Mechanisms of plasma disruption and runaway electron losses in tokamaks

S. S. Abdullaev, K.H. Finken², K. Wongrach², M. Tokar¹, H.R. Koslowski¹, O. Willi², L. Zeng³, and the TEXTOR team

¹ Forschungszentrum Jülich GmbH, Institut für Energie- und Klimaforschung - Plasmaphysik, D–52425 Jülich, Germany,
² Institut für Laser- und Plasmaphysik, Heinrich-Heine Universität Düsseldorf, Germany,
³ Institute of Plasma Physics, Chinese Academy of Sciences, 230031 Hefei, China

(Received 8 January 2015)

Based on the analysis of data from the numerous dedicated experiments on plasma disruptions in the TEXTOR tokamak mechanisms of the formation of runaway electron beams and their losses are proposed. The plasma disruption is caused by strong stochastic magnetic field formed due to nonlinearly excited low-mode number MHD modes. It is hypothesized that the runaway electron beam is formed in the central plasma region confined inside the intact magnetic surface located between \( q = 1 \) and the closest low–order rational [\( q = 4/3 \) or \( q = 3/2 \)] magnetic surfaces. The thermal quench time caused by the fast electron transport in a stochastic magnetic field is calculated using the collisional transport model. The current decay stage is due to the ambipolar particle transport in a stochastic magnetic field. The runaway electron beam in the confined plasma region is formed due to their acceleration in the inductive toroidal electric field. The runaway electron beam current is modeled as a sum of toroidally symmetric part and a small amplitude helical current with a predominant \( m/n = 1/1 \) component. The runaway electrons are lost due to two effects: (i) by outward drift of electrons in a toroidal electric field until they touch wall and (ii) by the formation of stochastic layer of runaway electrons at the beam edge. Such a stochastic layer for high–energy runaway electrons is formed in the presence of the \( m/n = 1/1 \) MHD mode. It has a mixed topological structure with a stochastic region open to wall. The effect of external resonant magnetic perturbations on runaway electron loss is discussed. A possible cause of the sudden MHD signals accompanied by runaway electron bursts is explained by the redistribution of runaway current during the resonant interaction of high–energetic electron orbits with the \( m/n = 1/1 \) MHD mode.

1. Introduction

One of the severe consequences of the plasma disruptions in tokamaks is the generations of the runaway electron (RE) beams (Wesson et al. 1989; Gill 1993; Schüller 1995; Gill et al. 2000, 2002) (see, also (Wesson 2004; Boozer 2012)). The REs generated during the disruptions of tokamak plasmas may reach several tens of MeV and may contribute to the significant part of post–disruption plasma current. The prevention of such RE beams is of a paramount importance in future tokamaks, especially in the ITER operation, since it may severely damage a device wall (Bécoulet et al. 2013).

At present there are several proposals to mitigate REs generated during plasma disruptions. The mitigation of REs by the gas injections has been discussed (see, e.g., Refs. (Hender et al. 2007; White et al. 2002, 2003; Bakhtiari et al. 2002, 2005; Granetz et al. 2007; Bozhenkov et al. 2008; Pautasso et al. 2009; Lehnen et al. 2009; Hollmann et al. 2010; Reux et al. 2010; Lehnen et al. 2011)). Suppression of REs by the resonant magnetic...
perturbations (RMPs) has been also intensively discussed since late 1990s (see, e.g., Refs. \cite{Kawanoetal1997, Tokuda&Yoshino1999, Helanderetal2000, Yoshino&Tokuda2000, Lehnenetal2008, Lehnenetal2009, Hollmannetal2010, Pappetal2011, Pappetal2012}). However, up to now there is no a regular strategy to solve this problem. One of the reasons is that the physical mechanisms of the formation of REs during plasma disruptions is still not well-known. The different scenarios of runaway formation during plasma disruptions are discussed in literature. Particularly, in Refs. \cite{Fulopetal2009, Fulop&Newton2014} the possible roles of whistler waves on the generation of REs and Alfvénic wave instabilities driven by REs have been discussed.

There were the numerous dedicated experiments to study the problem of runaway current generation during plasma disruptions triggered by massive gas injections (MGI) in the TEXTOR tokamak (see, e.g., \cite{Forsteretal2012, Zengetal2013, Wongrachetal2014}), in KSTAR tokamak \cite{Chenetal2013}, the JET tokamak \cite{Plyusninetal2006, Lehnenetal2011}, in DIII-D \cite{Hollmannetal2010, Conmauxetal2011, Hollmannetal2013}, Alcator C-Mod \cite{Olynyketal2013}, and others. In these works the dependencies of RE generation on the toroidal magnetic field, on the magnetic field fluctuations, on the species of injection gases have been investigated. Particularly, in KSTAR tokamak \cite{Chenetal2013} it has been found that there is no the toroidal magnetic field threshold $B_T < 2$ T as was indicated by previous experiments in other tokamaks. In Ref. \cite{Izzoetal2011, Izzoetal2012} MHD simulations have been performed to study the confinement REs generated during rapid disruptions by MGI in DIII-D, Alcator C-Mod, and ITER. Such simulations with two different MHD codes have been carried out by \cite{Izzoetal2012} to analyze shot–to–shot variability of RE currents in DIII-D tokamak discharges.

These numerous experiments show the complex nature of plasma disruption processes especially the formation of RE beams, and its evolution. One of the important features of this event is its irregularity and variability of RE beam parameters from one discharge to another one. This indicates the sensitivity of disruption processes and RE beam formations on initial conditions which is the characteristic feature of turbulent processes, particularly, the deterministic chaotic system. Therefore, \textit{ab initio} numerical simulations of these processes may be not always successful to understand their mechanisms because of the complexity of computer simulations of turbulent processes \cite{Kadanoff2004}. The problems of numerical simulations of plasma disruptions is comprehensively discussed by \cite{Boozer2012}.

