary state where the forces acting on the plasma are perfectly balanced. In the static limit ($\mathbf{v} = 0$), this reduces to the well-known force balance $\mathbf{J} \times \mathbf{B} = \nabla p$, where the magnetic force balances the pressure gradient. However, the existence of such an equilibrium is a necessary, but not sufficient, condition for the practical confinement of a fusion plasma: it must also be stable against small perturbations, since they are unavoidable. If the perturbed forces are such that they restore the original equilibrium, the plasma is stable; if, instead, they enhance the original perturbation, it is unstable, potentially leading to a plasma disruption.

Stability analysis is typically performed using the linearized ideal MHD equations. While one could directly solve the equation of motion to find the growth rates of instabilities, a more intuitive and powerful approach is the *energy principle* [14]. This method determines stability by examining the perturbed potential energy of the plasma, δW , resulting from an arbitrary displacement from equilibrium $\xi(\mathbf{r})$. The equilibrium is stable if and only if:

$$\delta W \equiv -\frac{1}{2} \int \xi \cdot F(\xi) d\mathbf{r} \geq 0$$

for all physically admissible displacements ξ . The force operator $\mathbf{F}(\xi)$ is given by Newton's second law ($\rho \ddot{\xi} = \mathbf{F}(\xi)$). Conversely, if there exists a single ξ that lowers the potential energy ($\delta W < 0$), it is energetically favorable for the system to move away from equilibrium, and the related instability will take place. To understand the physical drivers of these instabilities, it is useful to decompose the perturbed potential energy into three distinct contributions:

$$\delta W = \delta W_F + \delta W_S + \delta W_V$$

which represent, respectively, the fluid, surface, and vacuum contributions.

The vacuum energy perturbation δW_V represents the change in magnetic energy in the vacuum region surrounding the plasma, and it is given by:

$$\delta W_V = \int_{V_{\text{vac}}} \frac{|\hat{\mathbf{B}}_1|^2}{2\mu_0} d\mathbf{r}$$

where $\hat{\mathbf{B}}_1$ is the perturbed magnetic field in the vacuum volume outside of the plasma (V_{vac}). Since $\delta W_V \geq 0$, the vacuum response is always stabilizing. When the plasma boundary S_p moves, external magnetic field lines get distorted, requiring work to bend or compress the magnetic flux in the vacuum region. While this stabilizing effect is stronger in the presence of a perfectly conducting wall placed close to the plasma, even in the presence of a non-ideal wall this term remains positive and stabilizing.

The surface energy perturbation δW_S arises from the discontinuity of the total pressure gradient across the plasma-vacuum interface S_p . It is explicitly given by:

$$\delta W_S = \frac{1}{2} \int_{S_p} |\mathbf{n} \cdot \xi_\perp|^2 \mathbf{n} \cdot \left[\nabla \left(\frac{B^2}{2\mu_0} + p \right) \right] dS$$

where $\mathbf{n}$ is the outward unit normal vector and the double brackets $[\![\cdot]\!]$ denote the jump of a quantity across the interface. This term contributes to the potential energy only in the presence of equilibrium surface currents at the boundary. Consequently, for profiles where the current density vanishes smoothly at the edge, this term is zero. Unlike the vacuum energy, however, it can generally be either positive or negative.

Finally, the fluid energy perturbation δW_F represents the potential energy change within the plasma volume V_p and contains stabilizing (always positive) and possibly destabilizing (negative) terms. It is explicitly written as:

$$\delta W_F = \frac{1}{2\mu_0} \int_{V_p} \left[\underbrace{|\mathbf{Q}_\perp|^2 + B^2 |\nabla \cdot \xi_\perp + 2\xi_\perp \cdot \kappa|^2 + \mu_0 \gamma p |\nabla \cdot \xi|^2}_{\text{Stabilizing terms}} \right. \\ \left. \underbrace{- 2\mu_0 (\xi_\perp \cdot \nabla p) (\xi_\perp^* \cdot \kappa)}_{\text{Pressure-driven instabilities}} \right. \\ \left. \underbrace{- \mu_0 J_\parallel (\xi_\perp^* \times \mathbf{b}) \cdot \mathbf{Q}_\perp (\xi_\perp)}_{\text{Current-driven instabilities}} \right] d\mathbf{r}$$

where the subscripts $\parallel$ and $\perp$ respectively denote vector components parallel and perpendicular to $\mathbf{B}$, $\mathbf{Q} = \nabla \times (\xi \times \mathbf{B})$ is the perturbed magnetic field, $\kappa = \mathbf{b} \cdot \nabla \mathbf{b}$ is the curvature vector (pointing in the direction of the center of curvature of the magnetic field lines), $\mathbf{b} = \mathbf{B}/B$ is the unit vector along $\mathbf{B}$, and ξ^* is the complex conjugate of ξ (since the force operator $\mathbf{F}(\xi)$ is self-adjoint, δW_F is always real). The first three terms in the volume integral are always positive and, hence, stabilizing. They represent, respectively, the energy cost associated with bending magnetic field lines (shear Alfvén wave), compressing the magnetic field (compressional Alfvén wave), and compressing the plasma fluid (sound wave). Instabilities (or *modes*) can arise solely from the last two terms:

- **Pressure-driven instabilities:** These modes are driven by the pressure gradient ∇p and include *interchange instabilities*, similar to Rayleigh-Taylor instabilities, and *ballooning modes*, which impose a limit on the achievable plasma β (the ratio between the plasma and the magnetic pressures). Compared to current-driven instabilities, they are less likely to cause major disruptions, but they drive enhanced turbulent transport.
- **Current-driven instabilities:** These modes are driven by the current density parallel to the magnetic field $J_{\parallel}$. They are known as *kink modes* because they induce a macroscopic twisting of the plasma column and are classified according to the extent of the mode structure. *Internal kink modes* are confined within the plasma volume (leaving the boundary fixed) and are typically less catastrophic, manifesting primarily as periodic relaxation events of the core temperature and current profiles known as *sawtooth oscillations*. Conversely, *external kink modes* involve a displacement of the plasma-vacuum interface S_p and are particularly dangerous, often leading to major disruptions and setting severe limits on β .

The stability analysis provided by the energy principle relies on the ideal MHD approximation, which assumes a perfectly conducting fluid ($\eta = 0$), where magnetic topology cannot change. However, in realistic plasmas with finite resistivity, this constraint is relaxed, allowing magnetic field lines to break and reconnect in *magnetic islands*, giving rise to *resistive instabilities* such as *tearing modes*. These modes typically resonate on rational flux surfaces (where the safety factor q can be expressed as a rational number m/n) and constitute a significant threat to stable operation:

by forming magnetic islands that break the nested flux surfaces, they facilitate the sudden release of the plasma's stored energy, potentially resulting in a disruption.

Finally, the linear approach discussed here is limited to determining the onset of an instability, but it cannot predict the long-term evolution of the system far from equilibrium. To distinguish whether a mode will saturate into a benign quasi-stationary state or grow explosively into a disruption, a non-linear treatment is required. Furthermore, this treatment is essential to capture how the interaction with conductive structures can decelerate instabilities to the characteristic diffusion time of the wall. Accurately modeling these complex dynamics requires non-linear codes such as JOREK (Section 3.1).

2.2 Disruptions

Disruptions are a consequence of MHD instabilities and are defined as sudden, uncontrolled losses of global plasma confinement, resulting in the complete termination of the discharge [15]. During such an event, the majority of the energy stored in the plasma (both thermal and magnetic) is released on a millisecond timescale. In reactor-scale devices like ITER and EU-DEMO, the magnitude of these energies poses a severe threat to the machine integrity, and therefore must be avoided or mitigated. The rapid energy release during a plasma disruption drives three distinct damage mechanisms:

- **Thermal loads:** The rapid loss of thermal energy results in extreme heat fluxes on PFCs, potentially causing their ablation and melting.
- **Electromagnetic forces:** The decay of the plasma current induces *eddy currents* in the surrounding conductive structures. The interaction of these currents with the background magnetic field generates massive $\mathbf{J} \times \mathbf{B}$ forces that can deform the vacuum vessel and structural supports.
- **Runaway electrons:** The high toroidal electric field induced during the current decay can accelerate electrons to relativistic energies, creating high-energy beams capable of causing localized, deep damage to in-vessel components. Since this is the main focus of this thesis, it will be analyzed in more detail in the following sections.

It is important to distinguish between different degrees of confinement loss based on the severity of the instability. A *minor disruption* is characterized by a partial or complete loss of thermal energy from which the magnetic confinement is preserved, allowing the discharge to potentially recover. In contrast, a *major disruption* (the primary focus of this work) is an irreversible event where the loss of thermal energy is immediately followed by the complete collapse of the plasma current.

To provide the necessary context for the RE analysis presented in later chapters, it is essential to first understand the macroscopic evolution of disruptions. Section 2.2.1 details their physical triggers and their dynamics (thermal and current quenches). Section 2.2.2 focuses on the specific risks associated with the loss of vertical control, including vertical displacement events (VDEs) and halo currents. Finally, Section 2.2.3 reviews the disruption mitigation strategies that could alleviate these loads.

2.2.1 Disruption causes and dynamics

Disruptions are typically the consequence of crossing operational boundaries or of technical failures. The primary operational limits responsible for triggering disruptions [15] are:

- **The density limit:** The density should not exceed the Greenwald limit $n_{\text{GW}}(10^{20} \text{ m}^{-3}) = I_p / (\pi a^2)$, where I_p is in MA and a in m.
- **The q -limit:** The safety factor calculated at the edge (q_{95}) typically must remain above 2.
- **The β -limit:** The normalized toroidal β is constrained by the Troyon ideal MHD limit (3.5% MA $\text{m}^{-1} \text{ T}^{-1}$).

While the q and β limits imply immediate ideal MHD instabilities, the density limit underpins a deterioration of confinement and may indirectly initiate MHD instabilities close to the $q = 2$ surface. Although modern control systems aim to avoid these boundaries, the necessity of operating near high-performance limits to maximize fusion gain Q makes disruptions an inherent risk.

The onset of a major disruption is then followed by a multi-stage process driven by the non-linear evolution of MHD instabilities. Although the specific triggering

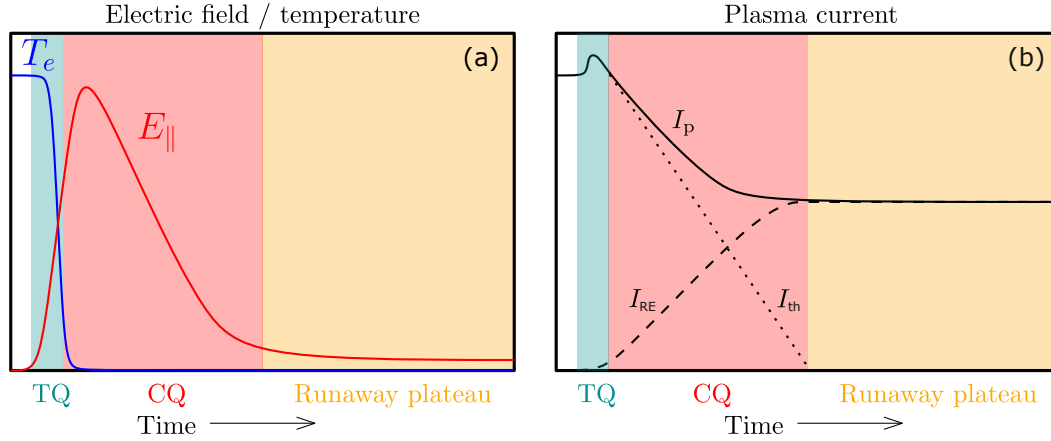


Fig. 2.1 Schematic evolution of plasma parameters during a disruption involving runaway electron generation. (a) The electron temperature T_e (blue line) collapses during the thermal quench phase (TQ, green shaded area), causing a rapid increase in resistivity and the induction of a large parallel electric field $E_{||}$ (red line) during the current quench phase (CQ, red shaded area). (b) The total plasma current I_p (solid line) shows the characteristic initial spike in the TQ followed by a decay. If a macroscopic RE beam is present, the decaying thermal current I_{th} (dotted line) is replaced by the relativistic RE current I_{RE} (dashed line), potentially leading to a runaway plateau (yellow shaded area). Adapted from [16].

mechanisms vary and there may be a precursor phase, the macroscopic sequence of events is commonly characterized by two distinct phases: the *thermal quench* (TQ) and the *current quench* (CQ):

- **Thermal quench:** During this phase (green shaded area in Figure 2.1), most of the core thermal energy of the plasma (W_{th}) is quickly lost and the plasma temperature T_e drops by several orders of magnitude, typically falling from keV levels to just ≈ 10 eV (blue line in Figure 2.1(a)). While the exact disruption trigger may vary (e.g., impurity radiation or MHD instabilities), the loss mechanism typically involves the stochastization of magnetic flux surfaces, which allows rapid parallel transport to exhaust the core energy. This energy is lost to the divertor and first wall (which experience a massive heat flux spike) on a timescale that scales with the machine cross-section, but which can be as short as ≈ 1 ms for devices like ITER [15]. Simultaneously, the current density profile flattens because of toroidal current redistribution, leading to a decrease in the plasma internal inductance l_i . As a consequence of poloidal flux conservation, the total plasma current I_p experiences a *current spike* of about 10% of the pre-disruption value [15] (Figure 2.1(b)).

- **Current quench:** Following the TQ, the temperature results from a balance between radiation losses and Ohmic heating. Radiation losses are dominated by impurities that can originate from plasma-wall interactions or be injected as part of a disruption mitigation strategy (Section 2.2.3). Generally, the plasma remains at tens of eV; in this cold state, it is highly resistive, since the Spitzer electrical resistivity scales as $\eta \propto T_e^{-3/2}$. This high resistivity causes the plasma current I_p to decay, dissipating the stored magnetic energy (W_{mag}) into Ohmic heating and radiation, while inducing eddy currents and EM forces on the surrounding structures. The duration of this phase (red shaded area in Figure 2.1) is on the order of tens of milliseconds. Crucially for this thesis, the rapid change in magnetic flux during the CQ induces a strong toroidal electric field ($E_{\parallel}$) in accordance with Faraday's law (red line in Figure 2.1(a)). It is this large electric field that is responsible for the acceleration of electrons into the runaway regime. The formation of a macroscopic beam of REs is not inevitable, as it occurs only if a sufficient initial RE seed population is present and the electric field forces overcome the collisional friction. In such scenarios, while the thermal current I_{th} (driven by thermal electrons) decays (dotted line in Figure 2.1(b)), the RE current I_{RE} grows (dashed line in Figure 2.1(b)). This may lead to the formation of a *runaway plateau* (yellow shaded area in Figure 2.1), where the RE beam effectively carries the entire instantaneous plasma current ($I_p \approx I_{\text{RE}}$). Depending on the specific scenario, the RE current can eventually carry a highly significant fraction of the pre-disruption plasma current (often up to 70% [15], although full conversion has been observed in devices such as the TCV tokamak [17, 18]). The physics governing these generation and transport mechanisms is detailed in Section 2.3.

2.2.2 Vertical displacement events and halo currents

Tokamaks typically operate with elongated plasma cross-sections, in order to increase the plasma performance by increasing the energy confinement time and improving the β -limit. However, this elongation naturally renders the plasma vertically unstable with respect to the $n = 0$ axisymmetric mode [19]. Since the elongation is produced by PF coils that carry currents in the same direction as the plasma, an attractive force exists between them. If the plasma undergoes a small vertical displacement dz from its equilibrium position, it moves closer to one set of coils and further from

the other, resulting in a net $\mathbf{J} \times \mathbf{B}$ force that accelerates the plasma in the direction of the perturbation, making the equilibrium unstable. Quantitatively, the vertical component of this destabilizing force F_z is proportional to the plasma current I_p and the gradient of the radial component of the external magnetic field $B_{r,\text{vac}}$ [1]:

$$F_z \propto I_p \frac{\partial B_{r,\text{vac}}}{\partial z} dz \quad (2.17)$$

During normal operation, this instability is mitigated by the presence of conductive structures surrounding the plasma (e.g., vacuum vessel, stabilizing coils). As the plasma moves, the changing magnetic flux induces eddy currents in these structures, which, in turn, generate a magnetic field that opposes the plasma motion, providing a stabilizing force $F_{z,\text{stab}}$:

$$F_{z,\text{stab}} \propto I_p^2 \left(\frac{dM_{\text{pw}}}{dz} \right)^2 dz \quad (2.18)$$

where M_{pw} represents the mutual inductance between the plasma and the conductive structures [20, 1]. However, these stabilizing currents decay on the resistive timescale of the wall $\tau_{L/R}$ up to hundreds of milliseconds, making this passive stabilization insufficient for steady-state operation and requiring an active feedback control system to maintain the vertical position [21]. This is feasible as long as the plasma current is high enough [22]. A vertical displacement event (VDE) occurs when this vertical control is lost, either due to a system failure or because the destabilizing force exceeds the maximum controllable force. When this happens, the plasma column experiences uncontrolled motion toward the PFCs (see Figure 2.2).

Depending on the causal sequence between the thermal quench and the vertical motion, VDEs are classified into two categories:

- **Hot VDE:** This event is initiated by a loss of vertical control while the plasma retains its thermal energy. As the plasma column is displaced (e.g., due to a control system failure), it is scraped off by the wall, reducing its minor radius a . Since the plasma current I_p initially remains roughly constant during the displacement, the edge safety factor q_{95} , which scales as a^2/I_p , decreases rapidly [23]. Because of the q -limit discussed in Section 2.2.1, when q_{95} drops below 2 it destabilizes low- n current-driven MHD modes. Further MHD modes are triggered by the sharp gradients at the edge [24]. The growth of these instabilities destroys the magnetic topology of the core plasma and triggers the TQ. Crucially, because the plasma impacts the wall while retaining high

current and temperature, this scenario generates the most severe EM forces and thermal loads. Consequently, the avoidance of hot VDEs is a mandatory requirement for the safe operation of a reactor-scale device.

- **Cold VDE:** In this scenario, also known as *vertically unstable disruption* [15], the causal sequence is effectively reversed compared to a hot VDE. The event begins with a major disruption (triggered by standard operational limits or MHD instabilities) that exhausts the plasma's thermal energy in the TQ phase. Vertical stability is lost only during the CQ phase, once the plasma has become cold and resistive and the current begins to decay. This instability arises because the stabilizing force $F_{z,\text{stab}}$ scales as I_p^2 , whereas the destabilizing force F_z scales linearly with I_p . Therefore, as the current drops, the restoring force vanishes faster than the destabilizing one, accelerating the plasma towards the wall. In the limit of fast current quenches ($\tau_{\text{CQ}} \ll \tau_{\text{L/R}}$), the vertical position becomes strictly a function of I_p [25], implying that wall contact consistently occurs at the same current level [26]. Sustaining this cold vertical displacement typically requires the presence of impurities to prevent the plasma from reheating, which would otherwise alter the CQ dynamics [27]. The defining distinction from hot VDEs is that the plasma loses its thermal energy before directly interacting with the wall; therefore, the localized thermal loads are generally lower than those observed in hot VDEs. However, significant halo currents and vertical forces can still be generated [15].

The dynamics of VDEs are of particular importance for the physics of REs. If a macroscopic RE beam is formed, it typically carries a significant fraction of the current and becomes the dominant current channel (runaway plateau in Figure 2.1). In the presence of such an RE beam, the dynamics of the VDE is fundamentally altered, as the plasma current stops decaying due to the negligible resistivity of the relativistic electrons. The subsequent VDE evolution determines the final impact location of the RE beam on the PFCs, dictating where the massive kinetic energy of the REs will be deposited. Furthermore, the shrinking of the plasma cross-section during a VDE influences the safety factor profile q , which in turn affects the stability of the RE beam against kinetic instabilities. These mechanisms will be detailed in Section 2.3.

As the plasma column moves during a VDE, the open magnetic field lines of the scrape-off layer (SOL, Figure 1.2) intersect the conductive first wall. When the

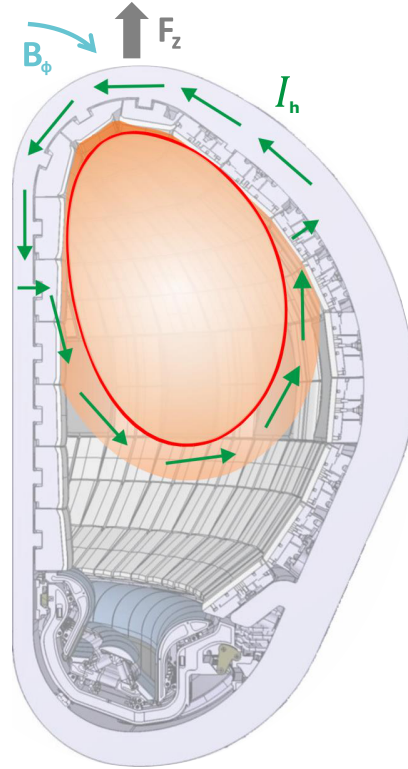


Fig. 2.2 Schematic representation of the electromechanical loads during an upward vertical displacement event (VDE). The halo current (I_h , green arrows) flows along the open magnetic field lines in the scrape-off layer, intercepts the first wall, and closes its circuit poloidally through the conductive vacuum vessel. The interaction between the poloidal component of this current and the toroidal magnetic field (B_ϕ , light blue arrow) generates a net vertical Lorentz force (F_z , gray arrow) that acts on the machine structure. Reproduced from [28]. © ITER Organization.

LCFS touches the PFCs, the resulting scraping of the plasma forces a fraction of the plasma current to flow along the open field lines (encompassing the so-called *halo* region [15]), enter the wall, flow through the conductive structure, and then re-enter the plasma [7]. This current is known as *halo current* (I_h) and is depicted in Figure 2.2 (green arrows). Additionally, these currents can be induced by the change in magnetic flux during the rapid current decay in the CQ phase and by the vertical motion of the plasma column [29].

The magnitude and spatial structure of halo currents are typically quantified by two parameters:

- **Halo fraction (HF):** It is defined as the maximum axisymmetric poloidal halo current normalized to the pre-disruption plasma current [30]:

$$\text{HF} = \frac{I_{h,\max}}{I_{p,0}} \quad (2.19)$$

In general, HF is found to be higher in hot VDEs compared to cold VDEs [15, 31]. Values up to $\approx 60\%$ have been observed in JET discharges [30].

- **Toroidal peaking factor (TPF):** Since halo currents are rarely axisymmetric, this parameter describes their toroidal asymmetry:

$$\text{TPF} = \frac{\max_{\phi}(J_h)}{\langle J_h \rangle_{\phi}} \bigg|_{I_{h,\max}} \quad (2.20)$$

where $\max_{\phi}(J_h)$ is the maximum value of the halo current density J_h at a given toroidal position and $\langle \cdot \rangle_{\phi}$ represents the toroidal average [30]. The TPF is usually calculated at the time of the maximum current ($I_{h,\max}$).

The product of these two parameters ($\text{HF} \cdot \text{TPF}$) is frequently used to estimate the maximum local poloidal halo current (and thus the worst-case local mechanical stress on machine components) as a function of the pre-disruption current $I_{p,0}$.

2.2.3 Disruption mitigation systems

Given the potential for irreversible machine damage, which would severely compromise sustained energy production, disruption avoidance is a paramount operational priority. However, when a disruption becomes inevitable (e.g., a hardware failure has occurred) a fast-acting disruption mitigation system (DMS) must be triggered to minimize the consequences. The fundamental strategy of modern DMS is to preemptively inject large quantities of high-Z impurities (such as neon or argon) or deuterium into the plasma before the natural TQ occurs, in order to trigger it. The presence of these impurities has the main purpose of uniformly radiating the stored thermal energy (W_{th}) over the entire first wall (FW) surface, preventing the localized deposition of heat onto specific portions of the PFCs. Additionally, the impurities cause a faster CQ, typically reducing the halo currents by more than 50% and reducing the TPF to ≈ 1 in modern devices such as ASDEX Upgrade, JET and DIII-D [15, 32, 33]. Recent non-linear MHD simulations and experimental analyses have clarified

the physics behind this EM load reduction [34, 35]: relatively strong toroidal currents in the SOL region effectively clamp the plasma vertical current centroid Z_{curr} near the equatorial plane, preventing the large vertical displacements required to generate large electromechanical forces ($F_z \propto I_p \Delta Z_{\text{curr}}$). Indeed, a significantly reduced ΔZ_{curr} is consistently observed in mitigated disruptions across various devices [36, 37, 38, 34]. Notwithstanding these toroidal halo currents, the poloidal component remains low in mitigated disruptions, reducing the resulting F_z on the structures [34].

Two primary techniques have been developed for injecting material inside the plasma environment:

- **Massive gas injection (MGI):** This method involves the rapid injection of gas through fast valves located at the vessel wall. Experimental observations indicate that the neutral gas itself does not penetrate deeply into the high-pressure core; instead, the rapid cooling at the edge triggers large-scale MHD instabilities [39], responsible for the rapid mixing of the impurities into the plasma center, facilitating the fast radiation of the thermal energy [15]. While MGI has been successfully employed in present-day devices [40, 41, 42], the effectiveness of its penetration and assimilation in reactor-scale plasmas is potentially limited [15].
- **Shattered pellet injection (SPI):** To overcome the penetration limitations of MGI, the shattered pellet injection (SPI) concept [43, 44] was developed and tested [43, 45, 46] and has been selected as the baseline mitigation strategy for ITER [45]. In an SPI system, a cryogenic pellet is accelerated to high speeds (typically hundreds of m/s) and shattered against a bent tube (the “shatter head”) just before entering the vacuum vessel. The resulting spray of solid fragments penetrates deep into the plasma core while ablating, promoting a more uniform assimilation of the material compared to surface gas injection [47, 44]. However, since the radiative cooling rate of impurities (such as neon) depends strongly on the electron temperature (T_e) and can be significantly lower in the core than at the edge during the pre-TQ phase, this more uniform assimilation does not necessarily guarantee a more efficient radiated power fraction.

While impurity injection is effective at mitigating thermal loads, it introduces a critical trade-off regarding REs. The rapid cooling induced by the impurities

dramatically increases the plasma resistivity and, consequently, the induced parallel electric field ($E_{\parallel}$). If this field exceeds a critical threshold, it can trigger the generation of REs [48, 49, 15]. Conversely, a sufficiently high density of free electrons (the so-called “Rosenbluth density” [15, 48]) can suppress the RE avalanche mechanism through collisional damping. This density threshold for ITER-scale tokamaks is typically orders of magnitude higher than the pre-disruption electron density [15, 50] and attempting to achieve it would impose considerable stress on the technological systems of the tokamak [50]. Furthermore, because of the strong RE avalanche gain of ITER (see Section 2.3.3), any RE seed would be efficiently multiplied into a multi-MA RE beam [51, 52]. Consequently, rather than aiming for the Rosenbluth density, the goal of the ITER DMS is to completely avoid RE generation before the avalanching phase. This relies on an appropriate tuning of the material injection: it must be sufficient to mitigate heat loads by triggering a TQ, but the composition and timing must be carefully orchestrated to prevent RE generation altogether. To achieve this, a potential strategy is to inject deuterium pellets prior to the high-Z impurity injection: 3D non-linear MHD simulations have shown that this increases the collisional drag and electron density before the TQ is triggered [53]. Because the primary objective of this pure deuterium injection is to maximize electron density for RE avoidance, deep core assimilation is crucial. However, achieving this deep assimilation remains physically challenging; recent 3D non-linear MHD simulations of JET discharges have demonstrated that strong outward $\mathbf{E} \times \mathbf{B}$ plasmoid drifts during deuterium SPI can substantially limit effective core assimilation [54].

If prevention fails and a macroscopic RE beam is formed, the strategy must shift from avoidance to mitigation, with the ultimate goal to dissipate the energy of the RE beam and ensure its benign termination. A primary method to reach this benign termination scenario relies on injecting a low-Z pellet to deliberately induce harmless instabilities that safely dissipate the RE beam [55]. During this process, the injected low-Z material (typically deuterium) triggers large-scale current-driven MHD instabilities, whose rapid growth stochastizes the magnetic flux surfaces, leading to a deconfinement of the REs [18, 56]. Recent simulations have shown that this magnetic stochastization broadens the wetted area of the RE beam on the PFCs, avoiding localized hotspots and mitigating the risk of melting [57, 58]. Crucially, because REs are rapidly expelled, the current is transferred back to the cold bulk plasma, ensuring that the remaining stored magnetic energy is harmlessly dissipated as radiation during the CQ [56].

Given the severe risk these relativistic particles pose to machine integrity and the complexity of the mechanisms governing their generation, acceleration, and eventual loss, the fundamental physics of REs is analyzed in detail in Section 2.3.

2.2.4 The DTT disruption mitigation system

At the DTT facility, a dedicated DMS will be implemented with a dual objective: ensuring robust machine protection against the severe loads of unmitigated disruptions and REs, and serving as a crucial experimental testbed. Indeed, one of the main goals of DTT is to support ITER and EU-DEMO regarding disruption and RE mitigation, playing a pivotal role in validating and optimizing mitigation strategies for reactor-scale plasmas [10].

From a technological standpoint, the DTT-DMS is planned to be implemented in phases, aligning with the operational phases of the device. While an MGI system is expected to be employed during the Day-0 operations, the ultimate objective is to commission an SPI as the primary mitigation technology for the Full Power scenario of DTT [9]. Because the detailed engineering and physical design of the DTT-DMS is currently ongoing, predictive modeling efforts are heavily relied upon. The RE avalanche assessments presented in this thesis serve as a foundational step in this effort. By mapping the parameter space for RE generation, this work will assist in designing a DMS capable of mitigating thermal loads while minimizing the risk of macroscopic RE beam generation.

2.3 Runaway electrons

While the MHD framework described in the previous sections treats the plasma as a thermal fluid, the phenomenon of REs is inherently kinetic. It arises from a fundamental property of Coulomb collisions in a fully ionized plasma, first predicted by Wilson [59, 60]: the collisional friction force acting on an electron decreases as its velocity increases. Specifically, in a plasma the friction force for suprathermal electrons scales as $F_{\text{drag}} \propto v^{-2}$. Consequently, since the driving electric force is independent of velocity, any electron exceeding a critical velocity will experience a net acceleration that grows with speed, causing it to effectively “run away” from the thermal distribution.