In this work we intend to approach to this problem from the point of view of Hamiltonian chaotic systems, mainly the magnetic stochasticity in a magnetically confined plasmas \cite{Abdullaev2014}. Based on the ideas of these systems and analyses of numerous experimental results, mainly obtained in the TEXTOR tokamak we propose possible mechanisms of formation and evolution of RE beams created during plasma disruptions. Since a self–consistent theoretical treatment of all these processes is very complicated we developed theoretical models for each stages of a plasma disruption. These models are used to estimate the characteristic times of the thermal and current quenches, the spatial size of runaway plasma beam and their decay times, the speed of RE radial drifts, the effect of magnetic perturbations.

It is believed that the plasma disruption starts due to a large–scale magnetic stochasticity caused by excited of MHD modes with low poloidal $m$ and toroidal $n$ numbers, $(m/n = 1/1, 2/1, 3/2, 5/2, \ldots)$. The heat and particle transports in the strongly chaotic magnetic field causes the fast temperature drop and ceases the plasma current. However, at the certain spectrum of magnetic perturbations, for example, at the sufficiently small amplitude of the $m/n = 1/1$ mode the chaotic field lines may not extend to the central
plasma region due to an intact magnetic surface located between magnetic surface \( q = 1 \) and the nearest low–order rational surface \( q = 4/3 \) [or \( q = 3/2 \)]. This intact magnetic surface confines particles in the central plasma region and serves as a transport barrier to particles during the current quench. Electrons in the confined region are accelerated due to large toroidal electric field and forms the relatively stable of RE beams. 

This occurs, for instance, when the plasma disruption initiated by the heavy Argon gas injection which does not penetrate deep into the plasma, therefore it does not excite the \( m/n = 1/1 \) mode with the sufficiently large amplitude. In contrary, the injection of the lighter noble gases neon and helium does not generate runaways. The reason is that light gases penetrate deeper into the plasma and excite the large–amplitude \( (m/n = 1/1) \) mode.

The existence of an intact magnetic surface and its location depends on the radial profile of the safety factor and the spectrum of magnetic perturbations. The latter sensitively depend on the plasma disruption conditions and vary unpredictable from one discharge to another during plasma disruptions. This makes RE formation process unpredictable and may explain a shot–to–shot variability of the parameters of RE beams.

Based on this mechanism we study the main three stages of the post–disruption plasma evolution: the fast thermal quench, the current decay stage (current quench), and the RE beam evolution. The physical processes during each of these stages will be studied by theoretical models. These processes are the formation of stochastic magnetic field, heat and particle transport in a stochastic magnetic field, the acceleration of electrons by inductive electric field, the lost mechanisms of REs, and the effect of internal and external magnetic perturbations.

The paper consists of eight sections. Mathematical tools and models employed to study the problems are given in Supplementary part. The numerous data obtained during the dedicated experiments in the TEXTOR tokamak are analyzed in Sec. 2. Possible mechanisms of plasma disruptions with a RE beam formation is proposed and analyzed in Sec. 3. The transport of heat and particles during the fast phase and the current decay stages of plasma disruption are studied in Sec. 4. The model of a post–disruption plasma beam is proposed in Sec. 5. Using this model a time–evolution of guiding–center (GC) orbits of electrons accelerating by the inductive toroidal electric field is studied in Sec. 6. Particularly, the change of RE confinement conditions with decreasing the plasma current and increasing the electron energy and the outward drift of GC orbits are investigated. The effect of external and internal magnetic perturbations on the RE confinement are discussed in Sec. 7. In the final Sec. 8 the we give the summary of obtained results and discuss their consequences.

2. Description of plasma disruptions

The TEXTOR is a middle size limiter tokamak with the major radius \( R_0 = 1.75 \) m, the minor radius \( a = 0.46 \) cm. The toroidal field \( B_0 \) can be varied up to 2.8 T, and the plasma current take up to 600 kA. In the experiments the plasma disruptions were triggered in a controlled way by gas injections using a fast disruption mitigation valve (DMV) [Bozhenkov et al. 2007, Finken et al. 2008, 2011, Bozhenkov et al. 2011]. Particularly, the disruptions with REs were triggered by argon (Ar) injection. The runaway-free disruptions were triggered either by helium or neon (He/Ne) injection performed by the smaller valve. The effect of the externally applied RMPs on the REs generations has been investigated using the dynamic ergodic divertor (DED) installed in the TEXTOR tokamak.

Figure 1(a) illustrates typical disruptions of the discharges of the TEXTOR tokamak.
with and without RE generations. Specifically, it shows the time evolution of plasma parameters (the loop voltage $V_{\text{loop}}$, the electron cyclotron emission (ECE), the soft X-ray (SXR) signal, the Mirnov signal, and scintillation probe signals) during disruptions of the discharges with REs (#119978) and without (#117444) REs.

There are also some discharges with untypical RE currents and shorter current decay times. The two examples of such discharges are shown in Fig. 1(b). We will discuss some features of these discharges at the end of the section.

The typical behavior of the plasma during the disruptions is following. The gas (Ar or Ne/He) was injected at the time instant $t = 2$ s. One can distinguish three stages of the disruption with the REs: the first (or fast) stage in which a sudden temperature drop occurs, in the second stage the plasma current starts to decay with a higher rate, and in the third stage the current decay slows down and the current beam with the REs is formed.

The first fast stage starts after a few milliseconds (between 2 ms and 5 ms) after the gas injection and ends with a sudden temperature drop (a thermal quench) in a time interval about one ms as seen from the ECE signals shown in a detail in Fig. 1 (c). The Mirnov signals indicating magnetic activities start just before of this time interval and they last a few milliseconds until a significant decay of the plasma current for the RE–free discharges or establishing the current with the REs (see Fig. 1 (a)). The magnetic activities are shown in Fig. 1 (b) for the two discharges with RE generations (#117507 and #119978) and the discharge without REs (#117444). One can notice that the growth of magnetic fluctuations after the sudden temperature drop.

The second stage of the plasma disruption begins with the current decay within a millisecond after the thermal quench. Particularly, for the discharges #117434 and #117444 the current decay starts in $0.47 \times 10^{-3}$ s and $0.87 \times 10^{-3}$ s, respectively, after the temperature drop (see Fig. 1 (a)). The time dependence of the current in this stage for all discharges is well approximated by the linear function of time $I_p = I_{p0} - bt$, where the coefficient $b = -dI_p/dt$ determines the average current decay rate. The scheme of determination of $b = [dI_p/dt]$ is shown in Fig. 2.

The values of the current decay rate $|dI_p/dt|$ during the current quench and RE plateau regimes, the initial RE current $I^{(RE)}_p$ for a number of discharges are listed Table 1. It also shows the time $t_{\text{max}}$ when the applied RMPs, i.e., the DED current $I_{\text{ded}}$ reaches its maximal value, and the toroidal mode $n$ of the RMPs.

At the beginning the current decay rate $|dI_p/dt|$ for all discharges are the same order and lies between $3.8 \times 10^4$ and $5.0 \times 10^4$ [kA/s] as listed in the 2-nd column of Table 1. For the discharges without REs the current completely disappears in a few millisecond. In this stage the loop voltage starts to rise due to inductive electric field opposing to the current decay.

In the third stage (RE plateau) of the disruption the rapid current decay is replaced by it’s slow decay and it starts the formation of the REs due to the acceleration of electrons in the inductive toroidal electric field and the secondary generation of REs. The values of the current decay rate $|dI_p/dt|$ along with the initial values of the plasma current $I^{(RE)}_p$ in this stage for several discharges are listed in the 3-rd and the 4-th columns of Table 1. The average values of $|dI_p/dt|$ for almost all discharges are confined in the interval (2.2, 5.6) MA/s, i.e., in one order lower than the current decay rate in the second stage. The values of $I^{(RE)}_p$ are also confined between 177 kA and 240 kA. These values of $|dI_p/dt|$ and $I^{(RE)}_p$ are close to the ones observed in the similar experiments in the DIII-D tokamak (see, e.g., Hollmann et al. 2010, 2013).

One should also note that in the RE plateau stage at certain time instants one observes
Mechanisms of plasma disruption and runaway electron losses in tokamaks

Figure 1. (a) Time evolution of the disruption of the TEXTOR shots #119978 (black solid lines) and #117444 (red curves) (from top to bottom): the plasma current, the loop voltage, the ECE signal, the SXR signal, the Mirnov signal, and scintillation probe signal. (b) The same but for the discharges #117859 (blue curves) and #120140 (black curves). (c) Initial stage of the temporal evolution of the plasma current (solid curve on the l.h.s. axis), ECE signals; (d) the Mirnov signals (all on r.h.s. axis) during a plasma disruption with (#117434 and #117507) and without (#117444) RE generations. \( I_{p}^{\text{RE}} \) is the initial value of the plasma current with REs. Disruptions for discharges #119978, #117507, #117859, #120140 are initiated by Ar injections, and #117444 by Ne injections.
Table 1. Parameters of discharges: 1-st column – the discharge number; 2-nd and 3-rd columns – the decay rates \(|dI_p/dt|\) [in MA/s] of the plasma current \(I_p(t)\) in the second and the third stages; 4-th column – the initial current of the RE beam \(I_p^{(RE)}\); 5-th column shows the parameters of the RMPs, a time \(t_{\text{max}}\) when the DED current reaches its maximum value \(I_{\text{ded}}\), the toroidal mode \(n\). Note, that the discharges \#117444 and \#117543 are RE-free.
Mechanisms of plasma disruption and runaway electron losses in tokamaks

also a sudden current drop accompanied by magnetic activity and RE bursts as seen from Figs. 1 (a) and (b) (see also e.g., Refs. [2000] [2012]). These events are probably related to the nonlinear interaction of high–energetic electrons with MHD modes which leads formation a stochastic layer at the beam edge open to the wall. We will discuss this phenomenon in Sec. 7.3. In the final termination stage one observes the quick RE current losses accompanied by magnetic activity.

As was mentioned above there are several untypical discharges for which the rates $|dI_p/dt|$ take highest or lowest values (see Fig. 1(b) and Table 1). Particularly, the current decay rate (in the 2-nd stage) for #117859 is lowest and highest for discharges #119877, #120140. Note that in the last two discharges the maximal DED current $I_{ded}$ is reached before the gas injection at $t = 2.0$ s. The RE current decay rate (in the 3-rd stage) for these discharges takes highest values. The quantity $I_p^{(RE)}$ takes the lowest value for the discharges #119877, #120140 and the highest value for #117859. One can notice strong spikes in the SXR signals of these discharges in compared to typical discharges (see Fig. 1(a)). Moreover, the above mentioned bursts of REs accompanied by magnetic activities are more pronounced in these discharges. We will discuss the peculiarity of these discharges in Secs. 3 and 5.

3. Formation of a confined plasma beam

It is believed that the plasma disruption is caused by a large scale magnetic stochasticity of field lines due to interactions of nonlinearly destabilized MHD modes ([1980] [1984] [1993] [2004] [2005] [2014]). The global stochasticity are mainly due to coupled MHD modes with low $(m,n)$ numbers: $(m = 1, \ldots, 5)$, $(n = 1, 3)$. The structure of a stochastic magnetic field depends of the amplitudes of MHD modes. Particularly, the $m/n = 1/1$ mode should play an important role on the structure of stochastic field lines near the plasma center. At low amplitudes of this mode the global stochastic field lines may not reach $q = 1$ magnetic surface and may form a confined region about the plasma center where REs can be generated. At high amplitudes of the $m/n = 1/1$ mode the stochastic field lines may cover entire plasma region with no confined particles.