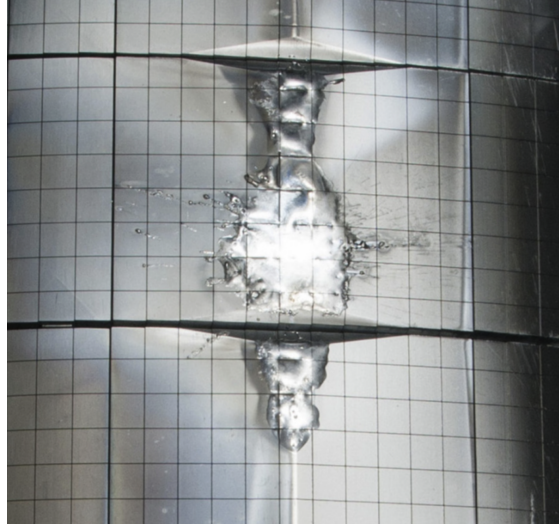


Fig. 2.3 Visual evidence of localized damage resulting from the impact of a collimated RE beam on the beryllium limiter of the JET tokamak. The concentrated energy deposition led to significant surface melting of the tiles. Reproduced from [61].

In the context of the disruptions described in Section 2.2, the conditions for this phenomenon are highly favorable. Following the TQ, the plasma becomes cold and highly resistive, causing the plasma current to decay (CQ). To conserve magnetic flux, this decay induces a very large toroidal electric field ($E_{\parallel}$), often substantially above the critical threshold required to continuously accelerate electrons out of the thermal distribution. Under these conditions, the magnetic energy stored in the tokamak is converted into the kinetic energy of an RE beam. The formation of a macroscopic RE beam poses a critical threat to the integrity of high current tokamaks, since the avalanche mechanism can amplify a small seed population by a factor of around 10^{20} in reactor-scale devices like ITER [48]. Unlike the thermal plasma, which disperses energy over a relatively wide area, these RE beams are highly collimated and can deposit megajoules of energy onto localized regions of the PFCs [15, 61]. Figure 2.3 provides experimental evidence of this damage mechanism, showing the melting of beryllium tiles in JET caused by the impact of an RE beam [61]. While this specific event was limited to surface melting, the significantly higher stored energies in reactor-scale devices would likely result in deep structural failure and possible loss of coolant accidents.

The following subsections detail the theoretical framework governing REs. Section 2.3.1 introduces the dynamics of individual particles and the critical field thresholds. Sections 2.3.2 and 2.3.3 describe the mechanisms by which these

electrons are generated, distinguishing between primary seed formation and secondary avalanche multiplication. Finally, Section 2.3.4 discusses the transport and eventual loss of these relativistic particles.

2.3.1 Runaway dynamics

To properly describe the evolution of the RE population during a disruption, it is necessary to move beyond the fluid description and adopt a kinetic approach. The behavior of electrons is governed by the relativistic version of Equation (2.1) for $\alpha = e$, which reads:

$$\frac{\partial f_e}{\partial t} + \mathbf{v} \cdot \nabla f_e - e (\mathbf{E} + \mathbf{v} \times \mathbf{B}) \cdot \nabla_p f_e = C_e + S_e \quad (2.21)$$

where S_e has been added to take into account secondary generation mechanisms (e.g., knock-on collisions), and $\frac{1}{m_e} \nabla_v$ has been substituted by ∇_p (p being the momentum) to correctly treat relativistic effects.

While the full kinetic evolution requires solving this equation, the fundamental condition for electron runaway is best understood by analyzing the parallel force balance. Following Dreicer [62, 63], who analyzed a system in which electrons and ions exist in thermal equilibrium within a homogeneous plasma, the simplified force balance reads:

$$\frac{dp_{\parallel}}{dt} - eE_{\parallel} = -F_{\text{drag},\parallel} = -p_{\text{th}} \nu_{ee} \Psi(p) \quad (2.22)$$

where $F_{\text{drag},\parallel}$ is the parallel component of the friction force, $p_{\text{th}} = \sqrt{2m_e T_e}$ is the thermal momentum (T_e being the electron temperature), ν_{ee} is the thermal electron-electron collision frequency, and $\Psi(p)$ represents the variation of collisional friction with respect to energy. The collision frequency is given by:

$$\nu_{ee} = \frac{m_e n_e e^4 \ln \Lambda}{4\pi \epsilon_0^2 p_{\text{th}}^3} \quad (2.23)$$

where m_e is the electron rest mass, n_e is the electron density, and $\ln \Lambda$ is the Coulomb logarithm. The function $\Psi(p)$ monotonically increases with velocity in the thermal range ($v \ll v_{\text{th}}$) before peaking at $\Psi(p_{\text{th}}) \approx 0.428$. Then, it decreases monotonically as $\approx p_{\text{th}}^2 / (mv)^2$ for suprathermal electrons [64]. This scaling ($F \propto v^{-2}$) at high velocities creates the possibility for “unlimited” acceleration if the parallel electric

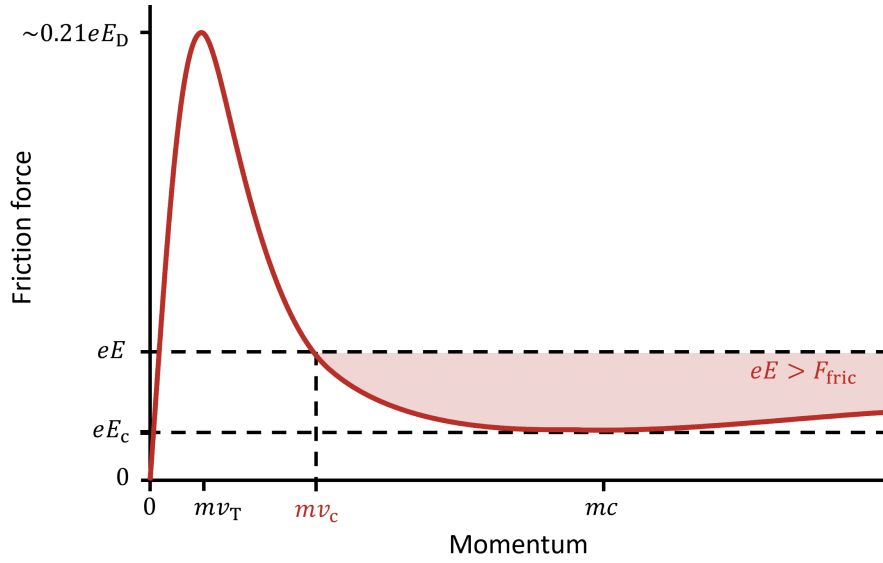


Fig. 2.4 Graphical representation of the force balance governing electron runaway. The red curve depicts the collisional friction force, which for suprathermal electrons decreases with increasing momentum. The dashed horizontal lines represent the driving force from the electric field (eE) and the critical electric field force (eE_c). A critical momentum mv_c exists at the intersection of the electric field force eE and the collisional drag force: for higher momenta (shaded red region), electrons experience a net positive acceleration and enter the runaway regime. The minimum field required to enable this process is the critical electric field E_c ; for fields below this threshold, no runaway generation occurs. Conversely, if E is high enough to exceed the slide-away value of $\sim 0.21eE_D$ (E_D being the Dreicer electric field), the whole electron population is converted into runaways. Reproduced from [65].

force $eE_{\parallel}$ exceeds the friction drag $F_{\text{drag},\parallel}$, as illustrated in Figure 2.4. For electrons whose momentum is much larger than the thermal one ($p \gg p_{\text{th}}$), the magnitude of the friction force can be written as:

$$F_{\text{drag},\parallel}(p \gg p_{\text{th}}) \approx \frac{n_e e^4 \ln \Lambda}{4\pi \epsilon_0^2 m_e v^2} \quad (2.24)$$

By defining the *Dreicer electric field* as:

$$E_D = \frac{2p_{\text{th}} v_{\text{ee}}}{e} = \frac{n_e e^3 \ln \Lambda}{4\pi \epsilon_0^2 T_e} \quad (2.25)$$

it is possible to write the maximum friction force, which occurs near $p = p_{\text{th}}$ (where $\Psi(p_{\text{th}}) \approx 0.428$), as:

$$\max(F_{\text{drag},\parallel}) \approx 0.428 p_{\text{th}} v_{\text{ee}} = 0.214 e E_D \quad (2.26)$$

If the applied force ($eE_{\parallel}$) is greater than this value, the entire electron population would theoretically run away.

Considering that for suprathermal electrons $F_{\text{drag},\parallel}$ decays as v^{-2} and, thus, has a minimum as v approaches the speed of light c , it is possible to find a threshold value for which there is no possibility of runaway. The electric field required to overcome this minimum drag and sustain the runaway of an already relativistic electron (i.e., with $v \rightarrow c$) is the *Connor-Hastie critical electric field* E_c [66]:

$$E_c = \frac{n_e e^3 \ln \Lambda}{4\pi \epsilon_0^2 m_e c^2} = \frac{T_e}{m_e c^2} E_D \quad (2.27)$$

For values of the electric field $E_{\parallel}$ between these two thresholds ($E_c < E_{\parallel} < 0.214 E_D$), a fraction of the electron population will run away and become relativistic (shaded red area in Figure 2.4). For a given $E_{\parallel}$, this occurs for velocities exceeding:

$$v > c \sqrt{\frac{E_c}{E_{\parallel}}} \equiv v_c \quad (2.28)$$

where v_c is the critical speed for runaway (the smallest velocity at which electrons can become REs). This can also be expressed in terms of the critical momentum using the relativistic definition $p = \gamma m_e v$, where $\gamma = (1 - \frac{v^2}{c^2})^{-1/2}$ is the Lorentz factor:

$$p_c = \frac{m_e c}{\sqrt{\frac{E_{\parallel}}{E_c} - 1}} \quad (2.29)$$

While Equation (2.22) establishes the fundamental threshold for RE generation, experimental measurements [67, 68] indicate that the “effective” critical field $E_{c,\text{eff}}$ necessary to maintain an RE beam consistently exceeds the Connor-Hastie limit, sometimes by an order of magnitude. This divergence stems from the simplified nature of the theoretical model, which excludes non-collisional energy losses and the perpendicular dynamics. In a real scenario, significant braking mechanisms such

as synchrotron emission [69], bremsstrahlung, and collisions with partially ionized high-Z impurities [70, 71, 72] increase the total drag. These additional energy sinks ultimately limit the maximum energy of the RE beam and they are detailed in Section 2.3.4.

2.3.2 Primary generation

Primary generation encompasses the physical processes capable of accelerating electrons directly from the thermal bulk distribution across the critical momentum threshold p_c (Equation (2.29)) into the runaway region. A defining characteristic of these mechanisms is that they do not require a pre-existing population of REs, unlike secondary generation. They provide, instead, the essential initial RE “seed” population. While the number of primary runaways may be relatively small compared to the final beam intensity in large devices, their presence is indispensable for triggering the subsequent exponential amplification via the secondary avalanche mechanism (described in Section 2.3.3). The principal primary mechanisms relevant to tokamak disruptions are the *Dreicer source*, *hot-tail*, acceleration driven by magnetic reconnection, and sources from nuclear processes such as beta decay and Compton scattering.

- **Dreicer source:** This classical mechanism for runaway formation, first described by Dreicer in 1959 [62], is a diffusive process occurring in velocity space. Since electrons with a velocity exceeding v_c are rapidly accelerated away, they leave a depleted region in the tail of the Maxwellian distribution. Collisional velocity diffusion in the bulk constantly attempts to refill this high-energy tail as particles become REs, resulting in a continuous, diffusive “leak” of electrons across the critical velocity boundary into the runaway region [73]. A simplified estimate for this generation rate can be expressed as

$$\frac{dn_{\text{RE}}}{dt} \sim \frac{eE}{\sqrt{m_e T_e}} n_e \exp\left(-\frac{E_D}{2E}\right)$$

under the assumption of a rigid Maxwellian distribution and electrons becoming REs within one collision time [73]. While the Maxwellian distribution is strictly not valid for suprathermal electrons, this expression effectively illustrates its fundamental scaling: it becomes exponentially weak in the regime where

the electric field is small compared to the Dreicer electric field ($E \ll E_D$). The normalized electric field typically required to produce a non-negligible RE population within relevant experimental timescales is $E/E_D \sim 0.02$ [74]. Consequently, while in many modern tokamaks Dreicer generation remains competitive with other primary mechanisms such as hot-tail, it is expected to be negligible in ITER post-disruption conditions [75, 72, 73].

- **Hot-tail:** In the context of mitigated tokamak disruptions, the hot-tail mechanism is often the dominant primary source [76, 77, 78, 52, 79]. This phenomenon arises from the extreme disparity between the collisional relaxation timescale of cold and suprathermal electrons. While thermal bulk electrons collide frequently and thermalize rapidly to the new cold temperature during a TQ, suprathermal electrons have much lower collision frequencies and therefore lose their energy much more slowly. Consequently, the high-energy (“hot”) tail of the Maxwellian distribution retains most of its high initial energy on the timescale of bulk cooling. As the bulk plasma cools, its resistivity increases rapidly, inducing a sharp rise in the toroidal electric field $E_{\parallel}$. If this electric field rises fast enough, the critical velocity v_c may drop below the velocity of these surviving high-energy electrons. These particles are then effectively accelerated by the electric field before they can thermalize, forming an RE seed [73]. This process creates a significant distortion in the electron distribution function, as illustrated schematically in Figure 2.5.
- **Magnetic reconnection:** Unlike the previous mechanisms driven by the global axisymmetric loop voltage, this source stems from the rapid conversion of magnetic energy into particle kinetic energy during local topological rearrangements. The physics of electron acceleration in reconnecting plasmas generally involves three distinct pathways [73]: (i) direct acceleration by the strong DC electric field generated at the X-points of magnetic islands; (ii) stochastic acceleration of resonant electrons by linear waves; and (iii) acceleration at shock fronts. While these processes are the primary engine for high-energy particles in astrophysical phenomena (such as solar flares [80]), the strong toroidal guide field in tokamaks constrains the dynamics. Consequently, a unified theoretical framework for this source in tokamaks is still under development, and experimental evidence specifically during the TQ remains scarce [73].

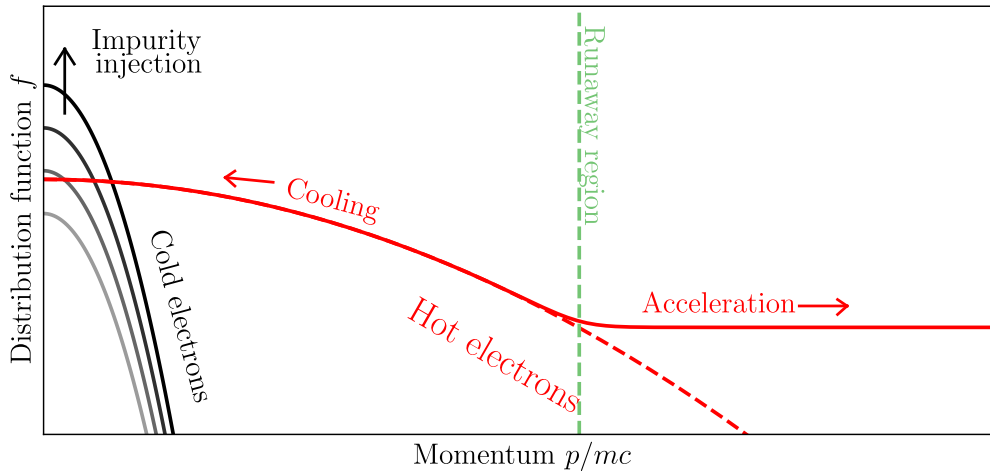


Fig. 2.5 Schematic representation of the electron distribution function in the early phases of a mitigated disruption. The injected impurities provide the plasma with cold electrons through ionization (gray and black curves) which slow down the hot electrons (red curve) through collisions. However, the higher-energy tail cools much more slowly and, if the induced electric field is strong enough to define a critical momentum (vertical dashed green line) that falls within this surviving hot tail, these particles are accelerated into the runaway regime. Reproduced from [16].

Furthermore, predictions are complicated by uncertainties regarding the degree of stochastization of flux surfaces during a disruption [81].

- **Nuclear sources:** In fusion reactor devices operating with DT fuel, such as ITER or EU-DEMO, nuclear processes provide an intrinsic source of high-energy electrons:
 - *Tritium beta decay:* Tritium is radioactive with a half-life of ≈ 12.3 years. Its β decay produces electrons with a continuous energy spectrum up to a maximum of 18.6 keV, providing a continuous volumetric source of suprathermal particles, which may become REs depending on the value of v_c .
 - *Compton scattering:* In a nuclear environment, γ rays (emitted by neutron-activated metallic walls) can undergo Compton scattering with bulk plasma thermal electrons, transferring sufficient momentum to knock them directly into the MeV energy range, instantly converting them into REs.

These nuclear sources ensure that, in future fusion power plants, the avalanche mechanism (Section 2.3.3) always has an RE seed to amplify regardless of the TQ dynamics.

2.3.3 Secondary generation (RE avalanche)

While primary generation mechanisms provide the initial seed population of REs, it is the secondary generation, or “avalanche” mechanism, that poses the greatest threat to high-current devices. In these machines, the avalanche mechanism acts as a massive amplifier of the initial RE seed, ultimately dominating the runaway dynamics and driving the formation of macroscopic RE beams.

The avalanche process is driven by large-angle, close Coulomb collisions between an existing high-energy RE and a thermal electron from the bulk plasma. This is distinct from the cumulative small-angle scattering that dominates the frictional drag. In such a large-angle collision, a primary RE transfers a significant fraction of its momentum to a thermal electron while still retaining sufficient energy to stay in the runaway region. If the energy transferred is sufficient to kick the secondary electron above the critical momentum p_c (Equation (2.29)), this secondary electron becomes a new RE. The probability of these events is governed by the differential Møller cross-section [82, 83] for the scattering between a relativistic electron and a thermal one (assumed stationary [84]):

$$\frac{d\sigma}{d\varepsilon} = \frac{2\pi r_e^2}{\gamma_0^2 - 1} \left[\left(\frac{\gamma_0}{\gamma - 1} \right)^2 + \left(\frac{\gamma_0}{\gamma_0 - \gamma} \right)^2 + 1 - \frac{2\gamma_0 - 1}{\gamma_0 - 1} \left(\frac{1}{\gamma - 1} + \frac{1}{\gamma_0 - \gamma} \right) \right] \quad (2.30)$$

where r_e is the classical electron radius and γ_0, γ are the Lorentz factors of the incoming RE and secondary electron, respectively [73]. The two leading terms correspond to the classical Rutherford cross-section, whereas the remaining components arise from the exchange interaction [73]. The total Møller cross-section can be found by integrating Equation (2.30) for all possible momenta, introducing a minimum $\gamma_{\min}$ (i.e. neglecting low-angle collisions below a certain threshold) to ensure that a collision actually takes place and that the integral does not diverge [85]:

$$\sigma = \int_{\gamma_{\min}}^{\gamma_0 + 1 - \gamma_{\min}} \frac{d\sigma}{d\gamma} d\gamma \quad (2.31)$$

Since the production rate of new REs is proportional to the instantaneous density of existing runaways, this mechanism leads to an exponential amplification of the RE current. The avalanche growth rate can then be written as [86]:

$$\Gamma \equiv \frac{1}{n_{\text{RE}}} \frac{dn_{\text{RE}}}{dt} = \frac{1}{2} n_e \sigma \frac{\sqrt{\gamma_0^2 - 1}}{\gamma_0} c = \frac{1}{2} n_e \sigma v_0 \quad (2.32)$$

where v_0 is the velocity of the incoming RE. The factor $1/2$ is introduced to avoid double-counting the electrons in phase-space integration, since they are indistinguishable [86]. It is important to note that the simplified source term in Equation (2.32) represents a specific limit, as this formulation formally assumes that the primary electrons are ultrarelativistic ($\gamma_0 \rightarrow \infty$) and have a negligibly small pitch angle. While this simplification is justified when the electric field significantly exceeds the runaway avalanche threshold, a truly general treatment requires solving the exact knock-on collision source alongside a proper sink term in the continuous kinetic equation to preserve particle conservation [73, 85].

Neglecting radiative limitations and considering a uniform and fully ionized plasma, Rosenbluth and Putvinski [48] calculated the avalanche growth rate to be

$$\Gamma \approx \frac{e}{m_e c \ln \Lambda_c} \frac{E_{\parallel} - E_c}{\sqrt{5 + Z_{\text{eff}}}} \quad (2.33)$$

where Z_{eff} is the effective charge of the plasma and $\ln \Lambda_c$ is the relativistic Coulomb logarithm, relevant when considering collisions between runaways and thermal electrons [87]. The relativistic Coulomb logarithm can be linked to the standard thermal electron-electron Coulomb logarithm $\ln \Lambda_0 \approx 14.9 + \ln(T_{\text{keV}}/n_{e20}^{0.5})$ [1] through this expression [87]:

$$\ln \Lambda_c = \ln \Lambda_0 + \frac{1}{2} \ln \left(\frac{m_e c^2}{T_e} \right) \approx 14.6 + 0.5 \ln \left(\frac{T_{\text{eV}}}{n_{e20}} \right) \quad (2.34)$$

where T_{keV} and T_{eV} are the temperatures normalized to 1 keV and 1 eV, respectively, and n_{e20} is the electron density normalized to 10^{20} m^{-3} . While the standard thermal Coulomb logarithm evaluates to $\ln \Lambda_0 \approx 8$ in typical cold, dense post-disruption plasmas ($T_e \sim 10 \text{ eV}$, $n_e \sim 10^{22} \text{ m}^{-3}$), the relativistic correction adds $0.5 \ln(m_e c^2/T_e) \approx 5.5$ to this baseline, yielding $\ln \Lambda_c \approx 13.5$. Conversely, under much hotter pre-disruption fusion conditions ($T_e \sim 1 - 25 \text{ keV}$, $n_e \sim (0.5 - 1) \cdot 10^{20} \text{ m}^{-3}$), the

relativistic correction diminishes, reducing the discrepancy between the two Coulomb logarithms, which typically range between 14 and 20.

Equation (2.33) highlights that Γ is driven linearly by the excess electric field ($E_{\parallel} - E_c$) and suppressed by the effective charge Z_{eff} . For a very strong electric field in the limit $E_{\parallel} \gg E_c$ (commonly encountered during a disruption) and considering a hydrogen plasma ($Z_{\text{eff}} = 1$) [74, 48], it is possible to write the avalanche growth rate as:

$$\Gamma \approx \frac{eE_{\parallel}}{2m_e c \ln \Lambda_c} \quad (2.35)$$

While Equations (2.33) and (2.35) are foundational, they rely on the assumptions of a uniform, fully ionized plasma. This is rarely the case during disruption mitigation via massive material injection, where the resulting partially ionized conditions can significantly enhance the avalanche growth rate. This occurs because the increased friction force on fast electrons only partially compensates for the additional bound target electrons that become available for the avalanche process via large-angle collisions [88].

In the fully ionized limit, the baseline avalanche gain, representing the ratio of final to initial REs for an infinitesimal RE seed population, can be obtained by integrating Equation (2.35) over the duration of the CQ, considering that $E_{\parallel}$ is driven by the current decay:

$$G_{\text{av}} = \frac{n_{\text{RE}}^{\text{final}}}{n_{\text{RE}}^{\text{seed}}} \approx \exp\left(\frac{I_p}{I_A \ln \Lambda_c}\right) \quad (2.36)$$

where $I_A \approx m_e c^3 / e \approx 0.02$ MA is the Alfvén current [48]. It is crucial to note that despite this massive exponential amplification, the final RE beam current is not always strictly proportional to the initial seed. In fact, in high-current scenarios where the avalanche gain is massive (such as the DT phase of ITER), the final runaway current has been shown to be only logarithmically sensitive to the magnitude of the seed [89]. This relatively weak dependence arises from the self-consistent dynamics of the electric field: when the rapidly growing RE current becomes comparable to the decaying Ohmic current, the induced electric field $E_{\parallel}$ is reduced, which in turn suppresses the avalanche growth rate [89]. Consequently, the final RE current in reactor-scale devices is largely decoupled from the precise magnitude of the initial RE seed. This presents a formidable challenge for disruption mitigation strategies,

since achieving a meaningful reduction in the final RE beam requires reducing the initial seed by several orders of magnitude.

Because the primary dependence of the avalanche gain G_{av} is on the pre-disruption plasma current $I_{p,0}$, there is a fundamental difference in how these dynamics unfold between present-day low-current machines and future high-current reactors [71, 72, 89, 52]. Employing the widely used historical value of $\ln \Lambda_c \approx 20$ [48, 15], which represents an optimistic scenario for a post-disruption plasma (since smaller values would lead to a larger avalanche gain), yields the simplified scaling $G_{av} = \exp(2.5 \cdot I_p)$ (I_p in MA) [15]. This allows us to distinguish the following scenarios:

- In small-size tokamaks like ASDEX Upgrade ($I_p = 1.4$ MA), the gain is negligible ($G_{av} \sim 30$), meaning RE beams are rarely observed unless deliberately triggered. The Day-0 scenario of DTT ($I_p = 2$ MA, $G_{av} \sim 150$) falls into this category.
- For medium-size tokamaks such as JET or DTT, the gain can increase by orders of magnitude, depending on the specific scenario. For instance, considering the full power scenario of DTT ($I_p = 5.5$ MA), the theoretical avalanche gain can reach $G_{av} \sim 1 \cdot 10^6$.
- In future reactor-scale devices, the gain becomes massive. Given the 15 MA plasma current of ITER, the gain exceeds 10^{16} , which implies that the conversion of the majority of the current into RE current is virtually unavoidable. While this is a theoretical value, its magnitude is nevertheless of great concern, especially considering future power plants such as EU-DEMO which will have a plasma current up to 19 MA ($G_{av} \sim 4 \cdot 10^{20}$).

This substantial difference in amplification factors positions DTT in a unique operational space. While the Day-0 scenario remains in a regime of relatively low avalanche multiplication, the full power phase introduces a shift where the avalanche mechanism becomes the dominant driver of the final RE beam intensity. This thesis will specifically investigate this transition, analyzing how the jump from the low- G_{av} initial phase to the high- G_{av} full-current scenario impacts the disruption dynamics and mitigation strategies.

2.3.4 Transport and losses

The final intensity of an RE beam depends not only on the generation rates described in the previous sections, but also on the confinement properties of the relativistic electrons. As discussed in Section 2.2, the rapid growth of MHD instabilities (such as tearing modes) during the TQ destroys the nested magnetic flux surface topology. This leads to a region of magnetic stochasticity where field lines “diffuse” radially across the minor radius. Since REs possess high parallel velocities ($v_{\parallel} \approx c$), they quickly stream along these chaotic field lines, leading to rapid radial diffusion. This transport is often described by the diffusion coefficient $D_{\parallel}$ obtained by Rechester and Rosenbluth [90]:

$$D_{\parallel} \approx \pi q R_0 \left(\frac{\delta B}{B} \right)^2 v_{\parallel} \quad (2.37)$$

where q is the safety factor, R_0 the major radius, and $\delta B/B$ the relative amplitude of the magnetic perturbation. This mechanism is particularly effective at deconfinement during the TQ, potentially suppressing the formation of the initial hot-tail seed if the stochastic region extends to the PFCs. However, non-linear MHD simulations, such as those performed with the JOEKE code [57], reveal that magnetic surfaces often reform as the MHD activity saturates during the CQ. This allows some of the RE population to be trapped again in the core and act as a new RE seed that can be avalanched, as illustrated in Figure 2.6. In addition to magnetic stochasticity, the particle orbits themselves play a critical role in confinement. As an RE accelerates, its orbit increasingly deviates from the magnetic flux surface. For multi-MeV electrons, this outward shift can be macroscopic, eventually causing the trajectory to intersect a solid component [91]. This results in REs being scraped off and losing all their energy even if the magnetic axis remains centered.

Furthermore, macroscopic vertical motion of the plasma during a disruption, such as in VDEs (discussed in Section 2.2.2), leads to a similar and highly consequential scraping-off effect. As the plasma cross-section is pushed onto the PFCs, toroidal flux surfaces are sequentially opened. Since REs stream rapidly along these open field lines, they are quickly lost to the wall and can no longer proliferate via the avalanche mechanism. Recent 2D non-linear MHD simulations with JOEKE [92] have demonstrated that this vertical displacement drastically reduces the poloidal flux available for RE acceleration, limiting the theoretical maximum avalanche gain in ITER to a range of 10^9 – 10^{12} (compared to earlier stationary-plasma estimates of

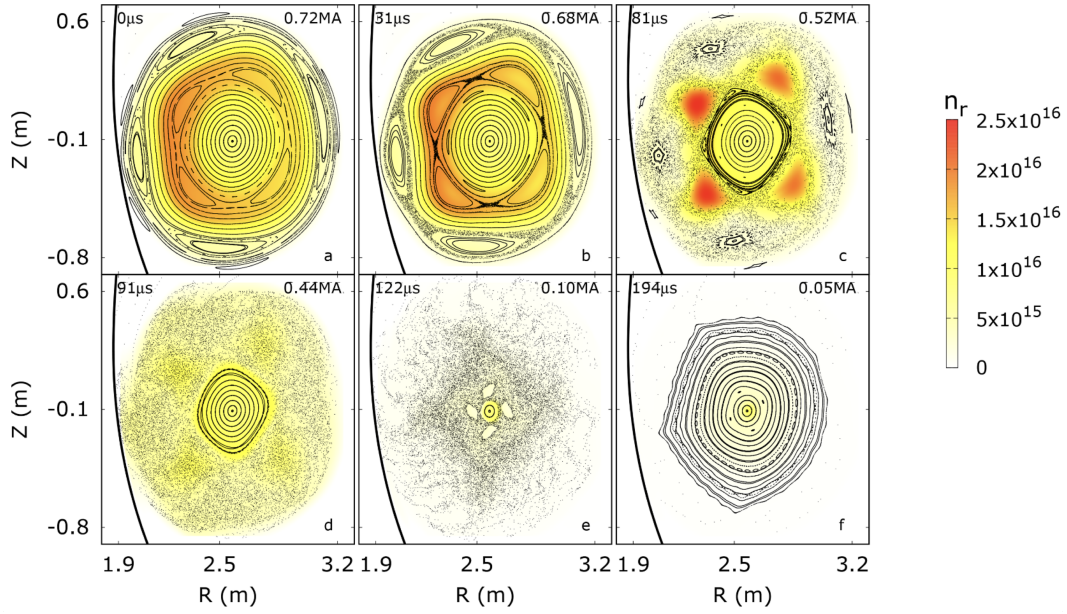


Fig. 2.6 Poincaré plots from a JOREK simulation showing the evolution of the RE density (n_r) and magnetic field stochasticity during a disruption. Starting from intact surfaces (a), the topology becomes highly stochastic as the disruption progresses (b-d), leading to the loss of REs. As the CQ proceeds, flux surfaces begin to re-form (e-f), effectively trapping the remaining REs in the core. Reproduced from [57].

10^{16} – 10^{20}). This dynamic has also been captured in computationally efficient 1D models: extensive studies utilizing the DREAM code demonstrate that a sufficiently slow RE avalanche allows the sequential destruction of flux surfaces to prevent the formation of a macroscopic RE plateau entirely, making possible the total RE avoidance even under the stringent conditions of a 15 MA ITER discharge [93].

Finally, the RE population is regulated by energy dissipation mechanisms that counteract the electric field acceleration. Two radiative processes are dominant:

- **Synchrotron radiation:** Unlike the cyclotron emission from thermal electrons which is largely reabsorbed by the plasma, synchrotron radiation emitted by relativistic electrons escapes the plasma, acting as a net energy sink [94]. The associated radiation reaction force scales strongly with the perpendicular momentum ($p_{\perp}$). Consequently, electrons with large pitch angles lose energy most rapidly. Although the runaway population is already strongly collimated by the parallel momentum gained from the electric field, this radiation can further reduce the average pitch angle, reinforcing the narrow, beam-like distribution.

- **Collisions with heavy ions:** Following mitigated disruption scenarios, the plasma often contains high densities of partially ionized impurities (e.g., argon, neon) or wall materials (like tungsten). REs can penetrate the electron clouds of these heavy ions, scattering off a nuclear charge effectively larger than the net ion charge. This enhances elastic scattering, which increases the collisional drag force, while simultaneously driving significant radiative energy dissipation through Bremsstrahlung [95, 89, 71]. While a relativistic electron can lose energy ionizing atoms, the released electron may contribute to the avalanche effect if it is released with a velocity larger than v_c .

The interplay between the generation, transport, and dissipation mechanisms discussed in this Section defines if a macroscopic RE beam forms and dictates whether it results in a benign dissipation of energy or a catastrophic impact against the PFCs.

Chapter 3

Numerical tools and methodology

The theoretical framework established in Chapter 2 highlights the inherent multi-scale and multi-physics nature of tokamak disruptions. The complex interplay between macroscopic MHD instabilities, magnetic field stochasticity, and disruption dynamics cannot be fully captured by analytical models or reduced 0D/1D codes. To transition from a qualitative understanding to a predictive description of these phenomena, self-consistent non-linear simulations are essential.

In this thesis, the non-linear extended MHD code JOREK [96, 97, 98] is employed as the primary numerical tool. Originally developed to investigate instabilities such as edge-localized modes and kink modes in X-point geometry [99], JOREK has evolved into a modular framework capable of simulating diverse large-scale plasma dynamics, including disruptions and VDEs. Its capability to handle realistic tokamak geometries and its use of fully implicit time-stepping schemes make it well-suited for the self-consistent modeling required in this work. The simulations presented here rely on a specific combination of physical models: the reduced MHD fluid description for the thermal plasma, extended by a fluid description for the REs, and a resistive wall model for electromagnetic (EM) boundary conditions. The reduced MHD model is particularly advantageous, as it retains the essential physics while allowing for significantly faster computation. This efficiency arises from the elimination of fast magnetosonic waves, which enables the use of larger time steps, as well as from the reduced number of variables [96].

This chapter details the numerical implementation and methodology used to obtain the results discussed in Chapters 4 and 5. Section 3.1 introduces the JOREK

code, focusing on its discretization and the derivation of the reduced MHD equations it uses, while Section 3.2 is devoted to the RE fluid model included in JOREK, with a special focus on the RE avalanche generation model. Section 3.3 details the coupling with the resistive wall code STARWALL [100, 101], which is essential for correctly modeling the VDE dynamics. Finally, Section 3.4 outlines the basic assumptions and physical setup of the simulations presented in the following chapters and the artificial thermal quench (ATQ) procedure, which is used to numerically initiate disruptions in a controlled manner.

3.1 The JOREK code

JOREK is a non-linear, resistive extended MHD code used for investigating magnetically confined plasmas. While it was originally developed to study instabilities in tokamak X-point geometries [99, 96], it has since evolved into a modular framework that supports a diverse range of physics models, including two-temperature fluids, neutrals, impurities, and REs. A thorough overview of these physical models and their extensive applications is provided in [96, 97] and their references.

The following sections detail the specific physical and numerical implementations adopted for the simulations in this thesis. Section 3.1.1 outlines the numerical discretization scheme, including the hybrid spatial representation and the time-stepping methods, while Section 3.1.2 summarizes the reduced MHD model, which forms the basis for the plasma evolution. A brief introduction to the construction of equilibria in JOREK is given in Section 3.1.3. The RE fluid model, which is self-consistently coupled to the reduced MHD equations and provides the computational basis for the RE avalanche, will be detailed in Section 3.2.

3.1.1 Spatial discretization and numerical implementation

JOREK employs a cylindrical coordinate system (R, Z, ϕ) , as illustrated in Figure 3.1. To solve the system of MHD equations described in the previous sections, the code employs a hybrid spatial discretization scheme that combines a 2D finite element grid in the poloidal direction and a Fourier expansion in the toroidal direction [96, 102]. In the toroidal direction (ϕ) , the physical variables are expanded in a Fourier series, exploiting the periodicity of the geometry. The number of harmonics n_{tor}

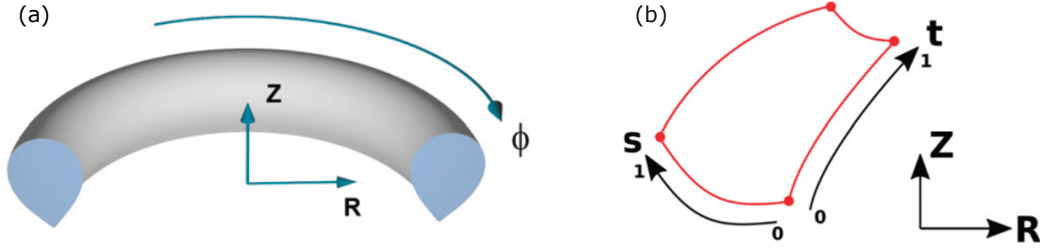


Fig. 3.1 (a) JOREK cylindrical coordinate system (R, Z, ϕ) used to describe the tokamak geometry. (b) Single 2D Bezier element in the (R, Z) space. The element-local coordinates are also shown: $(s, t) \in [0, 1]$. Adapted from [96].

can be selected for each simulation, with $n_{\text{tor}} = 0$ implying toroidal symmetry (i.e., 2D simulations in the poloidal plane) [96]. The poloidal plane (R, Z) is discretized using G^1 continuous Bezier finite elements consisting of cubic polynomials, although higher-order elements are available [96]. Since the coordinates use the same basis functions as the variables for their decomposition, these elements are also isoparametric. The equations are solved using the finite element method (FEM), which is briefly described in Appendix A. JOREK allows the use of a computational grid aligned with the magnetic flux surfaces of the initial equilibrium, which is more accurate for describing the anisotropy between parallel and perpendicular fluxes. Furthermore, it is possible to extend the grid to the PFCs, allowing for realistic wall geometries, which is crucial for SOL physics [96].

The system of equations is discretized in time using a fully implicit time-stepping scheme, typically utilizing the linearized Crank-Nicolson or the BDF2 Gears scheme [96]. By using an implicit scheme, it is possible to alleviate the Courant-Friedrichs-Lewy (CFL) condition, which restricts the largest usable time step. By writing the physical equations in the general form:

$$\frac{\partial \mathbf{A}(\mathbf{u})}{\partial t} = \mathbf{B}(\mathbf{u}, t) \quad (3.1)$$

where $\mathbf{u}$ represents the vector of physical variables, it is possible to advance the solution from time step n to $n + 1$ by linearizing it with respect to the increment $\delta \mathbf{u}^n = \mathbf{u}^{n+1} - \mathbf{u}^n$. The resulting general time-stepping scheme is given by:

$$\left[(1 + \xi) \left(\frac{\partial \mathbf{A}}{\partial \mathbf{u}} \right)^n - \Delta t \theta \left(\frac{\partial \mathbf{B}}{\partial \mathbf{u}} \right)^n \right] \delta \mathbf{u}^n = \Delta t \mathbf{B}^n + \xi \left(\frac{\partial \mathbf{A}}{\partial \mathbf{u}} \right)^{n-1} \delta \mathbf{u}^{n-1}, \quad (3.2)$$

where Δt is the time step size [96]. The specific numerical scheme is selected by the choice of the parameters θ and ξ :

- Setting $\theta = 0.5$ and $\xi = 0$ yields the Crank-Nicolson scheme.
- Setting $\theta = 1$ and $\xi = 0.5$ selects the BDF2 Gears scheme.

By combining high-order spatial discretization with robust time-stepping schemes, JOREK provides a reliable framework for simulating the complex, non-linear dynamics of disruptions and runaway electron generation.

3.1.2 The reduced MHD model

The foundation of the physical model in JOREK is the set of normalized MHD equations that describe the evolution of the magnetic potential $\mathbf{A}$, the plasma mean velocity $\mathbf{v}$, the total mass density ρ , and the total pressure p . The normalization is done with respect to μ_0 (vacuum permeability) and ρ_0 (initial mass density at the magnetic axis): while lengths are not normalized, time is normalized by a factor $\tau_{\text{norm}} = \sqrt{\mu_0 \rho_0}$, and the total density and pressure are normalized using ρ_0 and μ_0 , respectively. The normalization factors for the main physical variables are summarized in Appendix B [96]. The resulting equations take the form:

$$\frac{\partial \mathbf{A}}{\partial t} = -\mathbf{E} - \nabla \Phi, \quad (3.3)$$

$$\rho \frac{\partial \mathbf{v}}{\partial t} = -\rho \mathbf{v} \cdot \nabla \mathbf{v} - \nabla p + \mathbf{J} \times \mathbf{B} + \nabla \cdot \underline{\tau} + \mathbf{S}_v, \quad (3.4)$$

$$\frac{\partial \rho}{\partial t} = -\nabla \cdot (\rho \mathbf{v}) + \nabla \cdot (\underline{D} \nabla \rho) + S_\rho, \quad (3.5)$$

$$\frac{\partial p}{\partial t} = -\mathbf{v} \cdot \nabla p - \gamma p \nabla \cdot \mathbf{v} + (\gamma - 1) \nabla \cdot (\underline{\kappa} \nabla T) + (\gamma - 1) \underline{\tau} : \nabla \mathbf{v} + S_p \quad (3.6)$$

where Φ is the electric scalar potential, $\underline{\tau}$ represents the viscous stress tensor, while $\underline{D}$ and $\underline{\kappa}$ denote the particle and heat diffusion tensors, respectively. Crucially, these transport tensors can be highly anisotropic, distinguishing between transport parallel to the magnetic field lines and perpendicular transport. $\mathbf{S}_v$, S_ρ , and S_p represent source terms in the momentum, particle, and pressure equations, respectively. Unless specified otherwise (e.g., in a two-temperature model extension), the ion and electron temperatures are assumed to be equal ($T_e = T_i = T/2$). The total pressure is defined

as the sum of electron and ion pressures, $p = p_e + p_i = \rho T$. The system is closed by the definitions of the magnetic field $\mathbf{B}$ and the current density $\mathbf{J}$:

$$\mathbf{B} = F \nabla \phi + \nabla \times \mathbf{A} \quad (3.7)$$

$$\mathbf{J} = \nabla \times \mathbf{B} \quad (3.8)$$

where ϕ is the toroidal angle and $F \equiv RB_\phi$ is the toroidal flux function (R being the major radius), which varies spatially but is constant in time and is added for numerical reasons. The electric field $\mathbf{E}$ is determined by Ohm's law, which includes both drift-ordered (diamagnetic) and resistivity terms [96]:

$$\mathbf{E} = -\mathbf{v} \times \mathbf{B} + \eta(\mathbf{J} - \mathbf{J}_*) + F_0 \frac{\delta^*}{\rho} (\nabla_\perp p_i - \nabla_\parallel p_e). \quad (3.9)$$

Here, $\mathbf{J}_* = j_s \mathbf{e}_\phi$ (j_s and $\mathbf{e}_\phi$ being the toroidal current density source and the normalized toroidal basis vector, respectively) is an optional source term used to maintain the initial current profile or model bootstrap current, and η is the plasma resistivity, which follows the typical Spitzer dependence $\eta \propto T^{-3/2}$. The parameter δ^* is the diamagnetic coefficient, defined as:

$$\delta^* = \frac{m_i}{e F_0 \sqrt{\mu_0 \rho_0}} \quad (3.10)$$

where m_i is the ion mass and the constant F_0 is the toroidal flux function calculated with the major radius at the geometric center R_0 and considering the vacuum toroidal field $B_{\phi,0}$ [96].

Historically, reduced MHD models relied on an ordering parameter, typically the inverse aspect ratio $\epsilon = a/R$, to determine which terms of the full MHD equations to keep. It is possible to retrieve several sets of reduced equations depending on the ordering assumption, which consequently removes certain aspects of the physical phenomena under study. In contrast, the reduction in JOREK is Ansatz-based [103, 104]: instead of discarding terms based on their ordering with respect to a chosen parameter, this method postulates a specific structural form for the magnetic field and velocity variables that inherently eliminates the fast magnetosonic waves [105]. This approach allows for exact energy conservation at the equation level and retains more physical effects because of the lower number of assumptions. The drawbacks of this approach are that the equations become more complex and more difficult

to interpret. Moreover, it becomes harder to estimate errors without an ordering parameter. Nevertheless, it is still of great use because it allows for significantly faster computation, due to both the reduced number of variables and the ability to use larger time steps [96]. The reduction of MHD equations in JOREK is achieved through the fundamental assumption that the magnetic field is characterized by a dominant and approximately constant toroidal component B_ϕ , approximated as the vacuum toroidal field F_0/R (F_0 being constant in space and time) [106]. The magnetic field can then be written as:

$$\mathbf{B} = \frac{F_0}{R} \mathbf{e}_\phi + \frac{1}{R} \nabla \psi \times \mathbf{e}_\phi \quad (3.11)$$

where ψ is the poloidal flux. Equation (3.11) imposes a constraint on the velocity, which cannot be arbitrary. By defining the velocity stream function $u = \Phi/F_0$, the velocity can be written as:

$$\mathbf{v} = \underbrace{v_\parallel \mathbf{B}}_{\mathbf{v}_\parallel} + \underbrace{-R \nabla u \times \mathbf{e}_\phi}_{\mathbf{v}_E} + \underbrace{-\frac{\delta^* R}{\rho} \nabla p_i \times \mathbf{e}_\phi}_{\mathbf{v}_{\text{dia},i}} \quad (3.12)$$

where the first term is the velocity parallel to $\mathbf{B}$, the second term represents the $\mathbf{E} \times \mathbf{B}$ drift, and the last term captures the ion diamagnetic drift velocity [96].

By substituting the magnetic and velocity ansatz given by Equations (3.11)–(3.12) into the full MHD equations (Equations (3.3)–(3.6)), a closed system of equations is obtained. To avoid the numerical computation of higher-order spatial derivatives during the time-stepping, JOREK introduces two auxiliary variables: the toroidal current density j and the toroidal vorticity ω [96]. These are defined as:

$$j = \Delta^* \psi \equiv R^2 \nabla \cdot \left(\frac{1}{R^2} \nabla_{\text{pol}} \psi \right) \quad (3.13)$$

$$\omega = \Delta_{\text{pol}} u \equiv \nabla \cdot \nabla_{\text{pol}} u \quad (3.14)$$

where ∇_{pol} denotes the gradient in the poloidal plane: $\nabla_{\text{pol}} = \mathbf{e}_R \partial_R + \mathbf{e}_Z \partial_Z$ (∂_R and ∂_Z being the partial derivatives with respect to R and Z , respectively). Taking these additional variables into account, the base reduced MHD model solves for seven scalar variables: $\{\psi, u, j, \omega, \rho, T, v_\parallel\}$. The specific evolution equations for these variables are listed below [105, 106, 96].

1. By defining the Poisson bracket as $[a, b] = \nabla a \times \nabla b \cdot \mathbf{e}_\phi$, the evolution of the poloidal magnetic flux ψ can be written as:

$$\frac{1}{R^2} \frac{\partial \psi}{\partial t} = -\frac{1}{R} [u, \psi] - \frac{F_0}{R^2} \frac{\partial u}{\partial \phi} + \frac{\eta}{R^2} (j - j^*) + \frac{\delta^*}{2\rho} \nabla p \cdot \mathbf{B} \quad (3.15)$$

where j^* is the optional current source described earlier.

2. The poloidal momentum equation can be rewritten as an equation for the velocity stream function u :

$$\begin{aligned} R \nabla \cdot \left(R^2 \rho \frac{\partial \nabla_{\text{pol}} u}{\partial t} \right) &= [\rho R^4 \omega, u] - \frac{1}{2} [R^2 \rho, R^2 |\nabla_{\text{pol}} u|^2] \\ &\quad - [R^2, \rho T] + [\psi, j] - \frac{F_0}{R} \frac{\partial j}{\partial \phi} + \mu R^2 \Delta \omega \end{aligned} \quad (3.16)$$

The stress tensor has been simplified to a constant dynamic viscosity μ , even though it is possible to prescribe a different value in the parallel direction, $\mu_{\parallel}$.

3. For the parallel momentum, the equation is written for $v_{\parallel}$:

$$\begin{aligned} \rho B^2 \frac{\partial v_{\parallel}}{\partial t} + \frac{1}{R^2} \rho v_{\parallel} \nabla_{\text{pol}} \psi \cdot \nabla_{\text{pol}} \left(\frac{\partial \psi}{\partial t} \right) &= \frac{\rho}{2R} [\psi, v_{\parallel}^2 B^2] - \frac{\rho F_0}{2R^2} \frac{\partial (v_{\parallel} B)^2}{\partial \phi} \\ &\quad + \frac{1}{R} [\psi, p] - \frac{F_0}{R^2} \frac{\partial p}{\partial \phi} + \mu_{\parallel} B^2 \Delta v_{\parallel} \end{aligned} \quad (3.17)$$

where $\mu_{\parallel}$ is the value of dynamic viscosity parallel to $\mathbf{B}$.

4. The continuity equation for density ρ retains its full MHD formulation, but with the prescribed $\mathbf{B}$ and $\mathbf{v}$:

$$\begin{aligned} \frac{\partial \rho}{\partial t} &= R[\rho, u] + 2\rho \partial_Z u + 2\delta^* \partial_Z p - \frac{F_0}{R^2} \partial_\phi (\rho v_{\parallel}) - \frac{1}{R} [v_{\parallel} \rho, \psi] \\ &\quad + \nabla \cdot (D_{\parallel} \nabla_{\parallel} \rho + D_{\perp} \nabla_{\perp} \rho) + S_{\rho} \end{aligned} \quad (3.18)$$

where the particle diffusion tensor $\underline{D}$ has been divided into its components parallel ($D_{\parallel}$) and perpendicular ($D_{\perp}$) to $\mathbf{B}$. The former is only used when $v_{\parallel}$ is neglected, in order to avoid unphysical gradients in the parallel direction.

5. Like the continuity equation, the pressure equation retains its full MHD form as well, thus reading:

$$\begin{aligned}
\frac{\partial p}{\partial t} = & -v_{\parallel} \frac{F_0}{R^2} \partial_{\phi} p + \frac{v_{\parallel}}{R} [\psi, p] - R[u, p] \\
& - \gamma p \left(\frac{F_0}{R^2} \partial_{\phi} v_{\parallel} + \frac{1}{R} [v_{\parallel}, \psi] - \frac{2\delta^*}{\rho} \partial_z p + \frac{\delta^* R}{\rho^2} [u, \rho] - 2\partial_z u \right) \\
& + (\gamma - 1) \nabla \cdot (\kappa_{\parallel} \nabla_{\parallel} T + \kappa_{\perp} \nabla_{\perp} T) \\
& + (\gamma - 1) \underline{\tau} : \nabla (v_{\parallel} \mathbf{B} - R \nabla u \times \mathbf{e}_{\phi}) + S_p
\end{aligned} \quad (3.19)$$

As before, the heat diffusion tensor $\underline{\kappa}$ is divided into its parallel and perpendicular components ($\kappa_{\parallel}$ and $\kappa_{\perp}$, respectively). Typically, because of the high mobility of electrons along field lines, the parallel component is orders of magnitude larger and is usually modeled with a $T^{2.5}$ dependence, following the Spitzer-Härm formulation [3].

By combining these five equations (Equations (3.15)–(3.19)) with the definition equations for ω and j (Equations (3.13)–(3.14)), we obtain the seven equations that are simultaneously solved at every time step in the JOREK code [96].

JOREK allows for a flexible implementation of boundary conditions (BCs), supporting both Dirichlet BCs (where we fix the value of a variable at the boundary) and Neumann BCs (where the derivative normal to the boundary is specified). By default, for flux-aligned grids, the variables are fixed to their initial values at the boundary. Special attention is instead required for open field line regions intersecting the boundary, for instance considering grids that extend to the physical wall or the divertor region. Here, sheath boundary conditions are applied [96]: the parallel velocity component normal to the wall is set to the sound speed according to the Bohm criterion:

$$c_s = \sqrt{\frac{T_e + \gamma T_i}{m_i}} \quad (3.20)$$

where temperatures are in eV and $\gamma = 5/3$ is the adiabatic index for plasmas. The heat flux, instead, is governed by:

$$\mathbf{q} \cdot \mathbf{n} = \gamma_{\text{sh}} n_e c_s T_e \frac{\mathbf{B} \cdot \mathbf{n}}{B} \quad (3.21)$$

where $B = |\mathbf{B}|$ is the magnitude of the magnetic field and γ_{sh} is the sheath transmission factor, with typical values ranging from 7 to 8 [107]. The BCs also distinguish

between fixed-boundary cases, where the poloidal flux ψ and current j are constant in time on the JOEREK computational boundary (resembling an ideal wall condition), and free-boundary simulations, where these quantities evolve self-consistently. This free-boundary approach requires the coupling of JOEREK to a resistive wall code, which, in the framework of this thesis, is the STARWALL code, discussed in greater detail in Section 3.3.1.

3.1.3 Equilibrium construction

The initialization of the plasma state for the time-dependent simulations is performed using JOEREK's built-in Grad-Shafranov (GS) equilibrium solver [96]. It utilizes the same finite element grid and variable representation (Bezier elements) as the non-linear time evolution, described in Section 3.1.1. This consistency ensures that the discrete initial state accurately satisfies the equilibrium force balance, thereby avoiding discontinuous behavior that may arise from interpolation of equilibria from external codes.

The solver computes the equilibrium by solving the GS equation [108, 109] iteratively. The inputs required are the boundary conditions for the poloidal flux ψ and the profiles of temperature, density, and FF' . They are provided as functions of the normalized poloidal flux $\psi_n = (\psi - \psi_{\text{axis}})/(\psi_{\text{bnd}} - \psi_{\text{axis}})$, where ψ_{axis} and ψ_{bnd} are the values of ψ at the magnetic axis and on the plasma boundary, respectively [96].

After the equilibrium calculation, the reduced MHD variables are consistently initialized. The solution of the GS equation can be used to create a finite element grid aligned with the calculated flux surfaces, if needed.

3.2 The RE fluid model

To simulate the interaction between REs and the background plasma, the JOEREK code encompasses a hierarchy of RE models. While earlier studies primarily utilized test-particle approaches [110] or the fully coupled fluid model [111, 112] described here, very recent developments have introduced a new approach based on a full-f, full-orbit RE model self-consistently coupled to the MHD background

[58]. Although this model offers particular accuracy in regimes characterized by highly stochastic magnetic fields and is essential for detailed assessments of RE wall deposition patterns, the high dimensionality and temporal resolution required for such simulations impose significant additional computational costs, which can be prohibitive for long-timescale simulations of the entire disruption evolution.

In this thesis, we utilize the RE fluid model, whose formulation is primarily based on the development by Bandaru et al. [111] and its recent extension to include partially ionized impurities and deuterium neutrals [112]. By introducing the RE density n_{RE} as an additional physical variable, this approach significantly reduces the computational burden while retaining the essential physics of the mutual non-linear interaction between the RE beam and the MHD equations [96]. Unlike the test-particle model, the fluid formulation captures the EM back-reaction of the REs onto the bulk plasma, which is critical for the self-consistent modeling of the current quench phase of a disruption.

The efficacy and broad applicability of this RE fluid model have been demonstrated through its extensive application to RE beam generation and termination scenarios across major fusion devices, including JET [57], ITER [113, 114, 115, 92], and EU-DEMO [116]. These studies have utilized this model to address diverse physics objectives: from broad parameter scans in 2D predictive simulations [115], to highlighting the necessity of free-boundary conditions for accurately capturing avalanche multiplication [92]. Furthermore, utilizing 2D predictive modeling of RE formation across various impurity densities and 3D simulations of the RE beam termination against PFCs, the model has proven valuable for machine protection through the evaluation of the performance of sacrificial limiters [116].

The description provided in the remainder of this section summarizes the governing equations and physical assumptions employed in this work. For further details regarding the model formulation and implementation, the reader is referred to references [111, 112].

3.2.1 Governing equations for the fluid approximation

The fluid model treats REs as a distinct fluid species co-evolving with the background thermal plasma. Unlike full kinetic treatments, the RE fluid is characterized primarily by its number density, n_{RE} . The total current density $\mathbf{J}$ is decomposed into the Ohmic

current density carried by the thermal bulk plasma, $\mathbf{J}_{\text{th}}$, and the relativistic current density carried by the RE beam, $\mathbf{J}_{\text{RE}}$: $\mathbf{J} = \mathbf{J}_{\text{th}} + \mathbf{J}_{\text{RE}}$. The RE current density is given by:

$$\mathbf{J}_{\text{RE}} = -en_{\text{RE}}\mathbf{v}_{\text{RE}} \approx -en_{\text{RE}}\left(c\frac{\mathbf{B}}{B} + \frac{\mathbf{E} \times \mathbf{B}}{B^2}\right) \quad (3.22)$$

where the velocity of REs $\mathbf{v}_{\text{RE}}$ is assumed to have an $\mathbf{E} \times \mathbf{B}$ drift superimposed on the motion parallel to the magnetic field at the speed of light c .

The evolution of the RE density is governed by a continuity equation that accounts for both parallel advection and diffusive transport, as well as source mechanisms. Nevertheless, for simulations of disruptions involving stochastic magnetic fields (e.g., in disruption events), the parallel transport can be mimicked diffusively using a large RE parallel diffusivity $D_{\parallel, \text{RE}}$. Following this diffusive approach and the normalized reduced MHD equations introduced in Section 3.1.2, the evolution equation for n_{RE} is given by:

$$\frac{\partial n_{\text{RE}}}{\partial t} = R[n_{\text{RE}}, u] + 2n_{\text{RE}}\partial_Z u + \nabla \cdot (D_{\parallel, \text{RE}}\nabla_{\parallel} n_{\text{RE}} + D_{\perp, \text{RE}}\nabla_{\perp} n_{\text{RE}}) + S_{\text{RE}} + S_{\text{av}}, \quad (3.23)$$

where the notation is the same used in Section 3.1.2. The parallel diffusivity $D_{\parallel, \text{RE}}$ is set to a high value (typically $\sim 10^9 \text{ m}^2/\text{s}$) to mimic the fast transit of relativistic particles along the magnetic field lines. Conversely, the perpendicular diffusivity $D_{\perp, \text{RE}}$ is set at a much smaller value, primarily to ensure numerical stability [113]. The two source terms S_{RE} and S_{av} are, respectively, the primary and secondary generation mechanisms of REs, which are analyzed in the next section.

3.2.2 RE sources in the fluid model

The primary seed generation comprises tritium decay and Compton scattering from activated wall structures (where applicable). The Dreicer and hot-tail sources are currently not explicitly included in the standard RE fluid model formulation [112]. Consequently, they are typically treated as an initial condition or an explicit source injected during the thermal quench, representing the fraction of the pre-disruption current converted to RE current before the avalanche phase dominates. Given the strong sensitivity of the hot-tail generation to the kinetic details of the TQ, this approach introduces a significant degree of uncertainty regarding the exact magnitude of the initial seed population. To address this, a parametric scan of the initial RE

seed spanning several orders of magnitude is performed in the simulations presented in this work, ensuring the robustness of the results against variations in the seed amplitude.

The avalanche source, S_{av} , is modeled in JOREK accounting for the effects of partially ionized impurities. This extension is critical for simulations involving massive material injection, as bound electrons in high- Z impurities contribute to the collisional stopping power and the knock-on source term [70, 88], increasing the effective critical electric field and the avalanche growth rate compared to the standard Rosenbluth-Putvinski model [48], which assumed a fully ionized plasma. The avalanche source is defined as:

$$S_{av} = n_{RE}\Gamma = n_{RE} \frac{e}{m_e c \ln \Lambda} \frac{n_e^{\text{tot}}}{n_e} \frac{E_{\parallel} - E_c^{\text{eff}}}{\sqrt{4 + \bar{v}_s(p_*)\bar{v}_D(p_*)}}, \quad (3.24)$$

where $\ln \Lambda$ is the relativistic electron-thermal electron Coulomb logarithm, n_e is the free electron density, and n_e^{tot} includes both free and bound electrons: $n_e^{\text{tot}} = n_e + \sum_j n_j N_{e,j}$ ($N_{e,j}$ being the number of bound electrons for each species j). $\bar{v}_s$ and $\bar{v}_D$ are the normalized slowing down and deflection frequencies. The physics of the avalanche process is encapsulated in two key parameters that must be determined self-consistently: the effective critical electric field E_c^{eff} and the effective critical momentum p_* . They must be computed iteratively by solving a non-linear equation for each.