As was mentioned above in TEXTOR experiments plasma disruptions with REs were deliberately caused by the injection of Ar gas while the RE–free disruptions are triggered by He/Ne injection. Experiments show that the penetration lengths of atoms depends on their atomic weights ([2008]: He (or Ne) atoms penetrate deeper into plasma than Argon atoms. The injection of these gases may finally give rise to different spectra of amplitudes of MHD modes. One can expect that the amplitude of the $m/n = 1/1$ MHD mode excited by the He/Ne injection is higher than in the case of Argon gas injection.

Possible structures of a stochastic magnetic field before the current quench with the RE-free discharge and with the RE discharge are shown in Figs. 3 (a) and (b) by the Poincaré sections of magnetic field lines. The models for the radial profiles of the plasma current $I_p(\rho)$, the safety factor $q(\rho)$ of the pre-disruption equilibrium plasma, and the MHD magnetic perturbations are given in Sec. 3 of Supplementary part. The quantity $I_p(\rho)$ describes the plasma current flowing inside the magnetic surface of radius $\rho$,

$$I_p(\rho) = 2\pi \int_0^\rho j(\rho)\rho d\rho,$$

where $j(\rho)$ is the current density. The perturbation magnetic field simulating low–mode
number MHD modes is given by the toroidal component of the vector potential

\[ A^{(1)}_z(R, Z, \varphi, t) = -\frac{R_0^2}{R} \sum_{mn} m^{-1} a_{mn}(\rho) \cos (m\vartheta - n\varphi + \Omega_{mn} t), \]

\[ a_{mn}(\rho) = B_{mn} U_{mn}(\rho), \]

with the mode amplitudes \( B_{mn} \) and rotation frequencies \( \Omega_{mn} \). Here \( B_0 \) is toroidal field strength, \( R_0 \) is the major radius, and the functions \( U_{mn}(\rho) \) describes the radial profiles of modes. It is assumed that the perturbation field contains several MHD modes: \((m/n = 1/1), (m/n = 2/1), (m/n = 3/2), \) and \((m/n = 5/2)\). In the case shown in Fig. 3 (a) the normalized mode amplitudes \( b_{mn} = B_{mn}/B_0 \) are \((1, 1, 1, 1) \times \epsilon_{MHD}\), and in Fig. 3 (b): \((1/2, 1, 1, 1) \times \epsilon_{MHD}\). The toroidal field magnitude is \( B_0 = 2.5 \) T and the dimensionless perturbation parameter \( \epsilon_{MHD} = 10^{-4} \). As seen from Fig. 3 (a) for the large amplitude of the \((m/n = 1/1)\) mode the stochastic magnetic field extends up to the central plasma region destroying the separatrix of the \( m = n = 1 \) island. For the low–amplitude of the \((m/n = 1/1)\) mode shown in Fig. 3 (b) the stochastic magnetic field does not reach the \( q = 1 \) magnetic surface and covers the region outer the \( q = 1 \) magnetic surface. The last intact drift surface is located between the resonant surfaces \( q = 1 \) and \( q = 4/3 \).

There exits a critical perturbation level \( \epsilon_{MHD} \) which breaks the intact magnetic surface between the \( q = 1 \) and \( q = 4/3 \) surfaces thus leading to the total destruction of confinement of electrons and ions. This is in agreement with experimental observations of existences of critical magnetic perturbations starting of which the runaway beams are not formed (Zeng et al. 2013).

As seen from Fig. 3 particles in the plasma core are confined by intact magnetic surfaces located between resonant surfaces \( q = 4/3 \) and \( q = 1 \). Plasma beam confined in this area is relatively stable. It contains only the \( m/n = 1/1 \) MHD mode which does not lead to a global stochasticity. The radial transport of particles from the confined area
Mechanisms of plasma disruption and runaway electron losses in tokamaks

Figure 4. Radial profile of the plasma current $I_p(\rho)$ (solid curves 1 on l.h.s. axis) and the corresponding safety factor profile $q(\rho)$ (dashed curves 2 on r.h.s. axis). The rectangular (red) dots correspond to the experimentally measured values of $I_p^{(RE)}$ for several TEXTOR discharges.

The plasma parameters are $I_p = 350$ kA, $B_0 = 2.4$ T, $R_0 = 1.75$ m, $a = 0.46$ m. The values of $q_0 = q(0)$ are 0.75 and 0.8, respectively. The radii $\rho_1$, $\rho_2$, and $\rho_3$ are the positions of the rational magnetic surfaces $q(\rho_1) = 1$, $q(\rho_2) = 3/2$, and $q(\rho_3) = 4/3$, respectively.

can take place only due to small-scale turbulent fluctuations and therefore it has much smaller rate than those in the stochastic zone. The confinement time of these electrons is sufficiently long enough to be accelerated by the inductive electric field, thus creating a RE beam. The modeling of the current of this confined plasma will be discussed in Sec. 5.

This conjecture is supported by the experimental values of the plasma current $I_p^{(RE)}$ at the initial stage of the RE beam formation. Figure 4 shows the radial profile of the pre-disruption equilibrium plasma current $I_p(\rho)$ (curve 1) and the values of $I_p^{(RE)}$ on this curve for several TEXTOR discharges. The full plasma current is $I_p = 350$ kA. As seen these values of $I_p^{(RE)}$ lie in the region between the resonance magnetic surfaces $q(\rho_1) = 1$ and $q(\rho_3) = 4/3$ or $q(\rho_2) = 3/2$. Other experimental values of $I_p^{(RE)}$ measured in the TEXTOR and given in Ref. (Zeng et al. 2013) also lie in this region.