The slowing-down and deflection frequencies in Equation (3.24) are non-linear functions of the momentum p and, unlike in fully ionized models, they include sums over all impurity charge states to account for partial screening effects:

$$\begin{aligned} \bar{v}_s &= \frac{1}{\ln \Lambda_c} \left[\ln \Lambda^{\text{ee}} + \sum_j \frac{n_j}{n_e} N_{e,j} \left(\frac{1}{k} \ln(1 + h_j^k) - \beta^2 \right) \right] \\ \bar{v}_D &= \frac{1}{\ln \Lambda_c} \left[\ln \Lambda^{\text{ee}} + Z_{\text{eff}} \ln \Lambda^{\text{ei}} \right. \\ &\quad \left. + \sum_j \frac{n_j}{n_e} \left(\frac{2}{3} (Z_j^2 - Z_{0,j}^2) \ln [(p\bar{a}_j)^{3/2} + 1] - \frac{2}{3} \frac{N_{e,j}^2 (p\bar{a}_j)^{3/2}}{(p\bar{a}_j)^{3/2} + 1} \right) \right] \end{aligned}$$

In the above equations, $\ln \Lambda^{\text{ee}}$ and $\ln \Lambda^{\text{ei}}$ represent the electron-electron and electron-ion Coulomb logarithms, respectively. The summation index j runs over all impurity

charge-states present in the plasma. For each j , the number of bound electrons is given by $N_{e,j} = Z_j - Z_{0,j}$, where Z_j is the nuclear charge number and $Z_{0,j}$ is the ionization state (charge number of the ion). The effective ion charge is defined as $Z_{\text{eff}} = \sum_j n_j Z_{0,j}^2 / n_e$. The parameter $h_j = p\sqrt{\gamma - 1}/I_j$ relates the electron momentum p to the mean ion excitation energy I_j (normalized to the electron rest energy), while $\bar{a}_j$ represents the effective ion size for charge state j . Finally, $\beta = p^2/(1 + p^2)$ and $k = 5$ [112]. The effective critical momentum p_* , instead, is evaluated using:

$$p_* = \frac{\sqrt[4]{\bar{v}_s(p_*)\bar{v}_D(p_*)}}{\sqrt{E_{\parallel}/E_c}}, \quad (3.25)$$

where E_c is the Connor-Hastie critical field (Equation (2.27)). Since $\bar{v}_s$ and $\bar{v}_D$ depend on p_* , this equation is solved iteratively for p_* using a Newton-Raphson scheme from an initial guess, which proved to be a numerically robust method [112].

Simultaneously, the effective critical electric field E_c^{eff} , which defines the threshold for avalanche onset in Equation (3.24), is determined by solving a system of coupled equations:

$$\frac{E_c^{\text{eff}}}{E_c} \approx \bar{v}_{s0} + \bar{v}_{s1} \left[\left(1 + \frac{\bar{v}_{D1}}{\bar{v}_{D0}} \right) \ln \left(\frac{\bar{v}_{D0}}{2\bar{v}_{s1}} \right) + \sqrt{2\delta(E_c^{\text{eff}}) + 1} \right] \quad (3.26)$$

$$\delta(E_c^{\text{eff}}) = \frac{\bar{v}_{D0}}{\bar{v}_{s1}^2} \left[\frac{\bar{v}_{D0}\tau_{\text{syn}}^{-1}}{E_c^{\text{eff}}/E_c} + \phi_{\text{br}0} + \phi_{\text{br}1} \ln \left(\frac{\bar{v}_{D0}}{2\bar{v}_{s1}} \right) \right] \quad (3.27)$$

where τ_{syn} represents the normalized synchrotron radiation damping timescale, while the coefficients $\phi_{\text{br}0}$ and $\phi_{\text{br}1}$ describe Bremsstrahlung contributions dependent on the impurity mix. The slowing down and deflection frequencies are instead given by:

$$\begin{aligned} \bar{v}_{s0} &= 1 + \frac{1}{\ln \Lambda_c} \sum_j \frac{n_j}{n_e} N_{e,j} (\ln I_j^{-1} - 1), \\ \bar{v}_{s1} &= \frac{1}{2 \ln \Lambda_c} \left(1 + \sum_j 3 \frac{n_j}{n_e} N_{e,j} \right), \\ \bar{v}_{D0} &= 1 + Z_{\text{eff}} + \frac{1}{\ln \Lambda_c} \sum_j \frac{n_j}{n_e} \left((Z_j^2 - Z_{0,j}^2) \ln \bar{a}_j - \frac{2}{3} N_{e,j}^2 \right), \\ \bar{v}_{D1} &= \frac{1}{\ln \Lambda_c} \sum_j \frac{n_j}{n_e} Z_j^2. \end{aligned}$$

Although these expressions can be combined into a single cubic polynomial for the unknown ratio $x = E_c^{\text{eff}}/E_c$ (by eliminating $\delta(E_c^{\text{eff}})$), when the discriminant is negative, an iterative solution for x is mandatory, which is achieved using the Newton-Raphson method. Due to the potential for convergence difficulties in this non-linear regime, the implementation utilizes an array of initial guesses covering a wide range of values. The robust performance of this method has been consistently verified across all simulations performed with the model thus far [112].

3.2.3 Self-consistent coupling to MHD

The inclusion of the RE fluid necessitates specific modifications to the standard reduced MHD equations described in Section 3.1.2, ensuring a self-consistent description of the interaction between the RE beam and the background plasma. This is done via current-coupling, decomposing the total current density as done in Equation (3.22) and adding a correction term to the plasma momentum equation. While REs may carry a significant portion of the total current, their contribution to the mass density is negligible compared to the main ions.

Since REs are relativistic and effectively collisionless, they do not contribute to the resistive dissipation of the plasma. Consequently, the resistive terms are modified such that they depend solely on the current carried by the thermal bulk, $\mathbf{J}_{\text{th}} = \mathbf{J} - \mathbf{J}_{\text{RE}}$. This modification is crucial for capturing the physics of the CQ: as the thermal plasma cools and its resistivity η increases, the electric field rises, accelerating the REs. Eventually, the RE beam can carry the majority of the total current, effectively clamping the loop voltage and sustaining the discharge even in a highly resistive background plasma.

The correction to the MHD momentum equation caused by the presence of an RE population is due to their role in maintaining quasi-neutrality, which leads to an additional term in the force balance equal to $en_{\text{RE}}\mathbf{E} - \mathbf{J}_{\text{RE}} \times \mathbf{B}$. By utilizing the definition of the RE velocity vector $\mathbf{v}_{\text{RE}}$ used in Equation (3.22), this correction can be simplified by the fact that the Lorentz force component $-\mathbf{J}_{\text{RE}} \times \mathbf{B}$ partially cancels with the electric force term $en_{\text{RE}}\mathbf{E}$, leaving a net force directed along the magnetic field lines [111]:

$$en_{\text{RE}}\mathbf{E} - \mathbf{J}_{\text{RE}} \times \mathbf{B} \approx \frac{en_{\text{RE}}}{B} E_{\parallel} \mathbf{B}. \quad (3.28)$$

Incorporating this term, the corrected single-fluid momentum equation for the background plasma is:

$$\rho \frac{d\mathbf{v}}{dt} = -\nabla p + \mathbf{J} \times \mathbf{B} - \mathbf{v} S_\rho + \frac{en_{\text{RE}}}{B} E_\parallel \mathbf{B} \quad (3.29)$$

where S_ρ is the mass density source and d/dt is the material derivative (Equation (2.9)). It is important to note that after the TQ, when the parallel electric field $E_\parallel$ is large, the correction term can be of the same order of magnitude as the left-hand side, significantly altering the plasma dynamics [111].

3.3 Free boundary simulations

Accurately simulating VDEs requires a realistic description of the EM boundary conditions. A simple ideal wall would freeze the magnetic flux, preventing any vertical motion, while a pure vacuum boundary would fail to capture the stabilizing eddy currents that slow down the displacement. For the disruption and RE simulations presented in this thesis, accurately modeling this interaction is critical as the forces exerted by eddy currents in the DTT conductive structures play a fundamental role in determining the disruption dynamics and, consequently, the extent of the RE beam generation. To capture these effects, JOREK relies on a well-established coupling to the resistive wall code STARWALL [100, 101] as well as on a more recent coupling to the EM code CARIDDI [117, 118]. As mentioned in Section 3.1, this coupling replaces the standard fixed-boundary Dirichlet condition for the poloidal flux ψ and current density j with a physically rigorous free-boundary formulation. Their value at the computational boundary is allowed to evolve self-consistently, determined by the instantaneous distribution of currents inside the plasma and the induced eddy currents in the surrounding conductive structures.

The following subsections summarize the core numerical principles of the coupling scheme with the STARWALL code (Section 3.3.1) and detail the specific modeling assumptions used to represent the conductive geometry of DTT within the STARWALL framework (Section 3.3.2).

3.3.1 Coupling scheme with the STARWALL code

The coupling between JOREK and the resistive wall code STARWALL [101] is achieved using the virtual casing principle [119]. This formulation allows the EM BCs to evolve self-consistently, replacing the standard fixed-boundary condition with a dynamic response that accounts for the interaction between plasma currents and conductive structures. Effectively, the simulation domain is decomposed into a closed plasma region, solved by JOREK, and a vacuum region containing resistive walls, solved by STARWALL. Rather than solving for a full 3D volumetric current, STARWALL models the conductive structures as infinitesimally thin 2D triangular meshes. This section provides only the essential context required to understand the numerical coupling and its underlying physical approximations: for a complete and rigorous mathematical derivation, the reader is referred to the foundational works by Merkel et al. [100] and Hoelzl et al. [101].

According to the virtual casing principle [119], the magnetic field produced by the plasma current distribution inside the closed JOREK domain V_p can be represented by an equivalent surface current density, $\mathbf{j}_{\text{equiv}}$, flowing on the computational boundary surface δV_p (coupling surface). Using Ampère's law ($\nabla^2 \mathbf{A} = -\mu_0 \mathbf{J}$) and assuming the Coulomb gauge ($\nabla \cdot \mathbf{A} = 0$), this implies that the magnetic vector potential $\mathbf{A}$ evaluated at a position $\mathbf{r}_{\text{ext}}$ outside the plasma domain V_p can be computed by replacing the volume integral of the internal volumetric plasma current density $\mathbf{J}$ with a surface integral of the equivalent boundary current $\mathbf{j}_{\text{equiv}}$:

$$\mathbf{A}(\mathbf{r}_{\text{ext}}) = \frac{\mu_0}{4\pi} \int_{V_p} \frac{\mathbf{J}(\mathbf{r}')}{|\mathbf{r}_{\text{ext}} - \mathbf{r}'|} d^3 \mathbf{r}' = \frac{\mu_0}{4\pi} \int_{\delta V_p} \frac{\mathbf{j}_{\text{equiv}}(\mathbf{r}')}{|\mathbf{r}_{\text{ext}} - \mathbf{r}'|} dS' \quad (3.30)$$

To implement this numerically, the continuous integral relationship between the magnetic vector potential $\mathbf{A}$ and the equivalent surface currents must be discretized. By representing the coupling surface δV_p with 2D finite elements and applying FEM (Appendix A), this relation becomes a linear algebraic system:

$$\underline{\mathbf{M}}_{\text{pe}} \mathbf{A}^e = \underline{\mathbf{M}}_{\text{pp}} \mathbf{I}_p^{\text{equiv}} \quad (3.31)$$

where $\mathbf{I}_p^{\text{equiv}}$ and $\mathbf{A}^e$ denote the degrees of freedom for the equivalent surface current distribution and for the magnetic vector potential, respectively, mapped onto the JOREK finite element basis. The system relies on the self-inductance matrix $\underline{\mathbf{M}}_{\text{pp}}$

and $\underline{\mathbf{M}}_{pe}$, which projects the vector potential $\mathbf{A}^e$ onto the coupling surface [100]. A critical nuance in this formulation is that the vector potential $\mathbf{A}^e$ solved within JOEREK contains contributions from both the internal plasma and the external wall currents. To properly isolate the plasma-equivalent current $\mathbf{I}_p^{\text{equiv}}$, which is the term that actively drives the wall eddy currents, the external current contribution must be subtracted:

$$\mathbf{I}_p^{\text{equiv}} = \underline{\mathbf{M}}_{pp}^{-1} (\underline{\mathbf{M}}_{pe} \mathbf{A}^e - \underline{\mathbf{M}}_{pw} \mathbf{I}_w), \quad (3.32)$$

Here, $\mathbf{I}_w$ represents the degrees of freedom for the wall currents, and their specific coupling is captured by the mutual inductance matrix $\underline{\mathbf{M}}_{pw}$.

The time evolution of the currents outside the domain (the wall currents) is governed by a resistive-inductive equation, which accounts for the self-inductance ($\underline{\mathbf{M}}_{ww}$) and resistance ($\underline{\mathbf{R}}$) of the structures, as well as the voltage induced by the time-varying plasma currents ($\underline{\mathbf{M}}_{wp} \dot{\mathbf{I}}_p^{\text{equiv}}$):

$$\underline{\mathbf{M}}_{ww} \dot{\mathbf{I}}_w + \underline{\mathbf{R}} \mathbf{I}_w + \underline{\mathbf{M}}_{wp} \dot{\mathbf{I}}_p^{\text{equiv}} = \mathbf{U}_e. \quad (3.33)$$

where $\mathbf{U}_e$ represents the inclusion of externally applied voltages (e.g., from active coils). Because the matrices are strictly geometry-dependent, STARWALL computes them a single time during the initialization phase, allowing Equation (3.33) to be efficiently solved alongside the JOEREK time-step [101].

Finally, to close the JOEREK-STARWALL coupling, the tangential magnetic field $\mathbf{B}_{\text{tan}}$ at the computational boundary is calculated using the Biot-Savart law in the discretized FEM space, which reads:

$$\mathbf{B}_{\text{tan}}^e = \underline{\mathbf{M}}_{ep} \mathbf{I}_p^{\text{equiv}} + \underline{\mathbf{M}}_{ew} \mathbf{I}_w \quad (3.34)$$

where $\underline{\mathbf{M}}_{ep}$ and $\underline{\mathbf{M}}_{ew}$ respectively map the equivalent plasma surface currents and the external wall currents back to the tangential magnetic field on the JOEREK boundary. These matrices are specifically formulated to inherently account for the Fourier decomposition and the subsequent mapping onto the JOEREK finite element grid [120, 31]. Substituting Equation (3.32) into Equation (3.34) yields the final form:

$$\mathbf{B}_{\text{tan}}^e = \underline{\mathbf{M}}_{ep} \underline{\mathbf{M}}_{pp}^{-1} \underline{\mathbf{M}}_{pe} \mathbf{A}^e + (\underline{\mathbf{M}}_{ew} - \underline{\mathbf{M}}_{ep} \underline{\mathbf{M}}_{pp}^{-1} \underline{\mathbf{M}}_{pw}) \mathbf{I}_w. \quad (3.35)$$

which is used to dynamically solve the full free-boundary conditions of the simulations.

Although a more recent development extends this formalism to couple JOEREK with the EM code CARIDDI [117, 118], able to handle complex topological features such as 3D conducting structures and non-simply connected geometries, the simulations presented in this work rely on the well-established JOEREK-STARWALL coupling described above. The specific implementation of the DTT conductive structures in the STARWALL framework is detailed in the following section.

3.3.2 DTT structures in STARWALL

The description of the DTT conductive structures implemented in STARWALL follows the setup previously detailed by the author in [121]. This specific geometry is utilized for both the Day-0 (Chapter 4) and Full Power (Chapter 5) simulations analyzed in this thesis.

The numerical coupling described in Section 3.3.1 is applied here to the specific geometry of the DTT facility (Section 1.2). Accurate modeling of the conductive environment is essential for realistic disruption simulations, as the passive structures govern the timescales of VDEs, thereby directly influencing the RE beam generation and its final loss dynamics. A 2D poloidal cross-section of the implemented geometry is illustrated in Figure 3.2. For reference, the separatrix corresponding to the standard SN Day-0 equilibrium is also depicted in pink, visualizing the typical plasma volume relative to the vessel limits. Since the simulations discussed in this work are axisymmetric, the geometry of the conducting structures is treated with 2D toroidal symmetry as well, neglecting 3D features such as ports and the intrinsic toroidally varying nature of the components, which could introduce localized asymmetries. Further justification for the 2D approximation and its implications in the simulation is provided in Section 3.4.1. The model imported into STARWALL distinguishes between active coils and passive conducting structures. Due to the numerical constraints of STARWALL, which handles a single continuous surface with thin-wall triangles, the other structures are discretized as toroidal filaments. This results in the following configuration:

- **Vacuum vessel:** The two shells of the vacuum vessel are the primary passive conductors. The inner layer is modeled as a continuous axisymmetric thin wall (depicted in orange in Figure 3.2), while the outer shell is instead discretized using a set of toroidal filaments (red rectangles). The vessel material is stainless

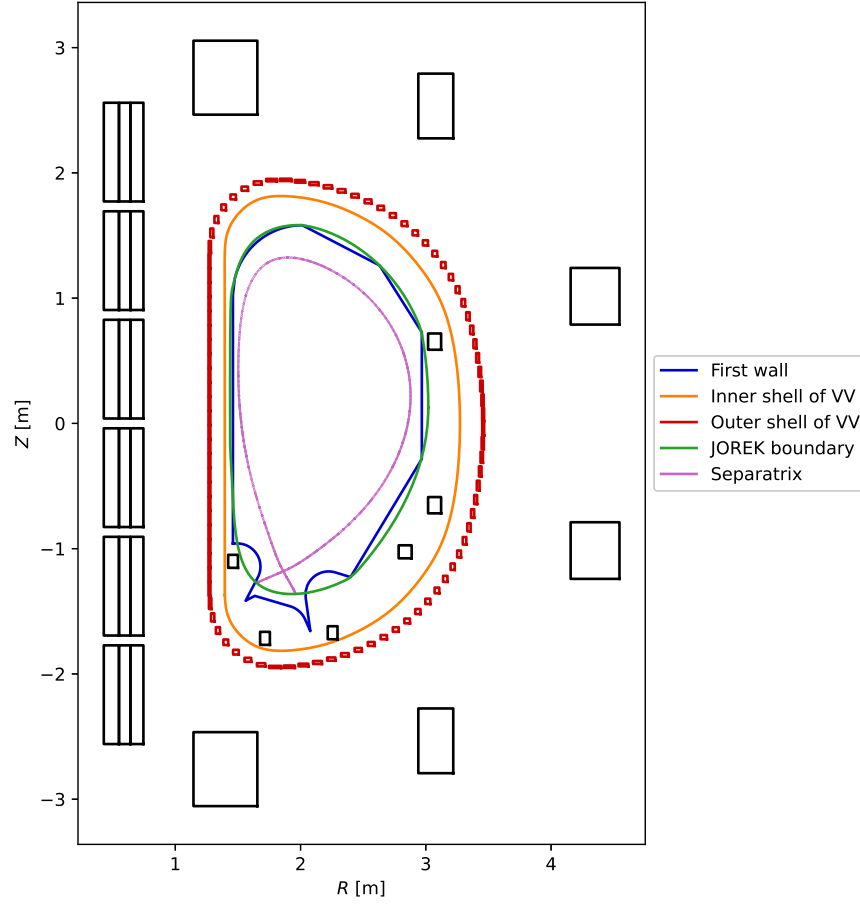


Fig. 3.2 Poloidal cross-section (R-Z plane) of the active and passive DTT conducting structures as implemented in the STARWALL code. The vacuum vessel is split into an inner shell (axisymmetric thin wall, orange line) and an outer shell (toroidal filaments, red). Active components, including the central solenoid, poloidal field coils, and in-vessel coils, are shown in black. The JOREK computational boundary (green) is defined to closely track the physical first wall (blue) while smoothing sharp corners to ensure numerical stability. The separatrix for the standard single null Day-0 scenario is included in pink for reference.

steel, characterized in the simulation by a resistivity $\eta_w = 0.75 \mu\Omega\text{m}$ and a uniform thickness of $d_w = 15 \text{ mm}$.

- **Active coils:** The EM environment includes the PF coils, the central solenoid, and the in-vessel coils, all shown in black in Figure 3.2. These are treated as current-carrying toroidal filaments that contribute to the magnetic equilibrium.

A critical aspect of the simulation setup is the definition of the JOREK computational boundary (shown in green in Figure 3.2). While the physical contour of the first wall

(FW, blue line) defines the material limit for the plasma volume, the computational boundary must remain numerically robust. Therefore, the boundary is constructed by smoothing out sharp geometric discontinuities to facilitate the numerical computation, while remaining as adherent as possible to the FW contour. This close proximity is crucial for the subsequent analysis, as it ensures that when the plasma touches the computational boundary during a VDE, it can be accurately interpreted as contact with the physical PFCs, given that the actual FW lies immediately beyond it. By solving Equation (3.33) using this geometric configuration, the simulations self-consistently evolve the plasma dynamics alongside the induced currents in the DTT structures.

3.4 Simulation setup and methodology

Having detailed the vacuum coupling via STARWALL in the preceding section, we now turn to the specific setup of the JOEREK simulations employed to study the RE dynamics in DTT. This chapter outlines the numerical framework applied to both the Day-0 (Chapter 4) and Full Power (Chapter 5) scenarios. While the specific initial equilibria and scenario-dependent parameters will be detailed in their respective chapters, the fundamental assumptions and methodology remain consistent across all simulations and are introduced here.

First, the specific model assumptions and the values of the main physical parameters used in this work are introduced in Section 3.4.1. Subsequently, Section 3.4.2 details the strategy employed to assess the sensitivity of the RE beam formation to the initial RE seed current and the impurity content, while Section 3.4.3 describes the artificial thermal quench (ATQ) procedure used to initiate the disruptions and establish the post-TQ conditions.

3.4.1 General modeling assumptions

The simulations presented in this work utilize the reduced MHD model available in JOEREK (Section 3.1.2), coupled with the RE fluid model (Section 3.2). The analysis is restricted to 2D (toroidally symmetric) simulations, thereby omitting non-axisymmetric perturbations. This choice is driven by several factors. First, 2D simulations provide a conservative upper estimate of the RE beam current:

since stochastic magnetic field lines in 3D simulations typically enhance RE losses (Section 2.3.4), a 2D approach represents a “worst-case” scenario regarding beam survival and current conversion. Second, the computational efficiency of the 2D model is essential to simulate the generation of an RE beam during the CQ for the wide range of parameters discussed in Section 3.4.2. Finally, these axisymmetric simulations may serve as a starting point for future 3D RE termination studies: as demonstrated in recent work [116], 2D simulations can be used to study the generation of a macroscopic RE beam, which is subsequently employed in 3D modeling to investigate its termination against PFCs.

Moreover, to reduce computational costs while capturing the essential macroscopic dynamics, several simplifications are applied to the background plasma fluid model:

- **Single-fluid approximation:** The plasma is treated as a single fluid where thermal equilibrium is assumed between ions and thermal electrons ($T_i = T_e$). Consequently, the total temperature is defined as $T = 2T_e$. This simplification allows for the resolution of global energy dynamics without the computational burden of the two-temperature extension available in JOREK.
- **Velocity simplification:** The parallel velocity of the background ions ($v_{\parallel}$) and the diamagnetic drift velocity (v_{dia}) are set to zero. Neglecting these terms significantly reduces the complexity of the model and the computational costs, since we can simplify Equation (3.12) to $\mathbf{v} \approx -R\nabla u \times \mathbf{e}_{\phi}$.
- **Absence of neutral particles:** The presence of neutral particles is neglected in this fluid description. Consequently, atomic physics processes such as charge exchange and wall recycling are not included in the model.
- **Impurity treatment:** Impurities are modeled as a distinct fluid species convected with the main plasma and characterized by a total mass density ρ_{imp} , which becomes an additional physical variable of the MHD model [122]. Consequently, the total mass density ρ is found by summing the plasma and the impurity densities. Additionally, impurities are assumed to be introduced in coronal equilibrium: although JOREK includes a marker-based module capable of evolving the full charge state distribution self-consistently [96, 123, 124], the coronal equilibrium assumption is chosen here to maintain computational tractability.

Regarding BCs, fixed Dirichlet conditions are applied to all fluid variables. However, as established earlier, the poloidal magnetic flux ψ and the total toroidal current density j at the boundary are determined self-consistently through the implicit coupling with the resistive wall code STARWALL, ensuring that the magnetic response of the conducting structures is accurately captured and enabling free-boundary simulations.

The physical coefficients used in the simulations are selected to balance physical accuracy with numerical stability. The electrical resistivity η_e and parallel thermal diffusivity $\chi_{\parallel}$ are modeled using standard Spitzer ($\propto T_e^{-3/2}$) and Spitzer-Härm ($\propto T_e^{5/2}$) scalings, respectively. However, to prevent numerical difficulties in regions of high temperature (pre-TQ), these coefficients are regularized using cut-off values. Cross-field transport for the thermal plasma is governed by a perpendicular thermal diffusivity $\chi_{\perp}$, which follows a prescribed radial hyperbolic tangent profile dependent on the normalized poloidal magnetic flux ψ_n . Regarding particle transport, the background plasma and impurities are governed by a constant isotropic particle diffusivity D . A dominant parallel diffusivity ($D_{\parallel, \text{RE}}$) is imposed to mimic the rapid transport of REs along magnetic field lines, which critically governs the RE loss rate upon contact with the PFCs. Its order of magnitude ($\sim 10^9 \text{ m}^2/\text{s}$) is anchored to optimal ranges identified via sensitivity studies in the literature for the same RE fluid model [112, 113]. Conversely, the perpendicular diffusivity ($D_{\perp, \text{RE}}$) is kept at a minimal value, with the objective of maintaining numerical stability without introducing artificial perpendicular diffusion [113]. Furthermore, it is important to clarify that while the behavior and order of magnitude of all the transport coefficients remain consistent across both scenarios, their specific physical values may differ slightly between the Day-0 and FP simulations. This is purely a consequence of using identical normalized input parameters, which scale differently based on the scenario (detailed in Appendix B). The specific physical values are summarized in the respective chapters (see Tables 4.1 and 5.1).

3.4.2 Parameter scanning strategy

A critical aspect of the methodology for both the Day-0 and FP scenarios is the systematic scanning of two physical quantities: the initial RE seed current and the amount of injected impurities. This approach allows for a comprehensive RE risk assessment in DTT.

1. **Initial RE seed:** The amplitude of the primary RE generation mechanisms (such as hot-tail and Dreicer generation) during the TQ is difficult to predict without dedicated simulations and can span several orders of magnitude [73]. Consequently, the initial RE seed current is treated as a free parameter in the framework of this thesis. By varying the seed current across a broad range (whose values are expressed as fractions of the pre-disruption plasma current), we comprehensively assess the sensitivity of the subsequent avalanche multiplication to the initial population in both DTT scenarios.
2. **Impurity content:** The quantity of impurities injected into the plasma is experimentally controlled by the disruption mitigation system. This parameter dictates the radiative cooling rate (thereby determining the CQ dynamics and the electric field magnitude) and influences the efficiency of the avalanche mechanism via knock-on collisions with bound electrons, as detailed in Section 2.3.3. Neon is selected as the impurity species for the simulations, as it is a primary candidate for the DTT DMS alongside argon [10].

Specific values for these scans will be provided in the dedicated chapters for each scenario. The objective of exploring the parameter space defined by these two variables is to characterize the RE avalanche dynamics across widely different regimes, thereby establishing the safety boundaries for DTT operations with respect to REs.

3.4.3 The artificial thermal quench (ATQ)

To investigate the formation of REs during the CQ, it is necessary to initialize the simulations with plasma conditions representative of a post-disruption state. While the JOREK code is capable of simulating the full, self-consistent TQ triggered by massive material injection in 3D [125, 126, 127, 128], such simulations are computationally demanding and are not essential for the objectives of this work, where the focus lies specifically on the CQ phase and RE avalanche generation. Consequently, we employ an artificial thermal quench (ATQ) methodology, which allows for the efficient and controlled establishment of post-TQ conditions (specifically low temperatures, flattened current profiles, and impurity injection), enabling a focused analysis of the subsequent RE generation and beam evolution. This procedure has already been exploited in past studies for different tokamaks [34, 113, 92, 116, 35], and

is designed to bypass the chaotic phenomenology of the TQ while preserving the macroscopic parameters required to study the CQ. The ATQ can be subdivided into three sequential steps, allowing the different aspects to be addressed independently, thereby facilitating the simulations. Importantly, throughout all these three steps, Ohmic heating is artificially disabled to prevent premature reheating of the plasma while the target conditions are being established:

1. **Thermal collapse:** In the first step, the perpendicular thermal diffusivity ($\chi_{\perp}$) is artificially increased by several orders of magnitude. This triggers a rapid collapse of thermal energy, driving the maximum electron temperature (T_e) down to a target value of approximately 10 eV, mimicking the cooling occurring during a TQ. Once the target temperature is reached, $\chi_{\perp}$ is reset to its physical value. In the temporal plots presented in subsequent sections, the conclusion of this phase is visually marked by a vertical dotted brown line.
2. **Current profile flattening:** To approximate the effects caused by stochasticity and magnetic reconnection during a disruption, a large hyper-resistivity in the core is temporarily applied. This serves to smooth gradients in the current density profile (j), flattening the safety factor profile q as well. This rapid flattening induces a transient I_p spike, a mechanism well-described in the literature [129, 130]. Upon achieving a sufficiently flattened current profile, the hyper-resistivity is returned to a negligible level for the rest of the simulations.
3. **Impurity injection:** In the final step, whose onset is indicated by a vertical dash-dotted brown line in the later time traces, neon impurities are introduced uniformly into the computational domain at a fixed constant rate of $\approx 1.35 \cdot 10^{22}$ particles/ms for both scenarios. This specific rate was chosen to ensure that the injection phase is rapid compared to the self-consistent CQ phase, while avoiding sharp transitions that could compromise numerical stability. Although, physically, impurity injection often precedes or triggers the TQ, in this artificial setup, impurities are introduced after the thermal collapse to establish the necessary power balance for the CQ, as they provide the radiative cooling required to counteract Ohmic heating during the subsequent phase. The injection continues until the desired total impurity content (a scanning parameter) is achieved.

Completion of these three steps concludes the artificial setup. At this point, Ohmic heating is reactivated, marking the onset of the self-consistent physical simulation during the CQ. The plasma then evolves freely according to the physical parameters and the interaction between the electric field, resistive dissipation, and RE generation. This approach ensures that the CQ and RE avalanche dynamics are studied in a controlled environment, initialized from a robust and reproducible post-disruption state.

Chapter 4

Runaway electron generation in the Day-0 scenario

The initial operational phase of DTT, referred to as the Day-0 scenario, represents a critical step in the commissioning of the machine. Although this early operational phase is characterized by a reduced plasma current of approximately $I_p = 2$ MA, significantly lower than the nominal 5.5 MA target of the Full Power scenario, RE risks cannot be dismissed entirely. Experiments have demonstrated that disruptions occurring at currents comparable to the DTT Day-0 level can still generate high-energy RE beams that result in significant damage to the first wall [61] (Figure 2.3). Consequently, a rigorous predictive assessment is required to confirm that the specific disruption dynamics of the Day-0 phase remain within safe limits, ensuring that the integrity of the PFCs is maintained as the facility progresses toward full power operations.