The profiles $I_p(\rho)$ and the safety factor $q(\rho)$ in Figs. 1 correspond to the values $q(0) = 0.75$ and $q(0) = 0.8$ of the safety factor $q(\rho)$ at the magnetic axis $\rho = 0$. It corresponds to the experimentally measured values of $q(0)$ in during the sawtooth crash in the TEXTOR tokamak (Soltwisch & Stodiek 1987; Soltwisch et al. 1987) (see also (Wesson 2004) page 372). The values of $q(0)$ measured after pellet injection in the DIII-D tokamak experiments are close to these values (Izzo et al. 2012). Small changes in $q(0)$ still keeps the RE plateau currents $I_p^{(RE)}$ in the interval $\rho_1 < \rho < \rho_3$. The highest and lowest values of $I_p^{(RE)}$ shown in Fig. 4 corresponding to the discharges # 117859 and #120140, respectively, lie at the border of region $\rho_1 < \rho < \rho_3, \rho_2$. They have the shortest duration time of RE currents (see Table 1 and Fig. 1 (b)). The presence of several low-order $m/n = 4/3$, $m/n = 3/2$, and $m/n = 1/1$ resonant magnetic surfaces within the RE beam may lead to excitations of the corresponding MHD modes. The interactions of these modes may lead to the quick loss of REs due to the formation of stochastic zone at the edge of the RE beam (see Sec. 7).

The formation of the RE beam inside the intact magnetic surface can be also confirmed by the spatial profiles of the synchrotron radiation of high-energy REs with energies
exceeding 25 MeV. Figure 5 shows the radial profiles of infrared radiation of the REs at the equatorial plane $z = 0$ for the two TEXTOR discharges. One can see that radiation is localized inside finite radial extent corresponding to the central region of plasma shown in Fig. 3.

Another indication of the formation of confined plasma beam is the rise of the temperature at the initial stage of the beam formation as seen in the ECE signals shown in Figs. 1 (a) and (b). It may occur due to the Ohmic heating of confined plasma by the induced toroidal electric field or by superthermal emission from high energy electrons. With converting thermal electrons into runaway ones the beam temperature goes down.

4. Thermal and current quench stages

The strong radial transport along the stochastic magnetic field lines causes the losses of heat and plasma particles from the stochastic zone. As seen from the ECE signals in Fig. 1 the temperature drops during the fast phase of disruption on timescale of order of several 0.1 ms, while the current decays is a timescale of order of $(4 \div 6) \text{ms}$ for the RE–free discharges and up to 0.1 s for the discharges with RE generations.

The temperature drop in the fast phase can be explained by the fact that the anomalously large heat transport in a stochastic magnetic field is mainly determined by the electron transport (see, e.g., [White 2014]). The current quench stage of disruption is determined by the particle transport in stochastic magnetic field. To study the processes we use models of a stochastic magnetic field and collisional test particle transport described Sec. 7 of Supplementary part.

Figures 6 (a) and (b) shows the typical Poincaré sections of field lines of this model in the runaway-free disruption case (a) and the case with RE generation (b). The perturbation amplitudes $\epsilon_{mn}$ of all MHD modes, except $(m = 1, n = 1)$ mode, correspond to the twice larger value of $\epsilon_{MHD}$ than the case shown in Fig. 3. For the $(m = 1, n = 1)$ mode $\epsilon_{mn}$ corresponds to the same value of $\epsilon_{MHD}$. The relation between $\epsilon_{mn}$ and $\epsilon_{MHD}$ is $\epsilon_{mn} = \epsilon_{MHD} b_{mn}/\Psi_a$, where $\Psi_a$ is the toroidal magnetic flux at the plasma edge (see Sec. 7 of Supplementary part).

In general the transport of heat and particles in the presence of RMPs is a three–dimensional problem. Particularly, a stochastic magnetic field with the topological structures like ones in Figs. 6 leads to poloidally and toroidally localized heat and particle...
Mechanisms of plasma disruption and runaway electron losses in tokamaks

Figure 6. Poincaré sections of field lines in a pre–disruption plasma caused by several MHD modes: (a) runaway-free discharges; (b) with runaway electrons. The dimensionless MHD mode amplitudes are $\epsilon_{mn} = 8.68 \times 10^{-3} b_{mn}$ with $b_{mn} = 1$ for all modes in (a), and $b_{11} = 1/4$, $b_{mn} = 1$ for all other modes ($n = 1, 2$, $m = 1 - 5$) in (b). The safety factor at the magnetic axis is $q(0) = 0.8$ and at the plasma edge $q_a = 4.7$.

deposition patterns on wall (Kruger et al. 2005). This is a general feature of open chaotic systems which has been observed in ergodic divertor tokamaks (see, e.g., Finken et al. 2005; Jakubowski et al. 2006; Abdullaev 2014). The problem can be simplified when we are interested only in radial transport rate. It can be done by introducing the radial diffusion coefficient averaged over a poloidal angle.

4.1. Heat transport

The electron heat conductivity in a stochastic magnetic field has been assessed by diverse approaches. We apply here the following formula for the electron heat diffusion $\chi_r$ deduced on the basis of calculations with a model for transport of test particles, by taking into account coulomb collisions with background plasma species (Abdullaev 2013) (see also Sec. 10.4 in Abdullaev 2014):

$$\chi_r(\rho, T_e) = \frac{v_\parallel D_{FL}(\rho)}{1 + L_c/\lambda_{mfp}},$$

(4.1)

where $v_\parallel \approx v_{Te} = 1.33 \times 10^7 T_e^{1/2}$ is the thermal velocity of electrons, $D_{FL}(\rho)$ the diffusion coefficient of field lines ($D_{FL}(\rho) \sim 10^{-5} \div 10^{-4}$m), $\lambda_{mfp} = 8.5 \times 10^{21} T_e^2(\rho)/n(\rho)$ the mean free path length of electrons with the temperature $T_e$ and density $n(\rho)$ measured in keV and m$^{-3}$, respectively, and $L_c \approx \pi q(\rho) R_0$ is the characteristic connection length.