This chapter investigates the CQ dynamics and RE generation within the Day-0 framework. We explore a wide range of post-disruption conditions, systematically varying the quantity of injected impurities and the magnitude of the initial RE seed population, with the primary objective of defining the operational boundaries where macroscopic RE beams might form and quantifying the maximum potential RE currents that could threaten the device integrity during this early phase. Section 4.1 outlines the specific modeling assumptions adopted for this scenario. It details the starting magnetic equilibrium used for the simulations, the transport coefficients governing the plasma and RE fluid evolution, and the definition of the parameter

space regarding initial RE seeds and impurity injection levels. Subsequently, Section 4.2 presents the results of the non-linear MHD simulations. The analysis begins by examining the disruption dynamics under intermediate seed conditions, proceeds to investigate the high-amplitude seed scenario, and concludes with a summary of the RE generation risks across the full parametric scan. The analysis and results detailed throughout this chapter are based on the work published in [121].

4.1 Scenario specifications

4.1.1 Equilibrium

The simulations discussed in this chapter are based on a simplified single null (SN) configuration of the Day-0 scenario (scenario A1 [131]), where the plasma current is limited to approximately $I_p = 2$ MA and the vacuum toroidal magnetic field B_ϕ is set to 3 T. The magnetic equilibrium was derived using the CREATE-NL plasma equilibrium solver [132], ensuring an up-to-date representation of the DTT geometry.

In establishing the initial equilibrium for the simulations, a selection was made between two available datasets. The first option offered realistic, experimentally relevant temperature profiles but relied on an outdated geometrical description of the machine (including the first wall and coil positions) and presented a safety factor $q < 1$ in the core, implying immediate susceptibility to internal MHD instabilities. The second option, which was ultimately selected for this study, utilizes a simplified, bell-shaped temperature profile (orange line of the bottom graph in Figure 4.1) but retains the correct, modern machine geometry. The decision to prioritize geometrical accuracy and initial MHD stability ($q > 1$) over realistic profiles is justified by the specific methodology employed to trigger the disruption. As detailed in Section 3.4.3, the ATQ procedure initiates the event by forcing a rapid thermal collapse: in the first computational steps, the electron temperature is driven from keV levels down to approximately 10 eV via enhanced perpendicular diffusivity. Because this collapse is profound and relatively quick, the specific nuances of the pre-disruption temperature profile (whether simplified or realistic) become effectively irrelevant to the subsequent CQ dynamics. Furthermore, the choice of an initially stable q -profile avoids the numerical complexities associated with modeling pre-existing internal kink modes or sawteeth. Since the second step of the ATQ involves a flattening

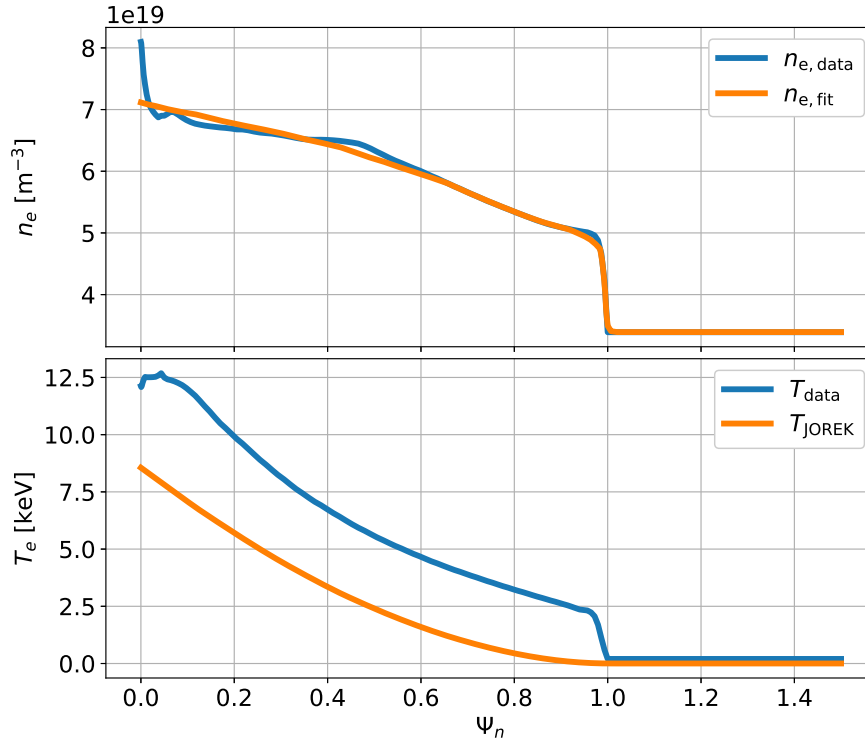


Fig. 4.1 Radial profiles of electron density (n_e , top) and temperature (T , bottom) as a function of the normalized poloidal flux ψ_n . The orange curves represent the equilibrium configuration implemented in the JOREK simulations for the Day-0 scenario. The density profile used in the simulation (orange) is a fit of the original data (blue). Regarding the temperature, the adopted equilibrium employs a simplified bell-shaped profile (orange) rather than the more complex profile of the older dataset (blue). This choice allows for the utilization of the most up-to-date DTT geometry while ensuring the initial magnetic equilibrium remains MHD stable ($q > 1$). The use of a simplified initial temperature profile is physically justified by the specific focus of this study on the CQ phase, during which the temperature evolution is determined self-consistently by the balance between Ohmic heating and impurity radiation, rendering the specific details of the pre-disruption temperature profile largely irrelevant to the final RE avalanche dynamics.

of the current density which inevitably reshapes the safety factor, starting from a stable $q > 1$ configuration allows for a clean initialization of the disruption physics without the overhead of managing core instabilities. Figure 4.1 illustrates JOREK input profiles for the electron density n_e and electron temperature T_e (orange lines). The initial plasma density is obtained through a fit of the equilibrium data for this scenario (blue line in the top plot). For completeness, the bottom graph also depicts the realistic temperature profile from the rejected, outdated equilibrium (blue line) to highlight the differences with the temperature of the chosen equilibrium.

Coeff.		Dependency	Value
η_e	$T_e < 500 \text{ eV}$	Spitzer $\propto T_e^{-3/2}$	$\eta_e = \eta_{e,\text{Sp}}$
	$T_e \geq 500 \text{ eV}$	Constant	$\eta_e^{\min} = \eta_{e,\text{Sp}} (@500 \text{ eV})$
$\chi_{\parallel}$	$T_e < 20 \text{ eV}$	Constant	$\chi_{\parallel}^{\min} = \chi_{\parallel,\text{SH}} (@20 \text{ eV})$
	$20 < T_e < 350 \text{ eV}$	Spitzer-Härm $\propto T_e^{5/2}$	$\chi_{\parallel} = \chi_{\parallel,\text{SH}}$
	$T_e \geq 350 \text{ eV}$	Constant	$\chi_{\parallel}^{\max} = \chi_{\parallel,\text{SH}} (@350 \text{ eV})$
$\chi_{\perp}$		Radial tanh profile	$\chi_{\perp,\text{core}} = 2 \text{ m}^2/\text{s}$
D		Constant	$1.23 \text{ m}^2/\text{s}$
$D_{\parallel,\text{RE}}$		Constant	$1.83 \cdot 10^9 \text{ m}^2/\text{s}$
$D_{\perp,\text{RE}}$		Constant	$1.83 \cdot 10^{-2} \text{ m}^2/\text{s}$

Table 4.1 Physical parameters adopted for the Day-0 simulations. The plasma resistivity (η_e) and parallel thermal diffusivity ($\chi_{\parallel}$) follow standard scalings but are regularized by constant cut-off values. The perpendicular thermal diffusivity ($\chi_{\perp}$) utilizes a radial hyperbolic tangent profile dependent on the normalized poloidal flux ψ_n , maintaining a core value and transitioning sharply to zero at $\psi_n = 0.92$ over a transition width of 0.01. The particle diffusivity (D) is isotropic for the thermal species, including impurities. In contrast, the RE fluid employs a highly anisotropic diffusion tensor. The value of the parallel component ($D_{\parallel,\text{RE}}$) was selected based on sensitivity studies of the RE fluid model in the literature, to optimally mimic fast parallel transport and minimize artificial perpendicular diffusion [112, 113]. The perpendicular component ($D_{\perp,\text{RE}}$) is maintained at a minimal level, in order to ensure numerical stability [113].

4.1.2 Physical parameters

The physical parameters governing MHD equations follow the general framework detailed in Section 3.4.1. The specific numerical coefficients adopted for the Day-0 scenario are summarized in Table 4.1. Regarding the temperature-dependent coefficients, the regularization cut-offs are set to 500 eV for electrical resistivity ($\eta_e^{\min}$), and 350 eV and 20 eV for parallel thermal diffusivity ($\chi_{\parallel}^{\max}$ and $\chi_{\parallel}^{\min}$, respectively). While $\eta_e^{\min}$ and $\chi_{\parallel}^{\max}$ are employed for numerical stability, $\chi_{\parallel}^{\min}$ is applied to avoid underestimating boundary heat fluxes [34, 35]. For particle transport, a constant isotropic diffusivity of $D = 1.23 \text{ m}^2/\text{s}$ is applied to the thermal species. The runaway electron fluid utilizes the optimized anisotropic diffusion tensor described in the methodology, with a dominant parallel component ($D_{\parallel,\text{RE}}$) to reproduce fast parallel

dynamics and a minimal perpendicular component ($D_{\perp, \text{RE}}$) to ensure numerical stability. The specific order of magnitude for $D_{\parallel, \text{RE}}$ was selected based on sensitivity studies in the literature for this specific RE fluid model, which demonstrate that it optimally mimics fast parallel advection without inducing spurious perpendicular diffusion [112, 113]. However, given the importance of this parameter in determining the RE loss rate once the plasma boundary touches the PFCs, sensitivity studies specifically devoted to the DTT environment could be explored in future work.

4.1.3 RE seed and impurity scan

Consistent with the scanning strategy defined in Section 3.4.2, we systematically vary the initial RE seed current and the impurity inventory to map the RE safety boundaries of the Day-0 scenario.

Regarding the initial RE population, three distinct seed currents were selected to bracket the possible primary generation regimes. These cases are defined within a relative framework for this parametric study and their values correspond to fixed fractions of the pre-disruption plasma current ($I_p = 2 \text{ MA}$):

- **Baseline seed scenario – 2 A** ($10^{-6} I_p$): Representing a minimal reference seed where primary generation is low.
- **Intermediate seed scenario – 200 A** ($10^{-4} I_p$): A transitional case serving to model moderate primary generation levels.
- **High-amplitude seed scenario – 20 kA** ($10^{-2} I_p$): An upper-bound worst-case scenario. This value implies an enhanced primary generation mechanism where a significant fraction (1%) of the total current is carried by relativistic electrons immediately following the TQ.

It is worth emphasizing that the selected upper bound of 20 kA already constitutes a massive suprathermal population. Any further increase in this parameter (e.g., to an order of magnitude higher) would represent a scenario where the device safety is threatened by the initial RE seed itself rather than by its subsequent amplification. Such conditions would shift the focus of the analysis from the avalanche dynamics during the CQ to the primary RE generation physics during the TQ, which falls outside the scope of this work.

Simultaneously, the content of the injected pure neon impurity, which is experimentally controllable via the DMS, is varied across the simulations. We define N_{imp} as the total number of injected neon particles and, for the Day-0 scenario, it ranges from approximately $1 \cdot 10^{19}$ to $3 \cdot 10^{21}$ particles within the computational volume ($\approx 52 \text{ m}^3$). This upper limit corresponds to an impurity content roughly equivalent to that of the background ions, marking a clear transition into a massive gas injection regime. This range provides a sufficiently broad parameter space to characterize the RE generation dynamics for this scenario of DTT. Exploring even more extreme impurity contents remains an open avenue for future studies.

4.2 Results

This section details the simulation results obtained for the DTT Day-0 scenario, analyzing the interplay between impurity content, initial RE seed currents, and the resulting CQ dynamics and RE generation. The investigation spans a range of conditions, from a benchmark case with no impurities or REs to scenarios with significant impurity injection and varying initial RE populations. A dedicated subsection for the baseline seed (2 A) is omitted because RE generation in this case remains negligible across all impurity levels, resulting in discharge dynamics that are virtually identical to the 200 A case. Consequently, the results for the 2 A seed are presented directly in the comparative summary (Section 4.2.3).

4.2.1 Intermediate RE seed (200 A)

We first examine the disruption evolution for an intermediate initial RE seed current of 200 A. Figure 4.2 displays the time traces for the plasma current (I_p) and RE current (I_{RE}), maximum electron temperature ($T_{e,\text{max}}$), and the particle inventories for ions and impurities.

The early phase of the simulation ($t \lesssim 1 \text{ ms}$) corresponds to the ATQ procedure, with its distinct phases marked by vertical brown lines in the time-evolution plots. The first phase concludes at the dotted line, marking the end of an initial drop in electron temperature to approximately 10 eV. Figure 4.3 illustrates this thermal collapse by comparing the poloidal cross sections of T_e at the simulation's onset and at the end of this initial phase. The second phase of the ATQ is characterized by the

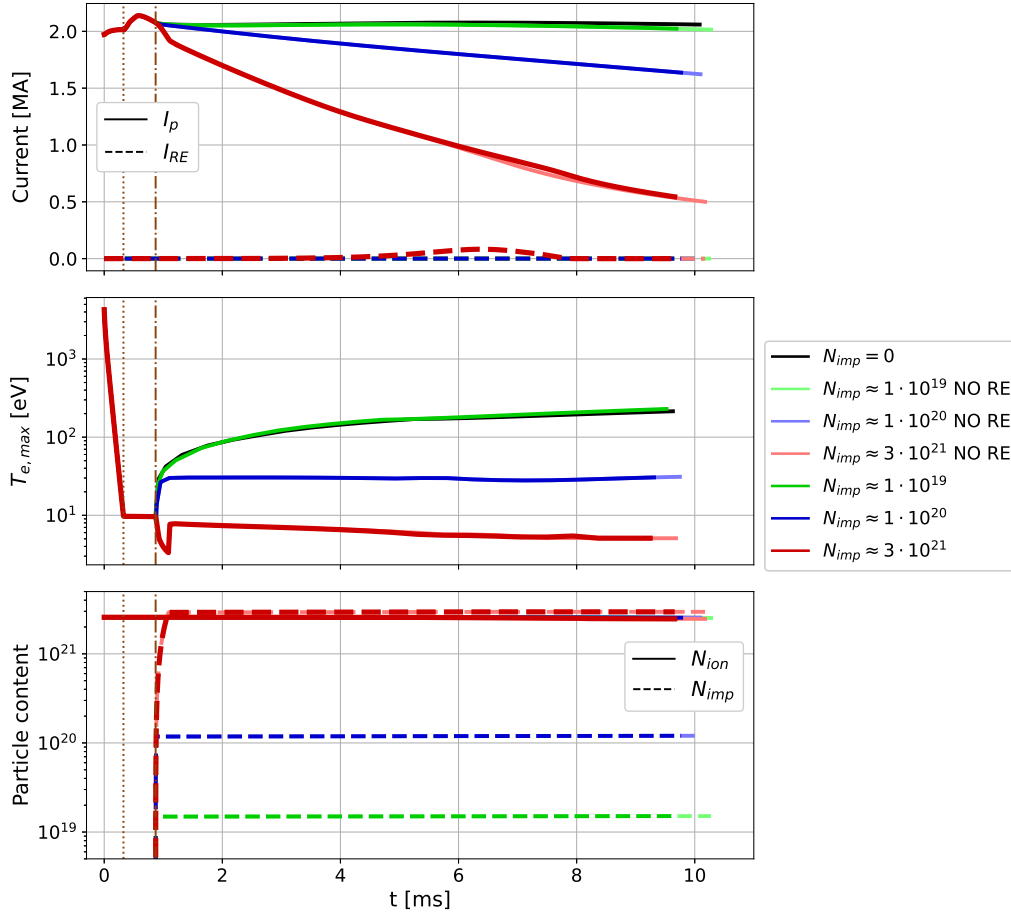


Fig. 4.2 Temporal evolution of key plasma parameters for the Day-0 scenario with an initial runaway electron (RE) seed of 200 A. From top to bottom: total plasma current (I_p , solid) and RE current (I_{RE} , dashed); maximum electron temperature ($T_{e,max}$); and total particle content for ions (N_{ion} , solid) and impurities (N_{imp} , dashed). Lighter shades of the same color indicate reference simulations performed without REs; these curves are largely obscured due to the negligible impact of the RE population at this RE seed level. The phases of the artificial thermal quench (ATQ) are delineated by vertical brown lines: the dotted line marks the end of the initial thermal collapse, while the dash-dotted line indicates the start of impurity injection. The sharp transient drop and subsequent recovery of $T_{e,max}$ around $t \approx 1$ ms for the highest impurity scenario (red curve) is due to this artificial setup, as Ohmic heating remains disabled until the impurity injection phase concludes.

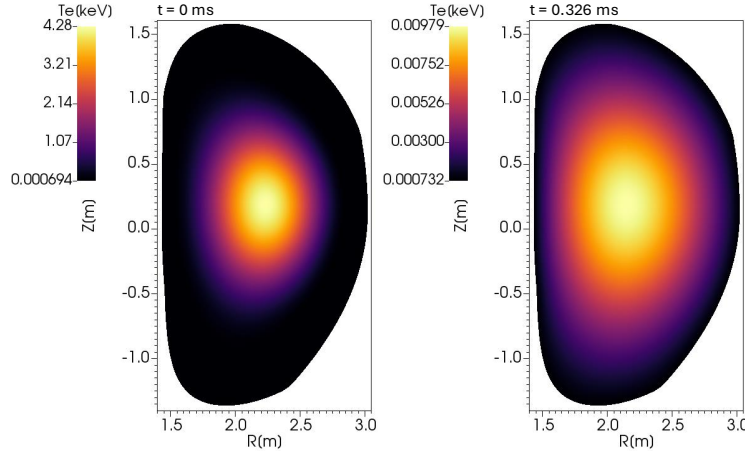


Fig. 4.3 Poloidal cross sections of the electron temperature (T_e) distribution at the beginning of the simulation ($t = 0$ ms, left) and at the end of the first step of the artificial thermal quench ($t = 0.326$ ms, right), a timestamp consistently marked by the brown dotted vertical lines in the time-evolution figures throughout this chapter. Despite the visual similarity of the relative profiles, note the difference in the color bar scales, highlighting a collapse in peak temperature from 4.28 keV down to approximately 10 eV.

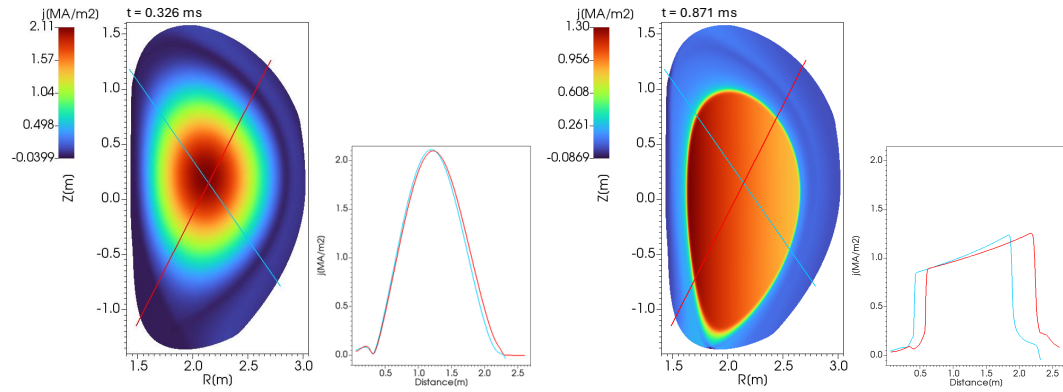


Fig. 4.4 Poloidal cross sections and corresponding 1D linear cuts of the current density (j) distribution. The left panels show the state at the end of the first artificial thermal quench step ($t = 0.326$ ms, marked by the brown dotted vertical lines in the time-evolution figures throughout this chapter), while the right panels show the state at the end of the second step ($t = 0.871$ ms, marked by the corresponding dash-dotted lines), immediately before impurity injection. The linear cuts clearly illustrate the flattening of the current profile inside the core during this phase.

flattening of the current profile and its associated I_p spike. Figure 4.4 highlights this current flattening inside the core through poloidal cross sections and linear cuts of the current density (j) at the start (dotted line) and at the end (dash-dotted line) of the second phase. The dash-dotted line also marks the onset of the third ATQ phase: impurity injection. During this period, the maximum electron temperature ($T_{e,\max}$) undergoes a further transient drop, which is most pronounced in the highest impurity case (red line) due to the extended duration of the injection phase. This drop is a direct consequence of the ATQ setup introduced in Section 3.4.3: as Ohmic heating remains artificially disabled while large quantities of impurities are introduced, the plasma experiences unopposed radiative cooling. Once the target impurity content is reached, the ATQ phase concludes and Ohmic heating is reactivated. Consequently, $T_{e,\max}$ quickly recovers to a physically self-consistent level dictated by the new balance between Ohmic heating and radiative losses. From this point onward, as the self-consistent CQ begins, the plasma current trajectories diverge depending on the injected quantity N_{imp} . For the lowest impurity case (green line), the plasma current decay is protracted, closely mirroring the impurity-free benchmark. This is a direct consequence of the insufficient radiative cooling, which fails to keep the electron temperature at levels that would maintain a high plasma resistivity. Conversely, higher impurity concentrations (e.g., red lines) drive strong radiative losses, clamping T_e to low values and enhancing resistivity, thereby accelerating the CQ.

Crucially, for this 200 A seed, the presence of REs has a negligible effect on the global CQ dynamics. The curves for simulations with REs (standard colors) and without REs (lighter shades) are virtually indistinguishable. Even at the highest impurity level ($N_{\text{imp}} \approx 3 \cdot 10^{21}$), where the CQ is the fastest and the electric field is strongest, the maximum I_{RE} remains strictly below 0.1 MA, indicating that the avalanche gain is insufficient to produce a macroscopic beam from this intermediate seed level.

4.2.2 High-amplitude RE seed (20 kA)

To probe the upper limits of RE avalanche risk in this scenario, simulations were performed with a significantly larger initial seed of 20 kA. The resulting current evolution is presented in Figure 4.5.

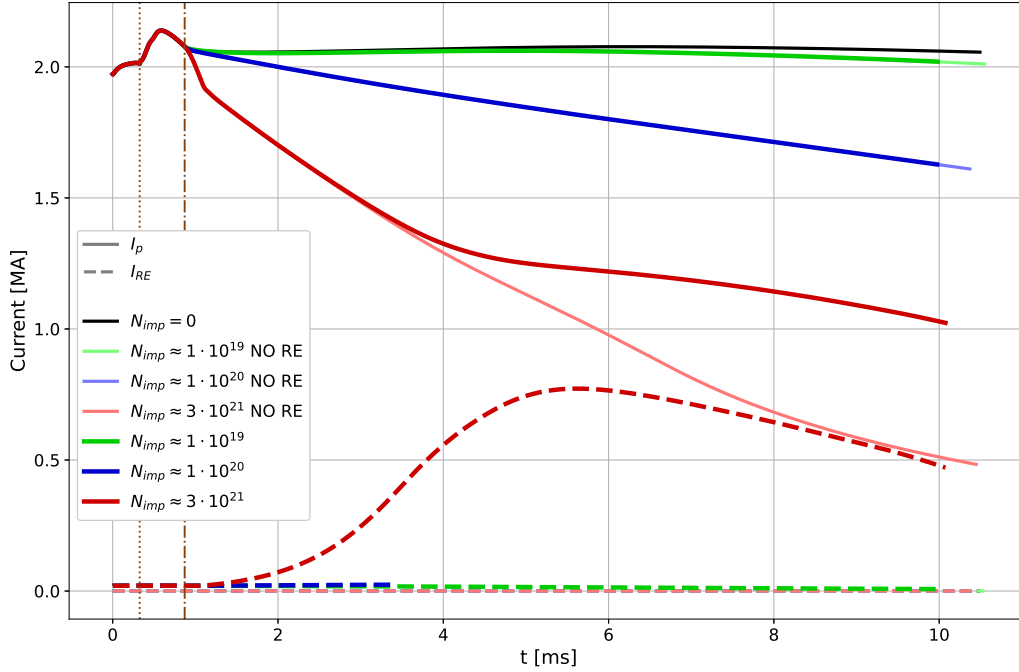


Fig. 4.5 Time traces of plasma current (I_p , solid) and runaway electron (RE) current (I_{RE} , dashed) for the high-amplitude initial seed case (20 kA). Varying impurity levels are shown by different colors. Lighter shaded lines represent reference cases without REs and the phases of the artificial thermal quench (ATQ) are delineated by vertical brown lines: the dotted line marks the end of the initial thermal collapse, while the dash-dotted line indicates the start of impurity injection. The end of the ATQ and the start of the self-consistent CQ phase is discernible by the distinct change in the I_p slope immediately following the dash-dotted line. Unlike the lower seed cases, a divergence between RE-inclusive and RE-free simulations is observed at the highest impurity level (red lines), where the growth of an RE beam notably alters the current decay dynamics.

A comparison with the previous case reveals a distinct change in behavior for the highest impurity scenario (dark red curves). While lower impurity levels (green and blue) still result in RE currents that do not affect the macroscopic I_p evolution, the combination of the highest seed (20 kA) and highest impurity content ($3 \cdot 10^{21}$) leads to a notable RE beam. In this specific case, I_{RE} peaks at approximately 0.8 MA. This significant RE population carries enough current to visibly slow the CQ compared to the RE-free reference case (lighter red line). Despite the strong RE growth, the beam does not carry the full plasma current. This is attributed to the presence of substantial halo currents in the SOL: the high impurity injection leads to comparable temperatures in the core and SOL, allowing large currents to be induced in the open field line region during the vertical displacement. These thermal halo currents

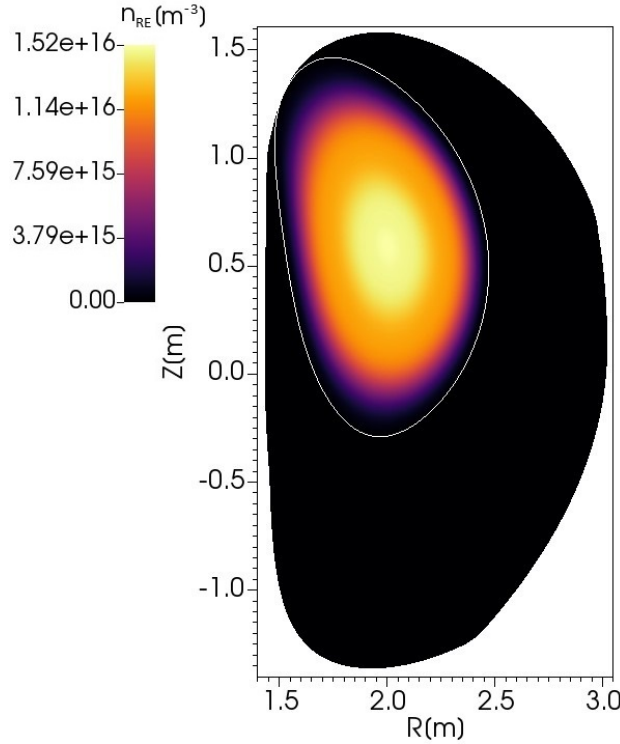


Fig. 4.6 Poloidal cross-section showing the runaway electron (RE) density distribution at the time of peak RE current ($t \approx 5.7$ ms) for the 20 kA RE seed and maximum impurity case. The white contour marks the separatrix, demonstrating that the REs are contained within the confined volume, which is being reduced by the vertical motion of the plasma.

cannot be converted into runaways because they flow in the open field line region. Consequently, they prevent I_{RE} from fully replacing the total plasma current I_p .

Simultaneously, the plasma experiences a vertical displacement that eventually forces it into contact with the PFCs. Figure 4.6 presents a snapshot of the poloidal RE density (n_{RE}) at the moment of peak RE current, demonstrating that the REs remain concentrated within the shrinking confined plasma volume bounded by the separatrix (indicated by the white line). A detailed analysis of these vertical movement dynamics and their full implications for RE generation will be discussed in the next section.

4.2.3 Summary of parameter scan

A consolidated view of the RE current evolution across all simulated scenarios is provided in Figure 4.7. To facilitate a direct comparison, the RE currents obtained for

the baseline seed (2 A, dotted lines) have been scaled by a factor of 100. Furthermore, the bottom panel introduces the time traces of the vertical position of the magnetic axis (Z_{axis}), illustrating the interplay between RE generation and the vertical motion of the plasma.

The results reveal distinct regimes of RE generation and varied VDE dynamics depending on the injected impurity content and starting RE seed. As the impurity content increases (and consequently the CQ and electron current drop faster; see Figure 4.2), the loss of vertical stability of the plasma shifts from a downward VDE to an upward VDE.

For the two lowest impurity levels ($N_{\text{imp}} \approx 1 \cdot 10^{19}$ and $1 \cdot 10^{20}$, represented by the green and dark blue lines), the plasma undergoes a downward vertical displacement, ultimately striking the PFCs at approximately 4.8 ms and 5.7 ms, respectively. In the lowest impurity case ($1 \cdot 10^{19}$), the avalanche mechanism is entirely absent. For the $1 \cdot 10^{20}$ case, a minor RE multiplication is initiated, but this marginal growth is immediately halted upon contact with the PFCs.

The intermediate impurity scenario ($N_{\text{imp}} \approx 1 \cdot 10^{21}$, represented by the light blue lines) marks a critical transition point. Here, the vertical position remains relatively stable for approximately 5 ms before rapidly accelerating downward, leading to a complete collapse in less than 2 ms. Notably, the RE avalanche in this scenario is not terminated by direct physical contact with the PFCs. Instead, its growth is suppressed by the severe shrinkage of the confined plasma volume resulting from the rapid loss of vertical stability as the plasma moves towards the X-point. This dynamic is clearly illustrated in Figure 4.8.

For the baseline and intermediate seed scenarios (2 A and 200 A), the system exhibits a linear response: the scaled 2 A curves (dotted) virtually superimpose on the 200 A traces (dashed lines), indicating that the resulting RE current scales proportionally to the initial seed magnitude. At high impurity levels (red and light blue curves), the traces exhibit a constant positive slope on the semi-logarithmic scale, confirming that the exponential avalanche mechanism is active during the initial phase of the CQ. However, starting from such low initial seeds, this amplification is insufficient to produce a macroscopic RE beam. For instance, in the highest impurity case ($N_{\text{imp}} \approx 3 \cdot 10^{21}$, red lines), the RE avalanche originating from the 200 A seed is terminated by the upward VDE, which drives the plasma into the PFCs and scrapes off the growing RE beam (indicated by the red 'x' marker on the dashed lines). The

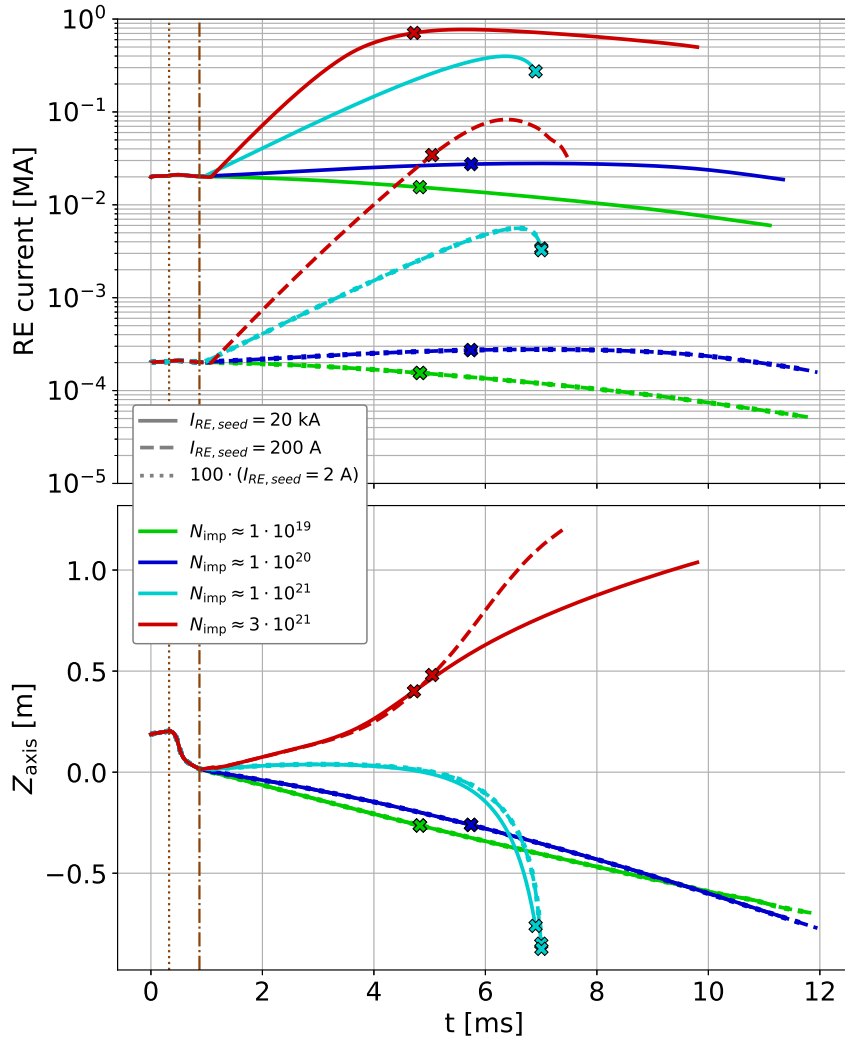


Fig. 4.7 Time evolution of the runaway electron current (top) and the vertical position of the magnetic axis (bottom) for the three initial seed regimes: baseline (2 A, dotted), intermediate (200 A, dashed), and high-amplitude (20 kA, solid). To improve readability and highlight scaling trends, the I_{RE} values for the 2 A seed (dotted lines) have been multiplied by a factor of 100. Color coding indicates the impurity injection levels, while the phases of the artificial thermal quench are delineated by vertical brown lines: the dotted line marks the end of the initial thermal collapse and the dash-dotted line indicates the start of impurity injection. The 'x' markers denote the timestamps when the plasma first contacts the plasma-facing components. The constant slopes clearly observed for high impurity levels (red and light blue) reflect exponential avalanche growth. The superposition of the scaled 2 A and 200 A traces points to a linear dependence on the initial seed magnitude, whereas the 20 kA case reveals non-linear effects. Note that the simulation for the 2 A seed at the highest impurity level (which would appear as a red dotted line) was omitted, as its negligible macroscopic impact is already bounded by the corresponding 200 A case.

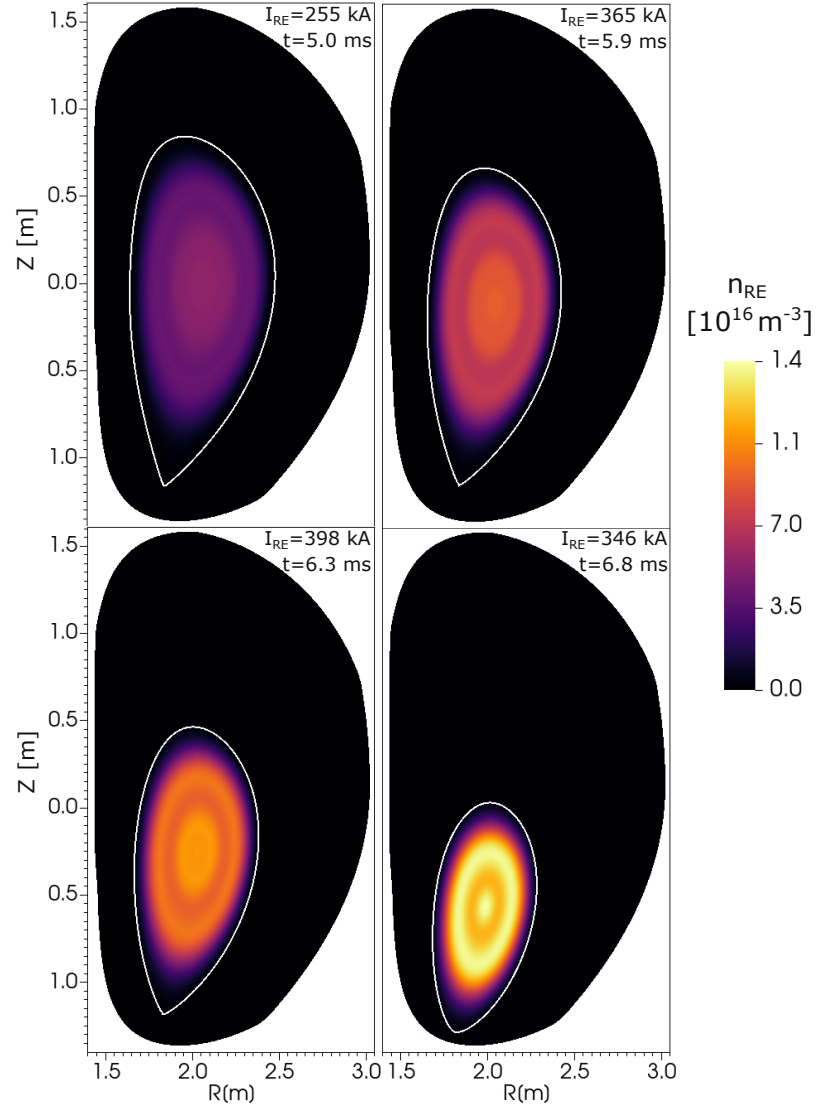


Fig. 4.8 Evolution of the runaway electron density profile (n_{RE} , expressed in 10^{16} m^{-3}) for the high-amplitude seed scenario ($I_{\text{seed}} = 20 \text{ kA}$) with an intermediate impurity content of $N_{\text{imp}} \approx 1 \cdot 10^{21}$ (corresponding to the light blue solid traces in Figure 4.7). The white contour denotes the separatrix, distinguishing the confined plasma from the open field line region. The sequence illustrates the volumetric collapse mechanism: as the vertical instability drives the plasma rapidly downward towards the X-point, the confined volume is progressively reduced. Consequently, although the local RE density in the core continues to rise due to the avalanche mechanism, the total RE current (I_{RE} , indicated at the top right of each panel) eventually decays after reaching its peak at $t \approx 6.3 \text{ ms}$, as the reduction in confined volume outpaces the avalanche growth prior to direct contact with the PFCs.

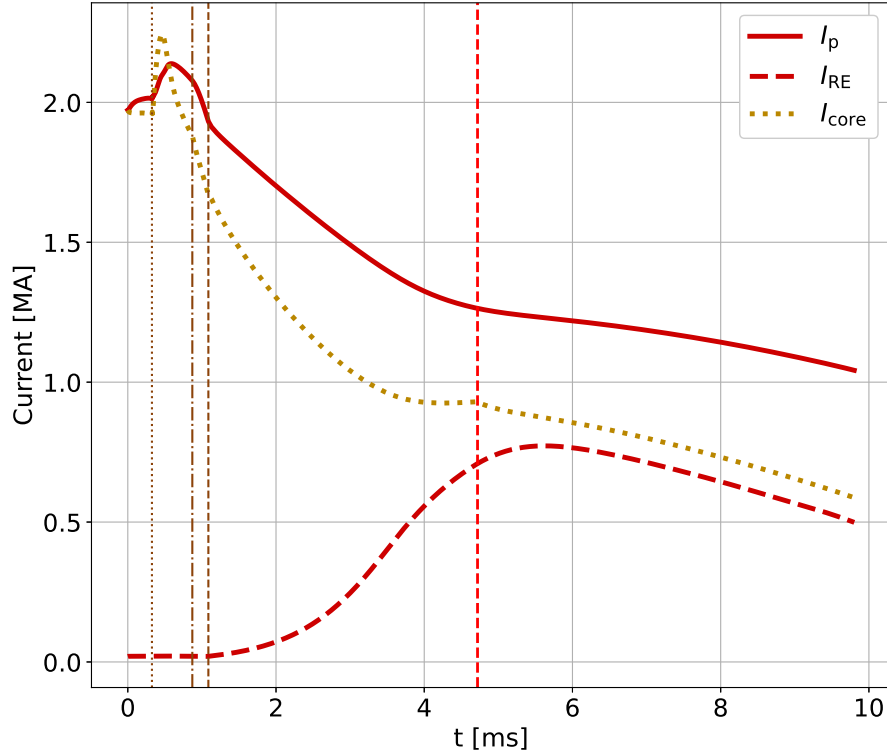


Fig. 4.9 Current decomposition for the high-amplitude seed scenario ($I_{seed} = 20$ kA) with an impurity inventory of $N_{imp} \approx 3 \cdot 10^{21}$ (corresponding to the solid red traces in Figure 4.7). The plot displays the total plasma current (I_p , solid line), the runaway electron (RE) current (I_{RE} , dashed line), and the total current flowing within the closed flux surfaces (I_{core} , gold dotted line). The phases of the artificial thermal quench are delineated by thin vertical brown lines: the dotted line marks the end of the initial thermal collapse, the dash-dotted line indicates the onset of impurity injection, and the dashed line signals the completion of the ATQ procedure and the start of the self-consistent current quench. Notably, during the artificial current flattening phase (between the dotted and dash-dotted vertical lines), a transient I_p spike occurs alongside the development of a negative edge skin current. An additional red dashed vertical line marks the timestamp of first contact with the plasma-facing components (PFCs). The exponential growth of the RE avalanche deviates from linearity and approaches saturation as I_{RE} nears the available confined current limit (I_{core}), an effect manifesting prior to PFC contact. The subsequent convergence of the I_{core} and I_{RE} curves demonstrates that the current within the confined plasma volume is eventually carried almost entirely by the RE beam, implying that the remaining current ($I_p - I_{RE}$) consists predominantly of halo currents flowing in the scrape-off layer.

dynamics of this upward VDE and the subsequent wall-scraping mechanism are highly analogous to those encountered in the Full Power scenarios. Consequently, a detailed 2D visualization of this specific interaction is deferred to the following chapter (Figure 5.8). Simulations for the baseline seed were not extended to the highest impurity level (red curves), as its macroscopic impact is already bounded by the intermediate case, which reaches less than 0.1 MA for the $N_{\text{imp}} \approx 3 \cdot 10^{21}$ case and does not visibly alter the CQ dynamics (see Figure 4.2). A 2 A seed under the same conditions would yield a peak RE current below 1 kA, which is practically negligible for the macroscopic evolution of the discharge.

A notable shift in the CQ dynamics becomes evident when analyzing the high-amplitude initial seed scenario (20 kA, represented by the solid red line) at the highest impurity level. Unlike the intermediate seed case, where the avalanche is terminated by direct physical contact with the wall, the high-amplitude seed scenario deviates from exponential growth prior to plasma contact with the PFCs (indicated by the red 'x' marker on the solid lines). This early limitation occurs because a significant fraction of the available thermal electron population within the core has already been converted into runaways. Consequently, the avalanche cannot sustain its exponential growth rate as the pool of thermal electrons inside the confined region is rapidly depleted, while electrons residing on open field lines cannot be converted due to prompt transport losses. As clearly illustrated in Figure 4.9, the majority of the total confined core current (I_{core}) is already carried by the RE population prior to wall impact (marked by the vertical dashed red line). This dynamic is also entirely consistent with the electric field suppression mechanism introduced in Section 2.3.3, where the rapid replacement of Ohmic current by the RE beam clamps the driving inductive field [89].

In conclusion, the Day-0 analysis suggests a relatively low risk of sustained RE beam formation during initial DTT operations within the scope of the simulated scenarios. While a plasma current of $I_p \approx 2$ MA is traditionally sufficient to generate damaging RE beams in contemporary tokamaks [61], our simulations indicate that the specific VDE dynamics and rapid volume shrinkage characteristic of this configuration play a crucial role in halting RE beam growth when the initial seed is within the baseline or intermediate regimes ($\lesssim 200$ A). Indeed, a macroscopic RE beam (≈ 0.8 MA) is observed only in the most pessimistic region of the parametric space analyzed here, which combines the high-amplitude seed scenario with massive impurity injection. Under all other investigated conditions, RE generation remains

minimal because the inherent dynamics of the Day-0 scenario naturally suppress beam growth through a combination of early PFC contact during VDEs, rapid volume shrinkage, or core current depletion limits (the latter being intrinsically more restrictive due to the smaller initial current compared to FP scenarios). Although these conclusions are bounded by the assumptions of the current numerical model, they suggest that by avoiding excessive impurity injection, the Day-0 operational phase can be operated safely with a significantly mitigated risk of RE-induced damage.

Chapter 5

Runaway electron generation in the Full Power scenario

Following the assessment of the initial Day-0 phase, this chapter extends the investigation to the DTT SN Full Power (FP) configuration (Scenario E1 [131]). This operational regime represents the target performance of the machine, characterized by a nominal plasma current of $I_p = 5.5$ MA, a toroidal magnetic field of $B_\phi = 5.85$ T, and an auxiliary heating power of approximately 45 MW.

As explained in Section 2.3.3, transitioning from 2 MA to 5.5 MA introduces a fundamental shift in the RE avalanche during the CQ, as the theoretical avalanche gain G_{av} scales exponentially with the plasma current. Consequently, the safety margins observed in the Day-0 scenario cannot be assumed to hold for FP operations. At these higher currents, the potential for converting even a small RE seed into a macroscopic RE beam increases dramatically, posing a significantly greater threat to the PFCs of DTT. This chapter aims to quantify these risks by performing a comprehensive parametric scan of impurity levels and initial RE seed currents, specifically tailored to the conditions of the FP scenario.

The chapter is organized as follows: Section 5.1 outlines the modeling setup for the FP scenario, detailing the modifications made to the magnetic equilibrium to ensure stability and defining the range of the scanning parameters. Section 5.2 presents the results of the JOREK simulations. The analysis is structured by seed magnitude, progressing from baseline seeds to test the avalanche threshold, to intermediate and high-amplitude scenarios. Subsequently, Section 5.3 provides a

comprehensive discussion and general assessment of the RE risks in DTT, featuring a direct comparison between the Day-0 and FP regimes that highlights the critical role of the plasma current in determining the final avalanche gain. The core analysis and results detailed in this chapter form the basis of the work collected in [133].

5.1 Scenario specifications

5.1.1 Equilibrium

Unlike the simplified Day-0 case, standard integrated modeling for this high-performance scenario typically predicts a safety factor $q < 1$ across a significant portion of the plasma core, creating susceptibility to large sawtooth crashes [10]. Although work on sawtooth control optimization is ongoing, here we employ a modified magnetic equilibrium based on the CREATE-NL solution [132] to isolate the CQ dynamics from these pre-existing instabilities. Figure 5.1 compares the original and modified profiles for electron density (n_e), electron temperature (T_e), FF' (F being the toroidal flux function), and the safety factor (q) (from top to bottom). The original CREATE-NL data are shown in blue, while their fitted approximations are depicted in orange. To ensure initial MHD stability, the T_e and FF' profiles were adjusted so that q could be maintained above unity in the core. These modified curves (green lines), together with the fitted n_e profile (which has not been modified), are adopted as input for recalculating the equilibrium inside JOREK (Section 3.1.3).

The modifications of the equilibrium profiles are methodologically consistent with the strategy employed for the Day-0 scenario. As detailed in Section 3.4.3, the first step of the ATQ induces a massive thermal collapse that reduces the core temperature by several orders of magnitude, effectively erasing the specific features of the pre-disruption temperature profile. Furthermore, the subsequent current flattening step reshapes the core q -profile regardless of its initial state. Consequently, utilizing this modified MHD-stable equilibrium allows for a robust initialization of the simulations without compromising the validity of the results during the CQ phase.

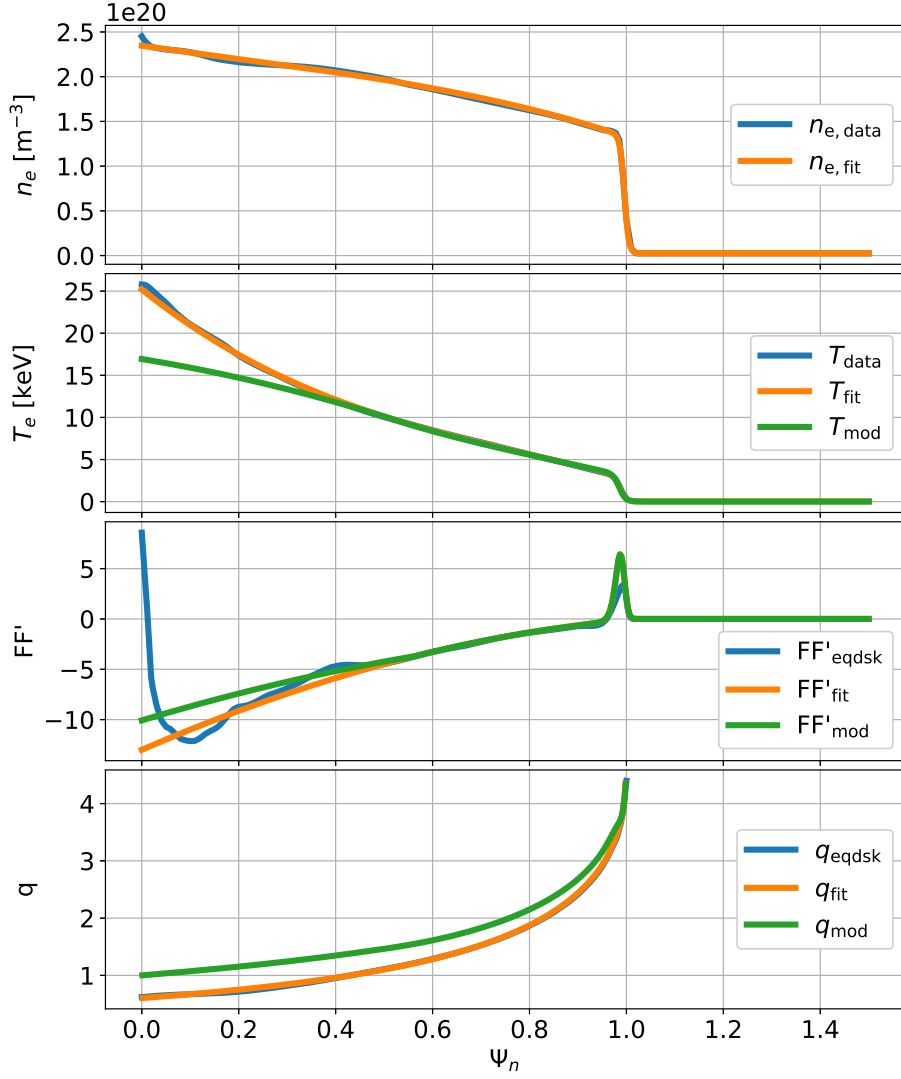


Fig. 5.1 Comparison of the equilibrium profiles for the Full Power scenario as a function of the normalized poloidal flux ψ_n . From top to bottom: electron density n_e , electron temperature T_e , FF' , and safety factor q . The blue curves represent the original CREATE-NL data, while the orange curves show the fitted approximations. While the electron density profile was left unmodified from the fit, the profiles for T_e , FF' and q were modified to eliminate pre-existing core instabilities. These modified, MHD-stable profiles ($q > 1$), shown in green, are adopted as the input for the Full Power simulations.

5.1.2 Physical parameters

The numerical parameters governing the plasma and RE transport follow the framework detailed in Chapter 3. Consistent with the Day-0 phase, the temperature-dependent coefficients ($\eta_e, \chi_{\parallel}$) are regularized by the same cut-off values, and a constant particle diffusivity $D = 2 \text{ m}^2/\text{s}$ is applied to the thermal species. The RE fluid retains the anisotropic diffusion tensor ($D_{\parallel, \text{RE}} \gg D_{\perp, \text{RE}}$) introduced previously. Notably, while the exact same normalized input parameter is used for the dominant parallel component ($D_{\parallel, \text{RE}}$) as in the Day-0 simulations, the higher initial central density of the FP scenario results in an adjusted physical value of $1 \cdot 10^9 \text{ m}^2/\text{s}$, compared to $1.83 \cdot 10^9 \text{ m}^2/\text{s}$ for the Day-0 scenario (see Appendix B for the details regarding JOREK normalization). As discussed in previous chapters, this order of magnitude aligns with optimal ranges established in the literature to properly model fast parallel dynamics without inducing spurious perpendicular diffusion [112, 113]. As with the Day-0 case, conducting dedicated sensitivity studies on $D_{\parallel, \text{RE}}$ to further characterize RE loss rates within the FP scenario could be explored in future work. The main coefficients used for the FP simulations are summarized in Table 5.1.

5.1.3 RE seed and impurity scan

The results presented in this chapter build upon a broad scan of the initial RE seed current and the quantity of injected impurities, in line with the strategy described in Section 3.4.2.

Regarding the RE seed, we account for the large uncertainties in primary generation mechanisms by treating I_{seed} as an independent variable spanning several orders of magnitude, as done for the Day-0 scenario. To maintain a consistent relative framework, the cases are defined as follows with respect to specific fractions of the initial plasma current ($I_p = 5.5 \text{ MA}$):

- **Baseline seed scenario – 5.5 A** ($10^{-6} I_p$): This represents a minimal reference seed, testing the system's susceptibility to avalanche from negligible initial populations.
- **Intermediate seed scenario – 550 A** ($10^{-4} I_p$): A transitional case representing moderate RE generation during the TQ.

Coeff.		Dependency	Value
η_e	$T_e < 500 \text{ eV}$	Spitzer $\propto T_e^{-3/2}$	$\eta_e = \eta_{e,\text{Sp}}$
	$T_e \geq 500 \text{ eV}$	Constant	$\eta_e^{\min} = \eta_{e,\text{Sp}} (@500 \text{ eV})$
$\chi_{\parallel}$	$T_e < 20 \text{ eV}$	Constant	$\chi_{\parallel}^{\min} = \chi_{\parallel,\text{SH}} (@20 \text{ eV})$
	$20 < T_e < 350 \text{ eV}$	Spitzer-Härm $\propto T_e^{5/2}$	$\chi_{\parallel} = \chi_{\parallel,\text{SH}}$
	$T_e \geq 350 \text{ eV}$	Constant	$\chi_{\parallel}^{\max} = \chi_{\parallel,\text{SH}} (@350 \text{ eV})$
$\chi_{\perp}$		Radial tanh profile	$\chi_{\perp,\text{core}} = 2 \text{ m}^2/\text{s}$
D		Constant	$2 \text{ m}^2/\text{s}$
$D_{\parallel,\text{RE}}$		Constant	$1 \cdot 10^9 \text{ m}^2/\text{s}$
$D_{\perp,\text{RE}}$		Constant	$1 \cdot 10^{-2} \text{ m}^2/\text{s}$

Table 5.1 Summary of the physical coefficients employed for the Full Power scenario. The resistivity (η_e) and parallel thermal diffusivity ($\chi_{\parallel}$) are based on standard Spitzer and Spitzer-Härm scalings, respectively, saturated at specific temperature thresholds. A radial hyperbolic tangent profile is used for the perpendicular thermal diffusivity ($\chi_{\perp}$), which drops to zero around $\psi_n = 0.92$ over a width of 0.01. The thermal particle diffusivity (D) is constant and isotropic. For the runaway electron fluid, a highly anisotropic diffusion tensor is employed. The large value for parallel diffusivity ($D_{\parallel,\text{RE}}$) was selected based on extensive benchmarks of the RE fluid model, as it provides the optimal balance between capturing fast parallel transport and minimizing numerical errors in the perpendicular direction [112, 113]. The specific physical value corresponds to the exact same normalized input used for the Day-0 scenario, adjusted here by the higher density normalization factor. The perpendicular component ($D_{\perp,\text{RE}}$) is kept at the minimum value required for numerical stability [113].

- **High-amplitude seed scenario – 55 kA** ($10^{-2}I_p$): The upper-bound scenario for this parametric scan. Similar to the Day-0 case, increasing the seed significantly beyond this value would shift the problem domain from avalanche dynamics to primary generation physics.
- **Comparison case – 20 kA**: This specific value is explicitly included to facilitate a direct comparison with the high-amplitude seed used for Day-0 simulations (Chapter 4), isolating the effect of the higher pre-disruption plasma current.

Simultaneously, the impurity content (N_{imp}) is varied from $1 \cdot 10^{20}$ to $6 \cdot 10^{21}$ particles. While the DMS for DTT is currently in the design phase and formalized reference parameters are still under development, this chosen range serves as a

relevant baseline for ongoing predictive studies. Specifically, the upper bound provides a scenario consistent with preliminary design estimates: the injection of a 12 mm mixed pellet with a 20% neon and 80% deuterium fraction (containing roughly 10^{23} neon particles) assuming a pre-TQ assimilation efficiency of less than 10%.

5.2 Results

This section details the simulation results for the DTT FP scenario ($I_p = 5.5$ MA). Following the ATQ procedure, the plasma enters the self-consistent CQ phase where the evolution of the electron temperature is determined by the balance between Ohmic heating (which is driven exclusively by the thermal current) and impurity radiation. We present the results organized by the magnitude of the initial RE seed, culminating in a direct comparison with the Day-0 scenario.

5.2.1 Baseline RE seed (5.5 A)

The lowest RE seed current considered for FP, $I_{\text{seed}} = 5.5$ A (a fraction of 10^{-6} of the initial plasma current), represents a scenario where primary generation mechanisms are relatively weak. Figure 5.2 illustrates the time evolution of the plasma and RE currents for varying impurity levels.

The simulations illustrate the substantially higher susceptibility of the FP scenario to avalanche amplification compared to the Day-0 case, consistent with the theoretical avalanche expectations discussed in Section 2.3.3. At low impurity levels ($N_{\text{imp}} \approx 1 \cdot 10^{20}$, represented by the black curves), radiative cooling is insufficient; consequently, T_e remains elevated, maintaining the electric field below the threshold required for significant multiplication, albeit at the cost of a partially mitigated disruption. Conversely, for all the other impurity inventories considered ($N_{\text{imp}} \gtrsim 1 \cdot 10^{21}$), the enhanced radiation clamps the post-TQ temperature to approximately 5–15 eV. The resulting surge in resistivity drives a stronger electric field, amplifying the baseline 5.5 A seed by over five orders of magnitude ($G_{\text{av}} \approx 1.3 \cdot 10^5$) and ultimately culminating in a macroscopic RE beam of roughly 0.7 MA. Notably, this growth exhibits distinct temporal characteristics: due to the microscopic initial seed, the RE current remains negligible during the first ≈ 2 ms of the CQ even at the highest

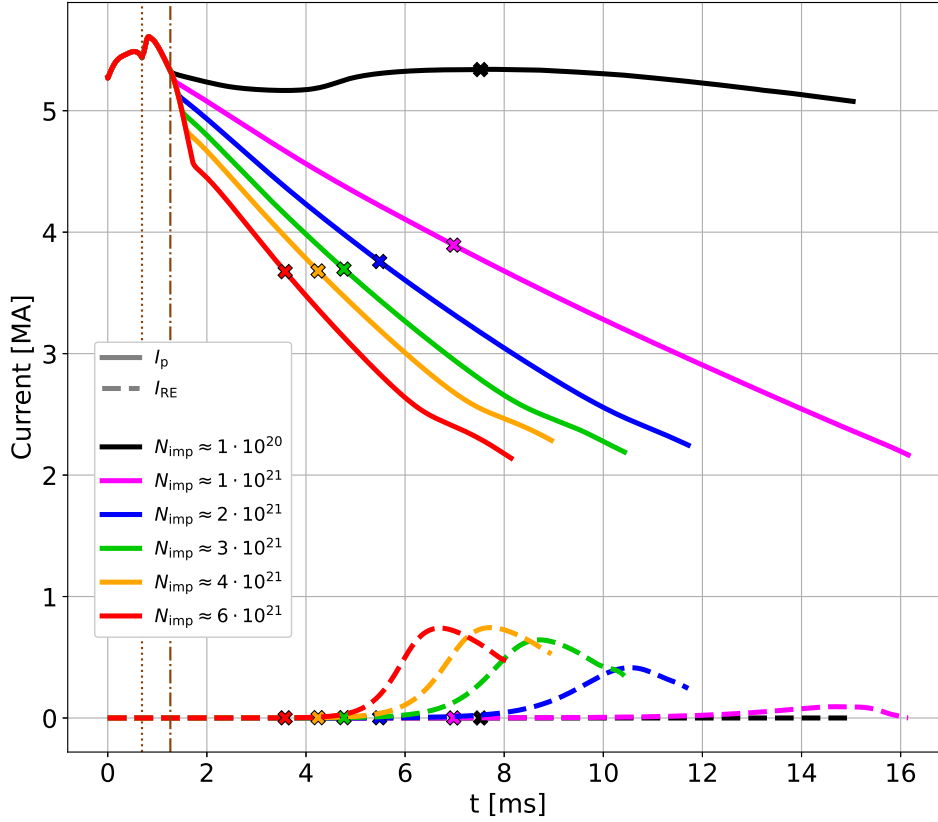


Fig. 5.2 Temporal evolution of the total plasma current (I_p , solid lines) and runaway electron current (I_{RE} , dashed lines) for the baseline initial seed case of 5.5 A. The colors distinguish the specific neon impurity inventory (N_{imp}), spanning a range from $1 \cdot 10^{20}$ to $6 \cdot 10^{21}$ particles. The phases of the artificial thermal quench (ATQ) are delineated by vertical brown lines: the dotted line marks the end of the initial thermal collapse, while the dash-dotted line indicates the start of impurity injection. The termination of the final ATQ phase (impurity injection) is discernible by the distinct change in the I_p slope immediately following the dash-dotted line; this transition marks the onset of the self-consistent current quench, which is progressively delayed at higher impurity levels because the fixed injection rate requires a longer duration to reach the target inventory. The 'x' markers denote the timestamps of first plasma contact with the plasma-facing components. A key finding is the formation of macroscopic RE beams (≈ 0.7 MA) in the high-impurity regimes ($N_{imp} \gtrsim 3 \cdot 10^{21}$, green, orange, and red traces). This massive amplification, occurring despite both PFC contact and the minimal baseline initial seed, serves as a clear demonstration of the high avalanche gain characteristic of the Full Power scenario.