A characteristic heat diffusion time one can estimate as $\tau_H = a^2/2\chi_r$, where for $\chi_r$ we assume its magnitude at the radial position $\rho = 0.566$ a. Before the disruption the local temperature here is of 0.6 keV. This provides $\chi_r = 287$ m$^2$/s and $\tau_H = 3.68 \times 10^{-4}$ s, i.e. of the order of the experimentally observed time for the plasma temperature drop during the thermal quench after disruption.

For a quantitative analysis we have modeled the time evolution of the radial profile for the electron temperature averaged over the poloidal $\theta$ and toroidal $\varphi$ angles, $T(\rho, t)$. 
S. S. Abdullaev et al.

Figure 7. Radial profiles of the electron heat conductivity $\chi(\rho, t)$ computed according to equation (4.1), (a), and of the electron temperature averaged over toroidal and poloidal angles, $T_e(\rho, t)$, found by solving heat conduction equation (4.2) numerically, (b), at different time moments after the disruption initiation.

This is done by solving numerically the following diffusion equation:

$$\frac{\partial T}{\partial t} = \frac{1}{\rho} \frac{\partial}{\partial \rho} \left[ \rho \chi(\rho, T) \frac{\partial T}{\partial \rho} \right],$$

(4.2)

where the heat diffusivity is given by equation (4.1) and the applied boundary conditions are: $\partial T(\rho)/\partial \rho = 0$ at $\rho = 0$ and $\partial T(\rho)/\partial \rho = -T/\delta_T$ at the plasma edge $\rho = a$, where $\delta_T$ is the characteristic length of temperature decay.

Below we consider an example of heat transport in a fully chaotic magnetic field shown in Fig. 6 (a). Figures 7 (a) and (b) show the radial profiles of the heat conductivity and the temperature at different times. One can see in this case the temperature drops and almost flattens within a time interval of order of 0.5 ms.

In the situation with partially stochastic magnetic field, see Fig. 6 (b), anomalous turbulent transport in the very plasma core, $\rho \leq 0.3$, with intact magnetic surfaces is by two orders of magnitude smaller than in the outer region. In this case the fast temperature drop in the central plasma region could be explained by the presence of the $m/n = 1/1$ MHD mode and the chaotic layer near the separatrix (see a review by Schüller (1995) for more details).

4.2. Current decay stage

We assume the magnetic field structure before the current quench has the similar to the ones shown in Fig. 6. However, the level of magnetic perturbations may be higher than before the thermal quench since there are strong magnetic fluctuations after the thermal quench as seen in Fig. 4 (d).

The timescale of the current decay is determined by the rate of radial particle transport in a stochastic magnetic field. This process has the ambipolar nature and it is strongly collisional due to the low plasma temperature. On other hand one expects that the toroidal electric field induced by the current decay also strongly affects on particle transport. Below we give a rough estimation of the particle transport rate based on the collisional test particle transport model.

In Table 2 we have listed the ambipolar diffusion coefficients $D_p$ and the characteristic diffusion times $\tau_p$ of particles at the different plasma temperatures in a stochastic magnetic field shown in Fig. 6. The typical plasma temperature after the thermal quench is about from 5 eV to 50 eV. The average particle confinement time $\tau_p$ at this tem-
Mechanisms of plasma disruption and runaway electron losses in tokamaks

Table 2. Ambipolar diffusion coefficients $D_p$ of particles and the diffusion times $\tau_p = a^2/2D_p$ from the stochastic zone at the different effective plasma temperatures. The plasma radius $a = 0.46$ m.

<table>
<thead>
<tr>
<th>$T_i$ [keV]</th>
<th>$D_p$ [m$^2$/s]</th>
<th>$\tau_p$ [s]</th>
</tr>
</thead>
<tbody>
<tr>
<td>0.005</td>
<td>0.0986057</td>
<td>1.072</td>
</tr>
<tr>
<td>0.050</td>
<td>0.386249</td>
<td>$2.739 \times 10^{-1}$</td>
</tr>
<tr>
<td>0.100</td>
<td>1.01251</td>
<td>$1.045 \times 10^{-1}$</td>
</tr>
<tr>
<td>0.500</td>
<td>6.46228</td>
<td>$1.637 \times 10^{-2}$</td>
</tr>
<tr>
<td>1.000</td>
<td>9.51915</td>
<td>$1.111 \times 10^{-2}$</td>
</tr>
<tr>
<td>2.000</td>
<td>13.1030</td>
<td>$8.074 \times 10^{-3}$</td>
</tr>
<tr>
<td>4.000</td>
<td>17.8366</td>
<td>$5.932 \times 10^{-3}$</td>
</tr>
<tr>
<td>5.000</td>
<td>23.7424</td>
<td>$4.456 \times 10^{-3}$</td>
</tr>
<tr>
<td>10.00</td>
<td>27.0265</td>
<td>$3.915 \times 10^{-3}$</td>
</tr>
</tbody>
</table>

5. Modeling of post–disruption plasma

The described scenario of plasma disruption with a RE beam allows one to model a post-disruption plasma. After establishing the runaway beam the current is localized inside the area enclosed by the last intact magnetic surface. In general the distribution of the current density $j$ would depend not only on the radial coordinate $\rho$ but also vary along the poloidal $\theta$ and the toroidal $\phi$ angles due to the presence of the $(m/n = 1/1)$ magnetic island. Such a post–disruption plasma current can be presented a sum of two parts,

$$j(\rho, \theta, \phi) = j_0(\rho) + j_1(\rho, \theta, \phi), \quad (5.1)$$
where \( j_0(\rho) \) is the current density depending only on the radial coordinate \( \rho \), and \( j_1(\rho, \theta, \varphi) \) is the helical current which is a periodic function of the poloidal \( \theta \) and toroidal \( \varphi \) angles.

The radial dependence of \( j_0(\rho) \) can be also modeled by assuming that after the disruption the current is uniformly distributed over confined area with the steep gradient at the beam edge \( a \). Calculations show that electron orbits does not significantly depend on the specific of the radial profile of \( j_0(\rho) \). For our calculations of GC orbits we choose the following profile

\[
  j_0(\rho) = \begin{cases} 
    J_0 \tanh \left[ \frac{(a^2 - \rho^2)}{\Delta_a} \right], & \text{for} \quad \rho < a, \\
    0, & \text{for} \quad \rho > a,
  \end{cases} \tag{5.2}
\]

where \( J_0 \) is the constant determined by the full current of the beam \( I_p^{(RE)} \) and \( \Delta_a \) is the steepness parameter. The current flowing inside magnetic surface \( \rho \), i.e., \( I_p(\rho) = 2\pi J_0 \int_0^\rho j_0(\rho')d\rho' \) is given by

\[
  I_p(\rho) = \begin{cases} 
    I_p^{(RE)} \left[ 1 - \frac{\ln \cosh \left( \frac{(a^2 - \rho^2)}{\Delta_a} \right)}{\ln \cosh \left( \frac{a^2}{\Delta_a} \right)} \right], & \text{for} \quad \rho \leq a, \\
    I_p^{(RE)}, & \text{for} \quad \rho > a,
  \end{cases} \tag{5.3}
\]

where \( I_p^{(RE)} \) is the full current of the confined area.

One should also note the fact that after the thermal quench the plasma beam is shifted inwardly because of drop of plasma pressure. In the modeling this fact can be taken into account by assuming that the radial position of the center \( R_a \) of the post–disruption plasma is different from the one of the pre–disruption plasma. The safety factor of the corresponding plasma is then given by

\[
  q(\rho) = q_{cyl}(\rho)C(\rho/R_a), \quad q_{cyl}(\rho) = \frac{2\pi \rho^2 B_0}{\mu_o R_a I_p(\rho)}, \tag{5.4}
\]

where \( q_{cyl}(\rho) \) is the safety factor of the cylindrical plasma, the function \( C(x) = 1 + A_1 x + A_2 x^2 + \cdots \) is a function which takes into account the toroidicity of plasma. The coefficients \( A_i \), \( i = 1, 2, \ldots \) depends on the plasma pressure (Abdullaev et al. 1999).\(^\text{[1]}\)


Figure 8 shows the radial profiles of \( I_p(\rho) \) (solid curves 1–3 on the l.h.s. axis) \(^{[5,3]}\) and the safety factor \( q(\rho) \) (dashed curves 1’ – 3’ on the r.h.s. axis) for the three discharge parameter, respectively. Solid black curve 4 corresponds to the pre–disruption plasma current profile. We set the toroidal field magnitude \( B_0 = 2.4 \) T, the beam center at \( R_a = 1.7 \) m. The plasma radius \( a \) is found from the condition \( I_{p0}(a) = I_p^{(RE)} \), where \( I_{p0}(r) \) is the current profile of the pre–disruption plasma. The vertical dashed color arrows show the radial positions of the \( q = 1, q = 4/3, \) and \( q = 3/2 \) magnetic surfaces and the vertical solid arrows indicate the plasma radii \( a \).

Note that that the red curves 1 and 1’ and green curves 3 and 3’ in Fig. 8 correspond to the discharges with the lowest and highest values of \( I_p^{(RE)} \) shown in Fig. 1(b). For the lowest value of \( I_p^{(RE)} \) the radial position of the \( q = 1 \) magnetic surface is very close to the plasma radius. For the highest value of \( I_p^{(RE)} \) the magnetic surfaces with \( q = 1, q = 4/3, \) and \( q = 3/2 \) are located inside the plasma region \( \rho < a \). However, the radial position of the magnetic surface \( q = 3/2 \) is at the plasma edge. For the typical discharges like the one shown by blue curves the magnetic surface \( q = 1 \) is located relatively far from the plasma edge.
Figure 8. Radial profiles of the plasma current $I_p(\rho)$ (solid curves 1–3 on l.h.s. axis), the safety factor profiles $q(\rho)$ (dashed curves $1'–3'$ on r.h.s. axis), and curve 4 corresponds to the pre–disruption plasma current. The red curves 1 and 1' correspond to $I_p^{(RE)} = 165 \text{kA}$, blue curves 2 and 2' correspond to $I_p^{(RE)} = 230 \text{kA}$, and green curves 3 and 3' correspond to $I_p^{(RE)} = 300 \text{kA}$. The vertical solid arrows indicate the radii of plasma beam $a$, the vertical dashed arrows indicate the positions of resonant magnetic surfaces $q = 1$ and $q = 3/2$. The toroidal magnetic field $B_t = 2.4 \text{T}$, $R_a = 1.7 \text{m}$, the pre–disruption plasma current $I_p = 350 \text{kA}$ and the radius $a_0 = 0.46 \text{m}$.

The Fourier expansion of the helical current, $j_1(\rho, \theta, \varphi)$,

$$j_1(\rho, \theta, \varphi) = \sum_{m,n} j_{mn}(\rho) \cos(m\theta - n\varphi + \phi_{mn}),$$

is mainly dominated by the $m/n = 1/1$ component. This assumption is based on the analysis of numerous disruptions in the JET tokamak (Gerasimov et al. 2014).