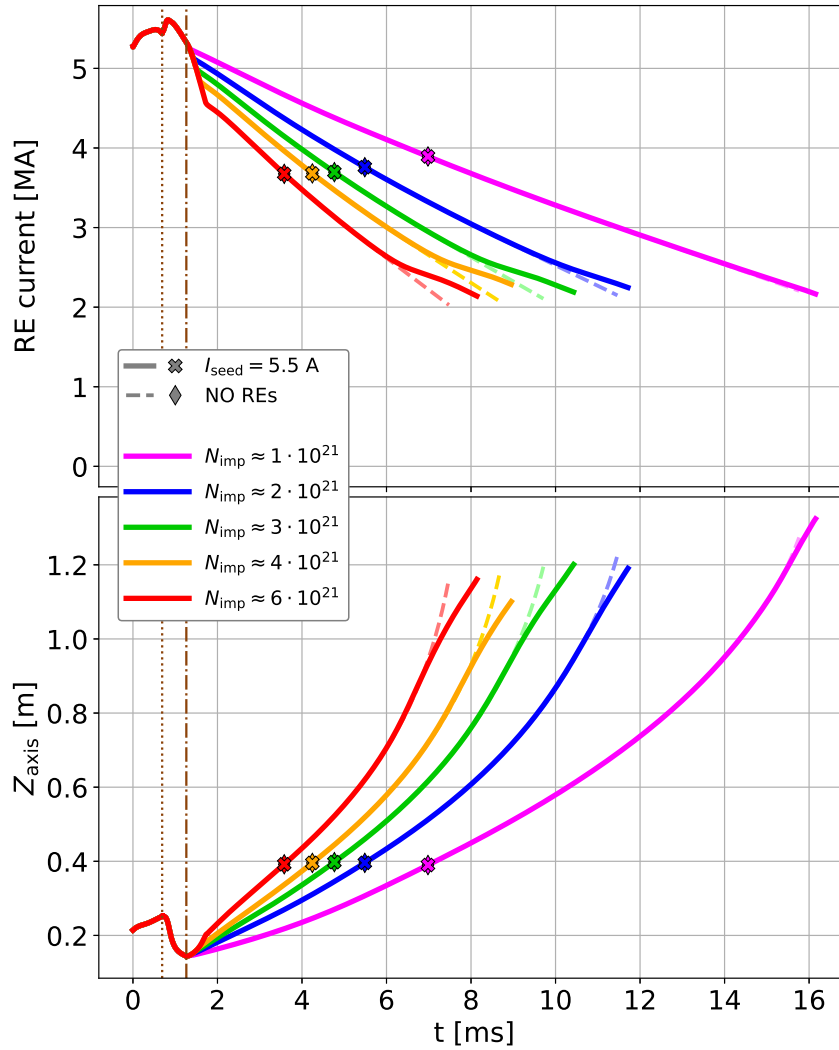


Fig. 5.3 Simultaneous evolution of the total plasma current (I_p , top panel) and the vertical position of the magnetic axis (Z_{axis} , bottom panel). The simulations initialized with the baseline initial seed ($I_{seed} = 5.5$ A) are represented by thick solid lines, whereas the thin dashed lines denote reference cases in which the avalanche source term was numerically suppressed. As in the preceding figure, color-coding indicates the impurity inventory, and vertical brown lines delineate the artificial thermal quench transitions. The 'x' and diamond markers denote the timestamps of first plasma contact with the plasma-facing components (PFCs) for the simulations with and without runaway electrons (REs), respectively. The bottom panel highlights that higher impurity injections drive more rapid vertical displacement events (yielding steeper slopes), thereby advancing the onset of PFC contact. Notably, a delayed dynamical impact of the RE population is observed: due to the relatively small initial seed, the avalanche requires several milliseconds to amplify the current to macroscopic levels sufficient to alter the current quench dynamics. Consequently, the inductive braking effect on the vertical motion, discernible as the divergence between the solid and dashed curves, manifests exclusively during the late stages of the disruption.

impurity concentrations, evolving into a macroscopically significant component of the total current only during the final stages of the disruption.

The impact of the RE beam on the macroscopic vertical stability of the plasma is shown in Figure 5.3. Because the RE population takes several milliseconds to grow large enough to influence the discharge dynamics, the vertical trajectory of the plasma (solid lines) only deviates from the RE-free reference cases (dashed lines) during the final phase of the disruption, thereby manifesting a braking effect on the vertical motion. Crucially, the plasma makes first contact with the PFCs while the avalanche process is still highly active. Consequently, the ongoing RE generation is constrained by the physical scraping-off of the plasma against the wall, which progressively shrinks the confined volume in a manner highly analogous to the dynamics observed in the Day-0 scenario. This competition establishes a temporary equilibrium between avalanche generation and physical losses to the wall, manifesting as a characteristic plateau in the RE current traces, before the RE current ultimately enters its final decay phase, driven by the continuous upward displacement of the plasma and the progressive reduction of its confined volume.

5.2.2 Intermediate RE seed (550 A)

Increasing the initial seed scenario to 550 A ($\approx 10^{-4} I_p$) yields qualitatively similar trends but results in notably higher beam currents and accelerated timescales. As displayed in Figure 5.4, this larger initial population effectively lowers the impurity threshold required for beam formation: significant avalanche multiplication is now observed at intermediate impurity levels that previously yielded negligible macroscopic RE generation in the baseline 5.5 A case. This advanced growth enables the RE beam to achieve macroscopic magnitudes much earlier in the CQ. For the high impurity case of $N_{\text{imp}} \approx 4 \cdot 10^{21}$, the RE current peaks at approximately 1.8 MA, constituting nearly 60% of the instantaneous plasma current.

To clarify the internal distribution of the plasma current, Figure 5.5 provides a detailed current decomposition for this specific simulation. The analysis reveals that the current within the closed flux surfaces (I_{core}) is heavily dominated by the RE population; indeed, the narrow gap between the I_{core} and I_{RE} curves confirms that the thermal current pool in the core has been almost entirely converted into runaways. Conversely, the difference between the total plasma current and the core current

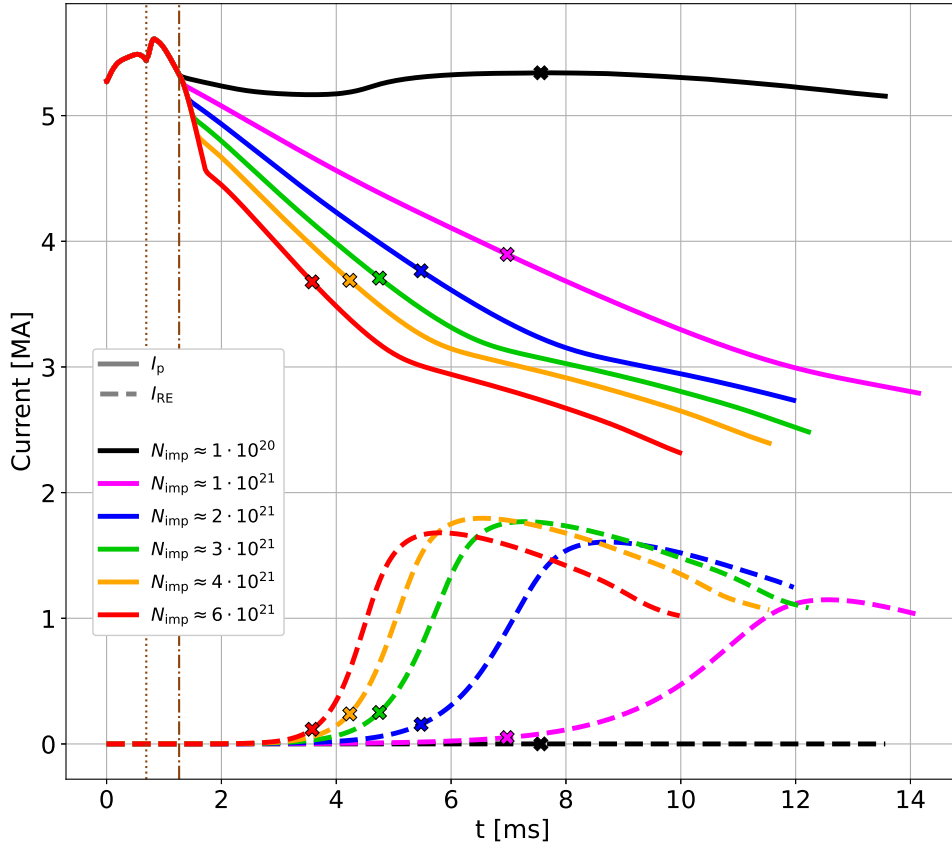


Fig. 5.4 Temporal evolution of the total plasma current (I_p , solid lines) and runaway electron (RE) current (I_{RE} , dashed lines) for the intermediate initial seed scenario ($I_{seed} = 550$ A). Color-coding and the vertical brown markers delineating the artificial thermal quench phases follow the convention established in previous figures, while the 'x' markers denote the timestamps of first plasma contact with the plasma-facing components. A comparison with the baseline seed scenario highlights a substantial reduction in the impurity threshold required for beam formation, with macroscopic RE currents emerging at significantly lower impurity inventories. Furthermore, for intermediate impurity levels, the magnitude of the resulting RE beam is notably higher, culminating in a peak RE current of approximately 1.8 MA, constituting nearly 60% of the instantaneous plasma current.

$(I_p - I_{core})$ represents resistive halo currents flowing along the open field lines of the cold scrape-off layer (SOL). As indicated by the separation between the solid (I_p) and dotted (I_{core}) traces, these halo currents play a major role, carrying roughly 1 MA during the late CQ phase.

Finally, it is key to clarify the mechanisms governing the temporary plateau and subsequent decay of the RE current. Much like in the baseline seed scenario, the initial contact with the PFCs occurs when the RE current is still relatively

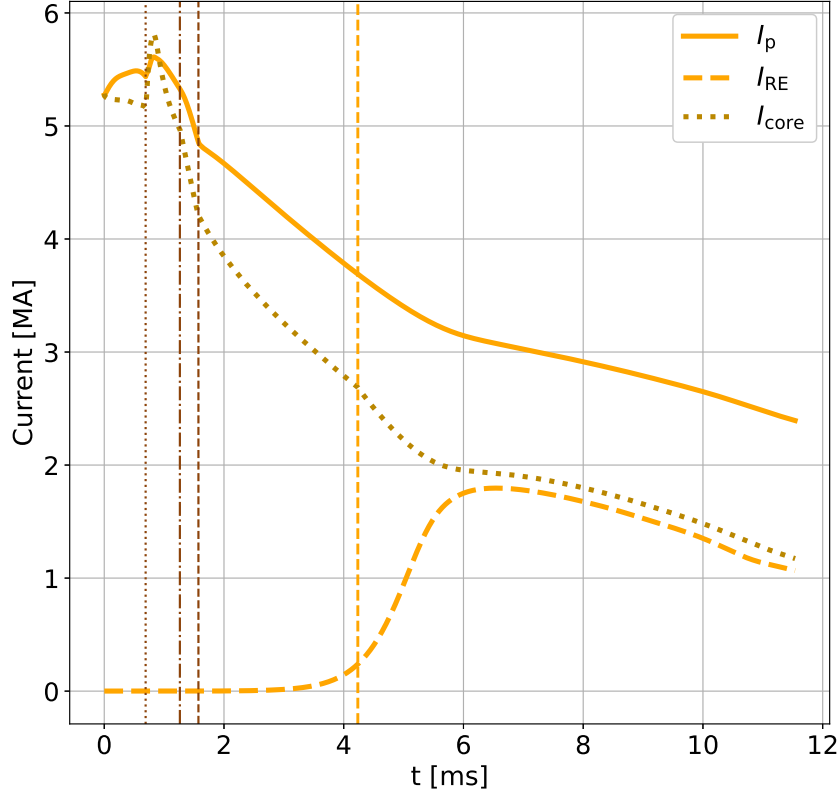


Fig. 5.5 Detailed decomposition of the total plasma current components for the intermediate seed case ($I_{\text{seed}} = 550$ A) with an impurity inventory of $N_{\text{imp}} \approx 4 \cdot 10^{21}$ (corresponding to the orange traces in Figure 5.4). The plot displays the total plasma current (I_p , solid line), the runaway electron current (I_{RE} , dashed line), and the total current flowing within the closed flux surfaces (I_{core} , gold dotted line). The vertical brown lines follow the artificial thermal quench conventions established in preceding figures, with an additional brown dashed vertical line at $t \approx 1.6$ ms indicating the completion of the ATQ procedure and the onset of the self-consistent current quench. An orange dashed vertical line at $t \approx 4$ ms marks the timestamp of first contact with the plasma-facing components (PFCs). Notably, during the artificial current flattening phase (between the dotted and dash-dotted vertical lines), a transient I_p spike occurs alongside the development of a negative edge skin current. A key observation is the convergence of the I_{core} and I_{RE} curves after $t \approx 6$ ms; this overlap demonstrates that the current within the confined plasma volume is carried almost entirely by the RE beam despite continuous wall-scraping against the PFCs, implying that the remaining current ($I_p - I_{\text{RE}}$) consists predominantly of halo currents flowing in the scrape-off layer.

low. Crucially, this wall interaction does not immediately terminate the avalanche mechanism: within the remaining closed flux surfaces, the avalanche continues to drive robust growth until the RE current virtually consumes the entire core current pool ($I_{\text{RE}} \approx I_{\text{core}}$). This explains the trend observed at the highest impurity inventories (e.g., between $4 \cdot 10^{21}$ and $6 \cdot 10^{21}$ particles), where a stronger RE avalanche results in a lower absolute peak RE current. While the avalanche successfully converts virtually all of the available core electrons to runaways in both cases, the higher impurity inventory accelerates the decay rate of the underlying core current. Consequently, the avalanche operates on a much more rapidly depleting current reservoir, bounding the absolute peak to lower values. As highlighted in Figure 5.5, it is only after this full core conversion is achieved that the continuous upward VDE and the associated scraping-off of the closed flux surfaces against the first wall become the dominant macro-dynamics, ultimately forcing the RE beam into its final decay phase.

5.2.3 High-amplitude RE seed (55 kA)

The upper bound of the parameter scan, $I_{\text{seed}} = 55 \text{ kA}$ ($= 10^{-2} I_p$), simulates a scenario characterized by exceptionally strong primary generation. As shown in Figure 5.6, the RE beam forms almost immediately after the CQ onset, and the maximum RE current reaches approximately 3.2 MA. Crucially, the simulation with the lowest amount of impurities ($1 \cdot 10^{20}$ particles, black line) eventually decays without avalanching, proving that a sufficient impurity inventory is a strict prerequisite for RE growth even when initialized with massive seeds. This absence of growth is due to the limited radiative cooling provided by this minimal impurity concentration: the post-thermal quench electron temperature remains sufficiently high to prevent the induced electric field from significantly exceeding the critical threshold for avalanche multiplication. Nevertheless, this inevitably results in a partially mitigated disruption characterized by severe thermal and halo current loads.

In contrast, for impurity levels of $N_{\text{imp}} \gtrsim 1 \cdot 10^{21}$, the conversion efficiency is remarkably high, and the RE current rapidly constitutes the vast majority ($\approx 80\%$) of the total plasma current. The presence of such a massive current carried by the runaway population has a profound effect on the CQ dynamics and on the vertical displacement of the plasma column, as seen in Figure 5.7. The macroscopic RE beam exerts a powerful braking effect that significantly slows down the upward vertical motion, thereby noticeably prolonging the current quench duration. Interestingly, a

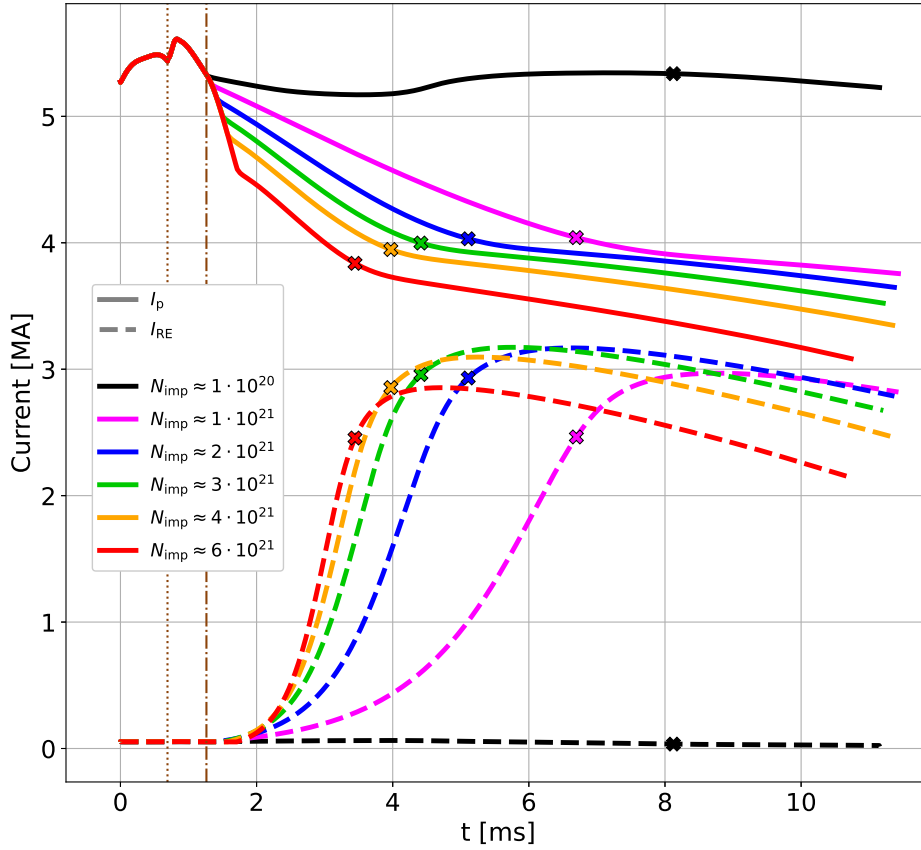


Fig. 5.6 Temporal evolution of the total plasma current (I_p , solid curves) and runaway electron (RE) current (I_{RE} , dashed curves) for the high-amplitude initial seed case of 55 kA. Color-coding and the vertical brown markers delineating the artificial thermal quench phases follow the conventions established in preceding figures. The 'x' markers denote the timestamps of first plasma contact with the plasma-facing components (PFCs). Under these high-amplitude seed conditions, the avalanche dynamics are exceptionally efficient: the RE fraction rapidly escalates to approximately 80% of the total current within a few milliseconds, even for moderate impurity inventories ($N_{imp} \approx 1 \cdot 10^{21}$). The resulting beams reach peak values of ≈ 3.2 MA, a magnitude sufficient to induce a prominent prolongation of the current quench by significantly reducing the total plasma current decay rate.

comparison with the RE-free reference simulations (dashed lines) reveals that in the absence of runaways, the plasma column touches the physical walls later and at a higher vertical position (Z_{axis}). This indicates that without the current-preserving effect of the RE beam, the rapid resistive collapse of the current drives a more severe shrinkage of the confined plasma volume, whereas the RE beam sustains the current and stalls this rapid volumetric contraction during the VDE.

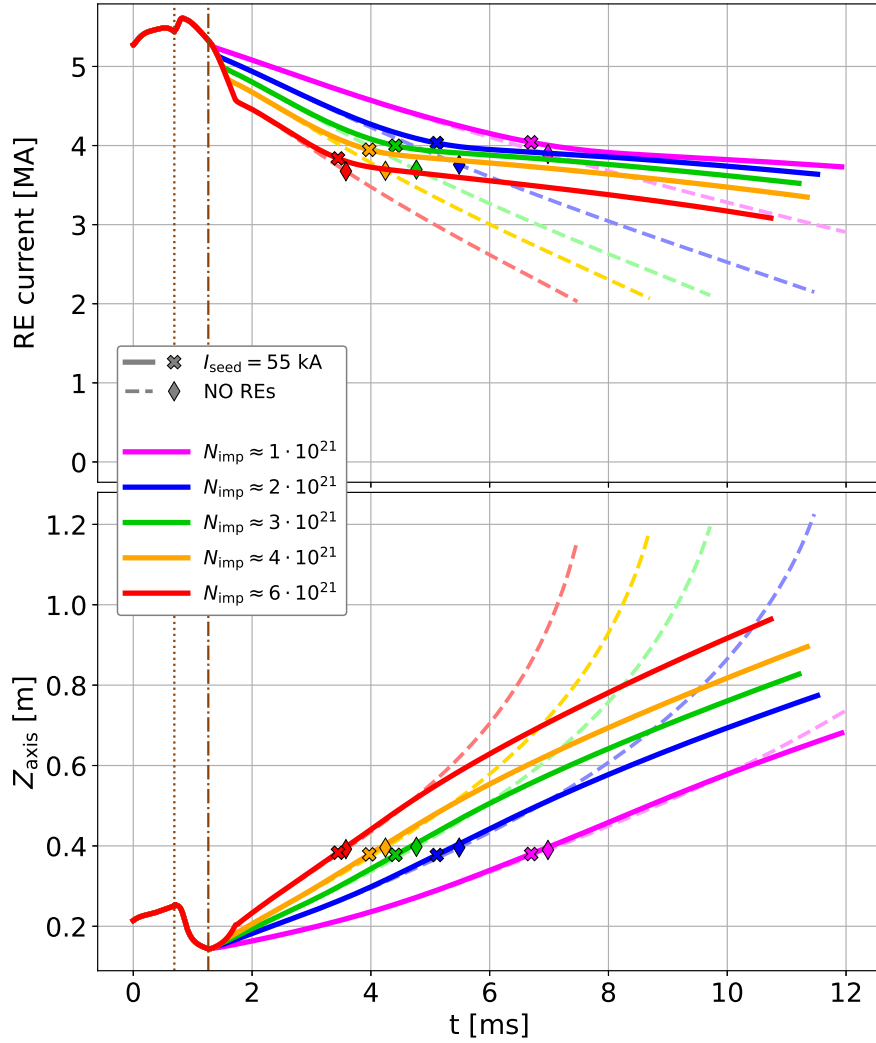


Fig. 5.7 Comparison of the total plasma current (I_p , top panel) and the vertical position of the magnetic axis (Z_{axis} , bottom panel). Thick solid lines represent the high-amplitude initial seed scenario ($I_{seed} = 55$ kA), whereas thin dashed lines denote reference simulations in which the avalanche mechanism was numerically suppressed. Color-coding and the vertical brown markers delineating the artificial thermal quench phases follow the conventions established in preceding figures. The 'x' and diamond markers denote the timestamps of first plasma contact with the plasma-facing components (PFCs) for the simulations with and without runaway electrons (REs), respectively. In this high-amplitude seed scenario, the RE beam carries a dominant fraction of the current, thereby significantly mitigating both the vertical displacement velocity and the current quench decay rate compared to the RE-free reference cases.

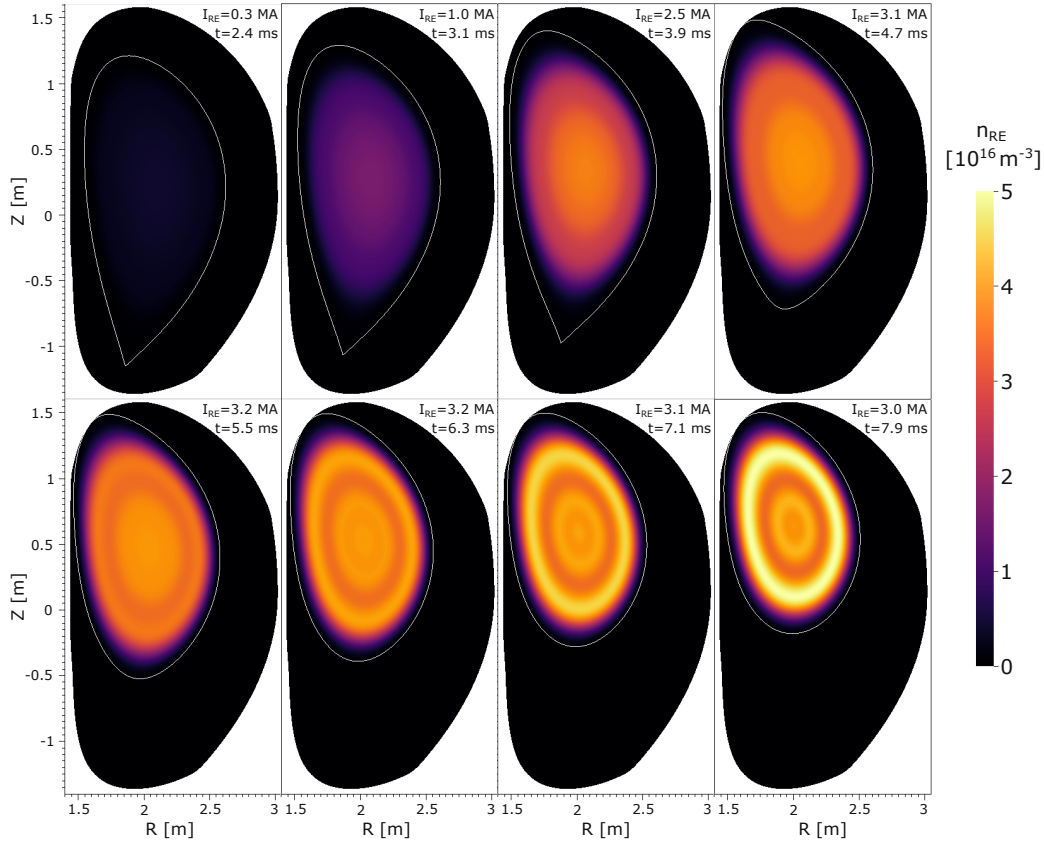


Fig. 5.8 Evolution of the runaway electron density profile (n_{RE} , expressed in 10^{16} m^{-3}) for the high-amplitude initial seed scenario ($I_{\text{seed}} = 55 \text{ kA}$) with an impurity inventory of $N_{\text{imp}} \approx 3 \cdot 10^{21}$ (corresponding to the green traces in Figures 5.6 and 5.7). The white contour denotes the separatrix, distinguishing the confined plasma from the open field line region. The sequence illustrates the wall-scraping mechanism: as the vertical instability drives the plasma upward, the confined volume is progressively reduced by interaction with the first wall. The onset of this physical contact with the plasma-facing components (PFCs) corresponds precisely to the green 'x' markers in Figures 5.6 and 5.7. Although the local RE density in the core continues to rise due to the avalanche mechanism, the total RE current (I_{RE} , indicated at the top right of each panel) eventually decays due to this geometric reduction of the confined plasma volume. Notably, the absolute peak of I_{RE} for this scenario occurs at $t \approx 5.8 \text{ ms}$, falling between the discrete snapshots displayed here.

Consistent with the mechanisms described for the intermediate seed scenario, a similar trend is observed at the highest impurity inventories ($N_{\text{imp}} \gtrsim 3 \cdot 10^{21}$), where the absolute peak of the RE current diminishes with increasing impurity content. In these regimes, the avalanche is so efficient that the RE current virtually consumes the entire core current pool ($I_{\text{RE}} \approx I_{\text{core}}$) well before the plasma column makes first contact with the PFCs. Because a higher impurity concentration accelerates

the underlying decay rate of the core current reservoir, the avalanche operates on a more rapidly depleting target, bounding the absolute peak RE current to lower values despite the stronger initial multiplication. Following this full core conversion into RE current, the slow decay of I_p is primarily driven by the scraping-off of the closed flux surfaces of plasma against the PFCs rather than resistive dissipation. Figure 5.8 visualizes this termination mechanism through the evolution of the RE density poloidal cross-section. In the early phase, the beam is generated in the core and remains well-confined. As the vertical instability progresses and drives the plasma column upward, the last closed flux surface (represented by the white contour) progressively shrinks due to the interaction with the upper first wall. This continuous scraping-off of the outer layers causes the loss of electrons to the wall (which is extremely rapid for REs) to dominate over avalanche mechanisms within the remaining confined region, governing the final decay phase of the plasma current.

5.3 Discussion and general assessment of RE risks in DTT

The simulations presented in the preceding chapters provide a comprehensive assessment of the behavior of REs in DTT. By utilizing the non-linear MHD code JOREK coupled with the resistive wall code STARWALL, we have explored two distinct operational phases: the initial Day-0 commissioning phase at a plasma current $I_p = 2$ MA, and the Full Power (FP) operation at $I_p = 5.5$ MA. The methodological approach, which relies on an artificial thermal quench (ATQ), allows for a focused analysis of the post-TQ dynamics. By varying the initial RE seed current (I_{seed}) and the impurity inventory (N_{imp}) over several orders of magnitude, we have mapped the operational boundaries of macroscopic RE beam formation. This section aims to provide a detailed interpretation of these results, synthesizing the physical mechanisms that differentiate the two scenarios and evaluating the implications for disruption mitigation strategies in DTT.

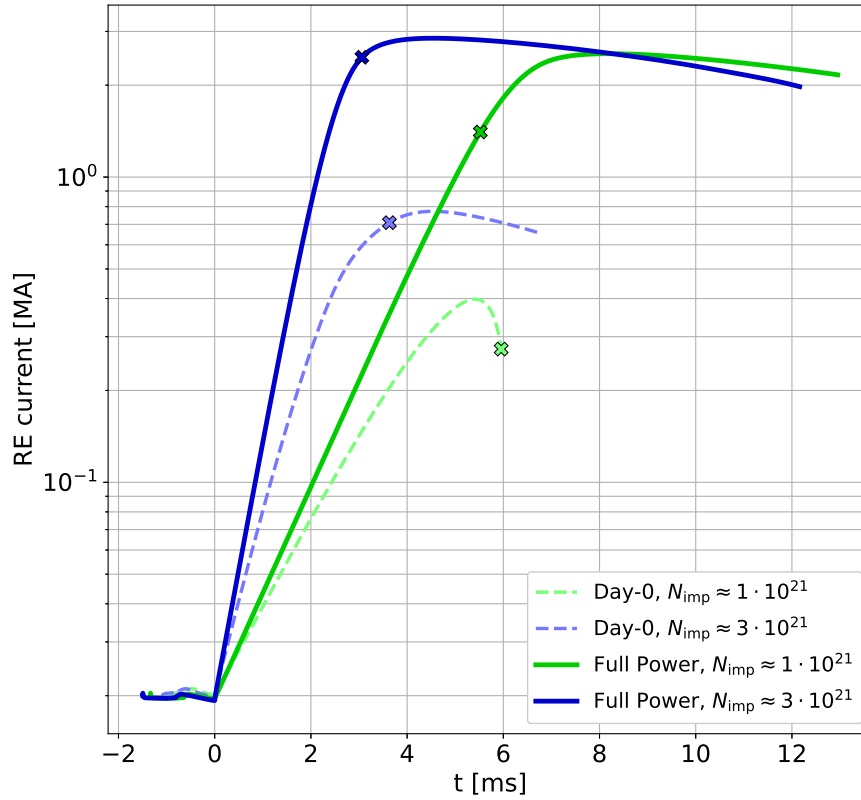


Fig. 5.9 Comparative analysis of runaway electron current evolution for the Day-0 scenario ($I_p = 2$ MA, dashed lines, adapted from Figure 4.5) versus the Full Power scenario ($I_p = 5.5$ MA, solid lines). Both configurations are initialized with an identical initial seed of 20 kA to isolate the effect of the plasma current on the RE multiplication factor. The green and blue traces correspond to impurity inventories of $N_{\text{imp}} \approx 1 \cdot 10^{21}$ and $3 \cdot 10^{21}$ particles, respectively. The 'x' markers denote the timestamps of first plasma contact with the plasma-facing components (PFCs). The use of a semi-logarithmic scale underscores the divergence in avalanche dynamics: the Full Power cases exhibit visibly steeper growth rates and higher peak currents compared to the Day-0 benchmarks, confirming the exponential sensitivity of the avalanche gain to the background plasma current magnitude.

5.3.1 Comparison: Day-0 vs. Full Power

To quantify the difference in RE risk between the commissioning phase and full-power operations, we compare simulations with identical initial seeds ($I_{\text{seed}} = 20$ kA) and impurity levels. Figure 5.9 presents this comparison on a semi-logarithmic scale. The Day-0 scenario (dashed lines) exhibits a comparatively RE-safe response, constrained by a two-fold effect of the lower initial plasma current. First, the lower I_p intrinsically limits the absolute reservoir of thermal current that can be converted into runaways.

As established in the previous chapters, even in high-amplitude seed cases where the core current becomes virtually entirely carried by REs, the absolute peak is strictly bounded by this smaller initial pool. Second, the lower plasma current reduces the avalanche multiplication factor. In the FP scenario (solid lines), both of these constraints are lifted: the 5.5 MA plasma provides a vastly larger absolute pool of convertible current alongside an enhanced avalanche gain. Visually, this is captured by the steeper slope of the solid curves (corresponding to a much faster avalanche growth rate) and their much higher peak values compared to the shallower rise and lower ceilings of the Day-0 lines. This dual effect leads to peak currents that are nearly an order of magnitude higher than in the Day-0 case for $N_{\text{imp}} \approx 1 \cdot 10^{21}$, reaching ≈ 3 MA. This direct comparison confirms that the safety margins inherent to the 2 MA phase vanish at 5.5 MA, necessitating robust mitigation strategies.

5.3.2 The role of avalanche multiplication

The main result emerging from the comparison of the two scenarios is the fundamental difference in their susceptibility to the RE avalanche phenomenon. While the Day-0 scenario appears robust against RE beam formation even for relatively high RE seeds, the FP scenario exhibits a drastically narrowed RE safety margin. As expected from the theoretical framework established in Section 2.3.3, this contrast confirms that the discharge dynamics are primarily governed by the exponential sensitivity of the avalanche mechanism to the higher plasma current.

The simulation results are consistent with the theoretical expectations based on the Rosenbluth-Putvinski avalanche scaling introduced in Equation (2.36). In the Day-0 case ($I_p = 2$ MA), the theoretical gain was estimated to be relatively modest ($G_{\text{av}} \sim 10^2$). JOREK simulations confirm this: even when initialized with substantial seeds (up to 200 A) and driven by strong electric fields induced by massive impurity injection, the avalanche multiplication is insufficient to bridge the gap to macroscopic RE currents. It is only when the upper bound of the seed scan (20 kA) is explored that a macroscopic RE beam forms. In contrast, the FP scenario ($I_p = 5.5$ MA) possesses a reservoir of poloidal magnetic flux sufficient to drive massive amplification. The analytical estimate of $G_{\text{av}} \approx 10^6$ presented in Section 2.3.3 is recovered by the simulations, where a microscopic seed of 5.5 A ($10^{-6} I_p$) is converted into a 0.7 MA beam ($G_{\text{av}} \approx 1.3 \cdot 10^5$) solely through avalanche multiplication. When the seed is increased by a couple of orders of magnitude (550 A), the FP scenario effectively

converts the vast majority of the confined thermal current into RE current. This implies that for FP operations, the uncertainty regarding the initial RE seed size becomes secondary: the avalanche gain is so high that even minimal RE seeds are sufficient to trigger a macroscopic RE beam. This observed weak dependence of the final RE current on the initial seed magnitude is highly consistent with the self-regulating electric field dynamics introduced in Chapter 2 [89].

Beyond the initial growth rates, the results also highlight fundamentally different saturation behaviors between the two scenarios. In the low-current Day-0 simulations, the avalanche gain is moderate. The RE beam reaches a maximum either because it successfully replaces the vast majority of the confined plasma current (thereby exhausting the driving electric field) or due to plasma volume shrinkage as it comes into contact with PFCs or moves towards the X-point. In the high-current FP simulations, however, the dynamics are distinct. For most of the combinations in our parametric scan, the avalanche is so potent that the RE current rapidly grows until it fully replaces the thermal plasma current within the closed flux surfaces ($I_{\text{RE}} \approx I_{\text{core}}$), even after making contact with the PFCs. Once this occurs, the core resistive effects vanish, and the confined plasma becomes highly conductive. The temporary RE plateau and subsequent decay are imposed by a geometric competition: the avalanche converting the electrons in the core into REs versus the VDE scraping off the confined plasma volume against the first wall.

5.3.3 Vertical stability and threat assessment

A consistent feature across both operational scenarios is the strong coupling between the RE beam evolution and the vertical displacement of the plasma. The conversion of the thermal core into a highly conductive RE beam essentially stalls the resistive current decay, sustaining I_p at MA levels for extended durations. During the subsequent upward VDE characteristic of the FP scenario, the plasma column is physically scraped off by the PFCs. While this continuous volume reduction forces the termination of the beam, the dynamic equilibrium established during this wall-scraping phase poses a severe threat to the machine. Instead of harmless dissipation, the energy of the macroscopic RE beam (whose current exceeds 3 MA in the worst-case FP scenarios analyzed here) is deposited directly onto localized areas of the first wall. This highlights the absolute necessity of avoiding this highly

efficient avalanche regime entirely, or otherwise ensuring the implementation of sacrificial limiters capable of withstanding these concentrated thermal loads.

5.3.4 Impact of impurities

The quantity of injected impurities (N_{imp}) is controlled experimentally according to the disruption mitigation strategy, introducing a complex trade-off between thermal load mitigation and RE avoidance. The parameter scans performed in this work allow us to categorize the operational space of DTT into distinct regimes based on impurity content.

At low impurity levels ($N_{\text{imp}} \lesssim 1 \cdot 10^{20}$), the radiative cooling is insufficient to sustain the low electron temperatures required for a high-resistivity plasma. In this regime, the CQ is slow, the induced electric field remains relatively weak, and macroscopic RE generation is suppressed regardless of the scenario or seed size. While this may eliminate the RE risk, it is unlikely to be a viable strategy for disruption mitigation. The resulting incomplete thermal collapse leads to a partially mitigated disruption, leaving the PFCs vulnerable to severe, unmitigated thermal loads and subjecting the machine structures to extreme halo current forces. Therefore, the management of these specific loads necessitates separate, targeted investigations.

An intermediate “transition” regime appears for moderate impurity levels ($1 \cdot 10^{20} < N_{\text{imp}} \lesssim 1 \cdot 10^{21}$). Here, the RE safety of the discharge is conditional. In the FP scenario, the electric field is strong enough to drive an avalanche, but the final macroscopic outcome depends heavily on the initial RE seed size. A small seed may not have enough time to multiply to dangerous levels before the plasma strikes the wall, while a larger primary seed will result in a significant beam. This regime presents a substantial challenge for predictive modeling, as the safety of the machine relies heavily on the accurate estimation of primary generation mechanisms (e.g., hot-tail or Dreicer generation), which are currently subject to large uncertainties.

Finally, for high impurity contents ($N_{\text{imp}} \gtrsim 2 \cdot 10^{21}$), the simulations indicate a transition to a regime fundamentally dominated by REs. In the FP scenario, this level of impurity injection leads invariably to the formation of a macroscopic RE beam, largely independent of the initial seed size. The enhanced resistivity drives a faster CQ that maximizes the induced voltage, while the presence of partially ionized high-Z impurities may further boost the avalanche growth rate via close collisions with

bound electrons. It is crucial to emphasize that the Day-0 scenario remains largely immune to this regime. Even at the highest impurity densities modeled (approaching the background ion density), the 2 MA plasma does not support the massive runaway growth seen at 5.5 MA, unless the initial seed is already exceptionally large ($10^{-2}I_p$). This confirms that the RE risk associated with massive gas injection is not a primary concern for the commissioning phase.

5.3.5 Strategies for Full Power operation

The FP scenario demands a much more nuanced operational approach compared to the Day-0 phase. The simulations indicate that relying on pure neon injection for thermal mitigation carries an unacceptable risk of triggering macroscopic RE beams. Consequently, future strategies must navigate a narrow operational window. On one hand, restricting the high-Z atom inventory below the critical avalanche threshold ($N_{\text{imp}} \approx 1\text{--}2 \cdot 10^{21}$) is required to prevent macroscopic RE beam formation. On the other hand, injecting pure impurity quantities below this threshold risks an incomplete thermal collapse. Such a partially mitigated disruption leaves the PFCs vulnerable to severe, unmitigated thermal loads and subjects the machine structures to extreme halo current forces. To decouple these inherently contradictory requirements, a two-step SPI strategy may be considered. This approach utilizes deuterium [53] or trace-neon SPI mixtures to ensure necessary core penetration while suppressing runaway generation [134, 135]. Ultimately, given the extreme strength of the avalanche source term identified for the FP scenario, preventing the formation of the RE beam directly at its source becomes of paramount importance for the Full Power scenario of DTT.

Chapter 6

Conclusions and outlook

6.1 Summary of main achievements

The DTT facility plays a pivotal role in the European fusion roadmap, serving as the bridge between ITER and EU-DEMO regarding power exhaust solutions. However, the high operational parameters required to reproduce reactor-relevant conditions inevitably introduce the risk of severe disruptions and runaway electron (RE) generation. This thesis has provided the first systematic predictive assessment of RE avalanche dynamics for DTT, focusing on two critical milestones: the commissioning phase (Day-0) and the Full Power scenario.

By employing the non-linear MHD code JOREK coupled with the resistive wall code STARWALL, we established a robust simulation framework. The use of an artificial thermal quench (ATQ) allowed us to bypass the complexities of the thermal quench (TQ) onset, isolating the physics of the current quench (CQ) and the subsequent RE avalanche. Through extensive parametric scans of initial RE seed currents and impurity injection amounts, this work has characterized the operational space of DTT and identified the primary constraints for RE-safe machine operation.

The central finding of this research is the quantitative definition of the impurity thresholds that separate benign operational windows from regimes dominated by macroscopic RE beams. The analysis reveals that the viability of disruption mitigation strategies is fundamentally determined by the machine's sensitivity to the impurity content:

- **Day-0 Scenario** ($I_p = 2$ MA): This phase is characterized by a weak dependence on the scanning parameters. The limited reservoir of poloidal magnetic flux restricts the avalanche potential, rendering the system inherently robust against RE avalanches. Our simulations showed that even under pessimistic assumptions that combine massive impurity injection (up to $N_{\text{imp}} \approx 3 \cdot 10^{21}$) with worst-case initial seeds, the resulting RE currents peak at levels below 1 MA. This effectively decouples the thermal mitigation problem from RE avoidance, indicating that the commissioning phase allows for a broader range of impurity injection parameters to optimize thermal protection with a significantly reduced risk of triggering macroscopic RE beams compared to the Full Power scenario.
- **Full Power Scenario** ($I_p = 5.5$ MA): In the Full Power regime, the operational window narrows drastically. The simulations identified a critical impurity threshold at $N_{\text{imp}} \approx 1 \cdot 10^{21}$. Below this value, the electric field is insufficient to drive a sustained RE avalanche, but the resulting incomplete thermal collapse leaves the PFCs vulnerable to severe, unmitigated thermal and halo current loads. Conversely, above this threshold, the high avalanche gain allows even seed populations as low as $10^{-6} I_p$ to be amplified into macroscopic beams. This creates a critical contradiction, where the massive high-Z injection required for thermal mitigation acts as the primary driver for macroscopic RE beam formation. Consequently, the Full Power scenario necessitates a disruption mitigation strategy that strictly limits high-Z injection in favor of alternative mixtures (e.g., deuterium-rich) to safely decouple thermal mitigation from RE generation.

6.2 Implications for DTT operations

The results presented herein have direct consequences for the ongoing design of the DTT disruption mitigation system (DMS). As previously discussed, the DTT-DMS will serve both as a machine protection system and an experimental platform to test ITER-relevant mitigation strategies [10]. While the massive gas injection (MGI) of pure high-Z impurities to mitigate thermal loads is viable for the Day-0 commissioning phase, the strategy must evolve for Full Power operations. The simulations suggest that relying on massive quantities of pure high-Z noble gases to

dissipate thermal energy will almost certainly trigger an RE event. Therefore, the operational window for safe disruption mitigation at 5.5 MA is narrow: the DMS (which will eventually transition to an SPI-based system) must inject enough material to mitigate thermal loads and electromagnetic forces, but must limit the quantity of high-Z atoms to prevent the generation of macroscopic RE beams.

These are intrinsically contradictory requirements, as the injection of low-Z atoms (such as hydrogen isotopes) does not significantly contribute to the radiation required to mitigate the thermal loads. To decouple these needs, a two-step shattered pellet injection (SPI) strategy was initially envisaged for the ITER DMS [53]. In this approach, a hydrogen-isotope SPI is injected first to drastically increase the core electron density (aiming to suppress RE generation rather than increase radiation), followed by a second injection with an impurity (e.g., neon) SPI to radiate out the thermal energy.

However, recent simulations indicate that this technique is limited by the strong outward $\mathbf{E} \times \mathbf{B}$ plasmoid drifts experienced by hydrogen-isotope SPI, which severely limit the core assimilation of such isotopes and, consequently, compromise the primary goal of the first injection [54]. To overcome this, an alternative strategy currently under investigation is the use of trace neon SPI to replace the pure hydrogen-isotope SPI: by adding a small fraction of neon (typically less than 1%), the outward plasmoid drift could be effectively suppressed, thereby improving core assimilation [134, 135].

In light of these recent findings, the strategy for the DTT DMS should consider these advanced injection schemes, which may be essential to successfully navigate the narrow operational window, ensuring sufficient core density to avoid REs while allowing for the safe dissipation of thermal loads.

6.3 Future outlook

While this work lays the foundation for predictive RE modeling in DTT, further research is required to translate these physical insights into a fully optimized RE avoidance and mitigation strategy. Based on the limitations and findings of this work, the following activities are proposed:

- **Sensitivity analyses and model refinements:** The results presented in this work rely on specific parameterizations and modeling assumptions that warrant

further exploration. Most notably, the parallel diffusivity of the RE fluid ($D_{\parallel, \text{RE}}$) plays a critical role in determining the RE loss rates upon physical contact with the PFCs; dedicated sensitivity studies tailored to the DTT geometry will help fully characterize its impact on the results. Furthermore, future efforts should aim to relax certain physical simplifications, such as transitioning from the simple coronal equilibrium assumption to a full charge state distribution. Finally, expanding the simulation parameter space to encompass more extreme impurity inventories and a broader continuous scan of initial RE seeds will help solidify and refine the operational boundaries identified in this study.

- **Self-consistent TQ modeling:** The ATQ method, while efficient for scanning parameters, enforces an axisymmetric thermal collapse. Future work will aim to simulate realistic TQs triggered by SPI using 3D modeling. This will provide physically self-consistent initial conditions for the CQ.
- **3D effects and stochastic losses:** The simulations presented in this work were toroidally axisymmetric, neglecting 3D MHD instabilities that may destroy magnetic flux surfaces and create stochastic field lines. While 3D magnetic perturbations introduce significant RE loss mechanisms through enhanced radial transport (potentially expanding the safety window), the resulting dynamics are highly complex. Recent studies indicate that these stochastic effects do not strictly guarantee a reduction in the final RE beam current [136]. Therefore, future 3D simulations are required to accurately quantify the interplay between these phenomena and to determine the net impact of magnetic perturbations on the RE safety limits of DTT.
- **Advanced kinetic modeling:** While the RE fluid approximation is robust for macroscopic dynamics, it cannot capture velocity-space phenomena such as the precise energy distribution of the RE beam or orbit-dependent losses. Future studies could leverage the newly developed full-f, full-orbit RE modules in JOREK [58] or hybrid fluid-kinetic schemes. These approaches will be particularly valuable for refining the estimates of the heat deposition on PFCs, where finite-orbit effects play a crucial role.
- **Optimization of gas mixtures:** The present study used pure neon to isolate the impurity dependence. Future simulations should investigate mixtures of deuterium and neon or other impurities (e.g., argon). The addition of deuterium

could provide the necessary electron density for preventing the avalanche of a macroscopic RE beam, potentially widening the safe operational window for the Full Power scenario.

- **Injection techniques for RE seed suppression:** Because of the strong avalanche gain expected in reactor-scale devices, the primary disruption mitigation goal is shifting toward the suppression of the initial runaway seed. Building upon the advanced injection schemes discussed above, future research for DTT should aim to develop and simulate dedicated injection techniques specifically designed to minimize RE seed generation. Future 3D modeling should assess how effectively different injection timings, shatter geometries, and novel material mixtures can penetrate the core and limit primary RE generation.
- **RE beam termination and sacrificial limiters:** The interaction of the RE beam with the first wall is the ultimate threat to the device. Detailed 3D simulations of the termination phase are needed to assess the spatial distribution of the heat loads during the vertical displacement. These studies will be crucial for engineering solutions, such as the strategic placement of sacrificial limiters designed to intercept the beam and protect the primary first wall structures.
- **Integrated disruption analysis:** Finally, while this work focused strictly on RE dynamics, a complete disruption safety assessment requires a multi-physics approach. Future analyses must integrate RE risks with the evaluation of thermal loads on PFCs during the TQ and the computation of electromagnetic forces on the vacuum vessel, ensuring a holistic protection strategy for DTT.

Appendix A

Finite element method (FEM)

The finite element method (FEM) is the core numerical technique used by JOREK to solve the MHD equations in the poloidal plane. The complex 2D poloidal cross-section requires a flexible discretization capable of handling realistic tokamak geometries. This appendix summarizes the mathematical foundation of the method, transitioning from the physical strong form to the numerical algebraic system. A rigorous derivation of the method and its theoretical foundations can be found in [137].

A.1 Weak formulation

The starting point is a general partial differential equation for a physical variable X , which can be written in terms of a differential operator $\mathcal{L}$ and a source term f :

$$\mathcal{L}X = f \tag{A.1}$$

Solving this *strong form* directly can be numerically difficult, as it requires the solution to strictly satisfy the equation at every point in the domain and possess continuous derivatives up to the second order (C^2 continuous), assuming $\mathcal{L}$ contains second-order derivatives. To overcome this, FEM employs the *weak form* of the equation, where it is multiplied by an arbitrary test function v_i^* and integrated over

the domain Ω . For our starting partial differential equation, the weak form is:

$$\int_{\Omega} v_i^* \mathcal{L}X = \int_{\Omega} v_i^* f \quad (\text{A.2})$$

A key advantage of this formulation arises when integration by parts is applied to the left-hand side: the operation shifts spatial derivatives from the unknown solution X to the test function v_i^* , relaxing the differentiability requirements on X . For instance, second-order derivatives in the strong form may only require first-order continuity (a C^1 continuous function) in the weak form.

A.2 Discretization and the Galerkin method

The second key step is the discretization of the spatial domain into a grid of elements with different shapes (such as triangles or, in the case of JOREK, cubic Bézier elements). The continuous variable X is then expanded as a linear combination of basis functions w_j^* :

$$X \approx \sum_{j=1}^N x_j w_j^*, \quad (\text{A.3})$$

where x_j are the unknown time-dependent coefficients and N is the total number of basis functions in the grid. The coefficients x_j are found such that the residual $\mathcal{R}$ is set to zero:

$$\mathcal{R} \equiv \int_{\Omega} v_i^* (\mathcal{L}X - f) = \int_{\Omega} v_i^* \left(\sum_{j=1}^N x_j \mathcal{L}w_j^* - f \right) = \sum_{j=1}^N x_j \int_{\Omega} v_i^* \mathcal{L}w_j^* - \int_{\Omega} v_i^* f = 0 \quad (\text{A.4})$$

Once the basis functions w_j^* and the test functions v_i^* are chosen such that the integrals can be computed analytically, we can write a linear system of equations that has the form:

$$\underline{M}\mathbf{x} = \mathbf{f} \quad (\text{A.5})$$

where $\mathbf{x}$ is the vector of unknown coefficients, and the components of $\underline{M}$ and $\mathbf{f}$ are defined as:

$$M_{ij} \equiv \int_{\Omega} v_i^* \mathcal{L}w_j^*, \quad f_i \equiv \int_{\Omega} v_i^* f$$

JOREK uses a specific approach referred to as the Ritz-Galerkin method, in which the test functions and basis functions originate from the same function space and are a specific class of cubic polynomials [96].

Appendix B

JOREK normalization

The physical model in JOREK relies on a set of normalized MHD equations, described in [96]. This normalization is performed primarily with respect to the central mass density ρ_0 at the magnetic axis (corresponding to a central number density n_0) and to the vacuum permeability μ_0 . This choice ensures that μ_0 is absorbed into the variables and does not appear explicitly in the final system of equations, simplifying the mathematical structure.

A distinctive aspect of the normalization applied in JOREK is that length scales are not normalized and remain in their physical SI units. Time, instead, is normalized by a factor τ_{norm} , defined as $\tau_{\text{norm}} = \sqrt{\mu_0 \rho_0}$. This normalization time is physically comparable to the standard Alfvén time, $\tau_A = a\sqrt{\mu_0 \rho_0}/B$, although it excludes the minor radius of the tokamak, a , and the magnetic field, B .

Consistent with this framework, the total mass density ρ is normalized by the initial central density ρ_0 , while the total pressure p is normalized using μ_0 . Several physical constants and ratios are also used in the definition of JOREK variables: k_B denotes the Boltzmann constant, and γ represents the adiabatic index (typically $\gamma = 5/3$ for plasmas). Furthermore, the impurity particle density n_{imp} utilizes the mass ratio $\mu_{\text{imp}} = m_i/m_{\text{imp}}$.

To distinguish between physical and normalized quantities, variables bearing the subscript “SI” refer to values in the standard International System of Units. Conversely, variables appearing without this subscript are expressed in the normalized units internal to the JOREK code. A selection of the most relevant normalization factors, following the conventions in [96], is detailed in Table B.1 on the following page.

SI vs. normalized (JOEKE) units		Description of the quantity
$R_{\text{SI}}(\text{m})$	$= R$	Major radius
$Z_{\text{SI}}(\text{m})$	$= Z$	Vertical coordinate
$\mathbf{B}_{\text{SI}}(\text{T})$	$= \mathbf{B}$	Magnetic field vector
$\mathbf{E}_{\text{SI}}(\text{V m}^{-1})$	$= \mathbf{E} / \sqrt{\mu_0 \rho_0}$	Electric field vector
$\psi_{\text{SI}}(\text{T m}^2)$	$= \psi$	Poloidal magnetic flux
$j_{\phi, \text{SI}}(\text{A m}^{-2})$	$= -j / (R \mu_0)$	Toroidal current density; $j_{\phi, \text{SI}} = \mathbf{j}_{\text{SI}} \cdot \mathbf{e}_\phi$
$n_{\text{SI}}(\text{m}^{-3})$	$= \rho n_0$	Particle density
$\rho_{\text{SI}}(\text{kg m}^{-3})$	$= \rho \rho_0$	Mass density = ion mass $\times$ particle density
$n_{\text{RE}, \text{SI}}(\text{m}^{-3})$	$= n_{\text{RE}} \sqrt{\rho_0 / \mu_0} / (eR)$	Runaway electron number density
$n_{\text{imp}, \text{SI}}(\text{m}^{-3})$	$= n_0 \rho_{\text{imp}} \mu_{\text{imp}}$	Impurity number density
$T_{\text{SI}}(\text{K})$	$= T / (k_B \mu_0 n_0)$	Temperature = electron + ion temperature
$T_{\text{eV}}(\text{eV})$	$= T / (e \mu_0 n_0)$	Temperature = electron + ion temperature
$FF'_{\text{SI}}(\text{T rad})$	$= FF'$	Toroidal flux function $F \equiv RB_\phi$; $F' = dF/d\psi$
$p_{\text{SI}}(\text{N m}^{-2})$	$= \rho T / \mu_0$	Plasma pressure
$\mathbf{v}_{\text{SI}}(\text{m s}^{-1})$	$= \mathbf{v} / \sqrt{\mu_0 \rho_0}$	Velocity vector
$v_{\parallel, \text{SI}}(\text{m s}^{-1})$	$= v_{\parallel} \cdot B_{\text{SI}} / \sqrt{\mu_0 \rho_0}$	Parallel velocity component, where $B_{\text{SI}} = \mathbf{B}_{\text{SI}} $
$u_{\text{SI}}(\text{m s}^{-1})$	$= u / \sqrt{\mu_0 \rho_0}$	$F_0 u = \Phi$ is the electric potential
$\omega_{\phi, \text{SI}}(\text{m}^{-1} \text{s}^{-1})$	$= \omega / \sqrt{\mu_0 \rho_0}$	Toroidal vorticity
$t_{\text{SI}}(\text{s})$	$= t \cdot \sqrt{\mu_0 \rho_0}$	Time
$\eta_{e, \text{SI}}(\Omega \text{ m})$	$= \eta_e \cdot \sqrt{\mu_0 / \rho_0}$	Electrical resistivity
$\mu_{\text{SI}}(\text{kg m}^{-1} \text{s}^{-1})$	$= \mu \cdot \sqrt{\rho_0 / \mu_0}$	Dynamic viscosity
$\nu_{\text{SI}}(\text{m}^2 \text{s}^{-1})$	$= \mu_{\text{SI}} / \rho_{\text{SI}}$	Kinematic viscosity (ρ_{SI} is the local value)
$D_{\text{SI}}(\text{m}^2 \text{s}^{-1})$	$= D / \sqrt{\mu_0 \rho_0}$	Particle diffusivity ($\parallel$ or $\perp$)
$K_{\text{SI}}(\text{kg m}^{-1} \text{s}^{-1})$	$= K \cdot \frac{\sqrt{\rho_0 / \mu_0}}{\gamma - 1}$	Heat diffusivity ($\parallel$ or $\perp$)
$\chi_{\text{SI}}(\text{m}^2 \text{s}^{-1})$	$= K_{\text{SI}} / \rho_{\text{SI}}$	Heat diffusivity ($\parallel$ or $\perp$)
$q_{\text{SI}}(\text{A s})$	$= q \sqrt{\mu_0 \rho_0}$	Particle charge

Table B.1 Connection between SI and normalized units for various plasma parameters, adapted from [96]. Variables with the subscript ‘SI’ denote physical units, while those without are normalized. The normalization factors n_0 and ρ_0 are the central number and central mass densities, respectively, while μ_0 is the vacuum permeability. Additionally, k_B is the Boltzmann constant, γ is the adiabatic index ($\gamma = 5/3$ for plasmas), and $\mu_{\text{imp}} = m_i / m_{\text{imp}}$.

Nomenclature

ASDEX	Axially Symmetric D ivertor EX periment
ATQ	Artificial Thermal Q uench
BC	B oundary C ondition
BDF2	B ackward D ifferentiation F ormula of 2nd -order
CFL	C ourant- F riedrichs- L ewy
CQ	C urrent Q uench
CS	C entral S olenoid
DC	D irect C urrent
DMS	D isruption M itigation S ystem
DTT	D ivertor T okamak T est
ECRH	E lectron C yclotron R esonance H eating
EM	E lectrom m agnetic
EU-DEMO	E Uropean DEMO nstration Power Plant
FEM	F inite E lement M ethod
FP	F ull P ower
FW	F irst W all
GS	G rad- S hafranov

HF	Halo Fraction
ICF	Inertial Confinement Fusion
ICRH	Ion Cyclotron Resonance Heating
ITER	International Thermonuclear Experimental Reactor
JET	Joint European Torus
LCFS	Last Closed Flux Surface
MCF	Magnetic Confinement Fusion
MGI	Massive Gas Injection
MHD	MagnetoHydroDynamics
NBI	Neutral Beam Injection
PFC	Plasma Facing Component
PF	Poloidal Field
RE	Runaway Electron
RH	Remote Handling
SI	Système International d'unités (International System of Units)
SN	Single Null
SOL	Scrape-Off Layer
SPI	Shattered Pellet Injection
TF	Toroidal Field
TPF	Toroidal Peaking Factor
TQ	Thermal Quench
VDE	Vertical Displacement Event

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