We should assume that the value of the safety factor at the beam axis $q(0)$ is less than unity. This assumption is supported by a number of experimental measurements of the current profile after the sawtooth crashes in the TEXTOR, the TFTR, and JET tokamaks (Soltwisch et al. 1987; Yamada et al. 1991; Soltwisch & Koslowski 1995; O’Rourke 1991; Koslowski et al. 1996; Soltwisch & Koslowski 1997).

This model of the post–disruption plasma current describes only the initial stage of RE beam. During acceleration of electrons in the toroidal electric the RE orbits drift outward and their form evolves from the circular one to oval one. This process changes in turn the RE beam form and its current. The self–consistent description of time evolution of RE beam is difficult and it requires a separate study.

6. Evolution of GC orbits during acceleration

6.1. Outward drift of RE orbits

First we consider the case of the axisymmetric plasma beam neglecting the helical magnetic perturbations. The inductive toroidal electric field generated due to the current decay during the plasma disruption accelerates thermal electrons. This is an adiabatic process since the characteristic time of significant variation of energy is much larger than the transit time of electrons. Therefore the GC orbit slowly drifts outward without changing the area of GC orbit in the poloidal plane which is an adiabatic invariant $J$ or the action variable (see Sec. 6.1 of Supplementary part)). With increasing electron energy the topology of GC orbits also slowly changes from the circular one to the oval one.
Starting from the certain critical energy $E_{cr}$ the adiabaticity of the process breaks and the GC orbit bifurcates by creating the unstable stagnation point (or X-points) inside the plasma region. With the further increase of energy the GC orbit crosses the separatrix (a homoclinic orbit associated the X-point) and becomes unconfined. The value $E_{cr}$ depends on the plasma current $I_p$. The described phenomenon is an addition mechanism of confinement loss of REs. Figure 9 (a) shows a typical evolution of a GC orbit in the presence of the toroidal electric field with the constant beam current $I_p = 100 \text{ kA}$ and the loop voltage $V = 40 \text{ V}$.

One should note that the formation of the separatrix of RE GC orbits during the acceleration process in tokamaks has been first predicted in Ref. (Zehrfeld et al. 1981). The numerical study of this process in a realistic tokamak configuration has been carried out in Ref. (Wongrach et al. 2014b). Particularly, it was shown that with increasing electron energy the area confined by the separatrix decreases and it vanishes when the energy exceeds a certain critical value $E_{cr}$, i.e. such electrons cannot be confined. The critical energy $E_{cr}$ is proportional to the square root of the plasma current $I_p$, $E_{cr} \propto \sqrt{I_p}$.

The described evolution of RE orbits is in agreement with the experimental observation of the IR radiation patterns observed in the experiment in the TEXTOR tokamak (Wongrach et al. 2014b) and in the DIII-D tokamak (Hollmann et al. 2013). The observations clearly show the evolution of the spatial form of RE beam from crescent ones into oval ones with increasing the electron energies.

The example of the time-evolution of GC orbits in the plasma beam with a time-varying current $I_p(t)$ and the loop voltage $V(t)$ corresponding to the TEXTOR discharge #117527 is shown in Fig. 9 (b). To simplify the calculations of orbits we have assumed that the loop voltage $V(t)$ is uniform in the poloidal section, i.e., it does not depend on the radial coordinate $r$ and equal to the experimentally measured value at the limiter. However, this assumption only approximately describe the situation. To find more the exact magnitudes of the toroidal electric field during the runaway current decay one should solve the corresponding Maxwell equations.

One of the important parameter of the GC orbit is the effective safety factor $q_{eff}$ defined as a ratio $q_{eff} = \Delta \varphi/2\pi$ where $\Delta \varphi$ is the increment of the toroidal angle $\Delta$ per one poloidal turn. It is a function of the action variable $J$ and particle energy $E$. For
Mechanisms of plasma disruption and runaway electron losses in tokamaks

Figure 10. Time–evolution of the effective safety factors $q_{eff}$ (l.h.s. axis) and electron energies (r.h.s. axis) during the acceleration in the discharge #117527. Blue curves correspond to the orbit launched at the coordinate $(R = 160, Z = 0)$ cm, and red curves to the one with $(R = 165, Z = 0)$ cm. Horizontal lines correspond to $q(t) = m$ where $m = 1, 2, \ldots$ are the integer numbers.

Low–energy electrons the quantity $q_{eff}(J, E)$ coincides with the safety factor $q(\rho)$ of the equilibrium magnetic field. With increasing the electron energy the effective safety factor strongly deviates from $q(\rho)$. With approaching RE energy $E$ to the critical one $E_{cr}$ it diverges as

$$q_{eff}(J, E) \propto -\ln |E - E_{cr}|. \quad (6.1)$$

Figure 10 shows the typical time evolutions of the effective safety factors $q_{eff}$ of two GC orbits during the electron acceleration in the conditions of the TEXTOR discharge #117527.

Figure 11 shows the time–evolution of the outward drift velocity $v_{dr}$ calculated numerically for the three different RE beam currents. It is quite well described by formula derived in Ref. (Abdullaev 2015)) (see also Sec. 4 of Supplementary part)

$$v_{dr} = \frac{E_{\varphi}}{B_z^*} \left( \frac{R_0}{R} - \frac{T_{av}}{T} \right), \quad (6.2)$$

where

$$B_z^* = B_z + F(E), \quad T_{av} = \frac{2\pi q_{eff} R_0}{v_{\varphi}}. \quad (6.3)$$

In (6.2) and (6.3) the quantity $B_z$ is the $Z$– component of the poloidal magnetic field at the equatorial plane $z = 0$, $E_{\varphi}$ is the toroidal electric field strength, $T$ is the transit time of orbit, $v_{\varphi}$ is the toroidal velocity, $F(E)$ is the term depending on a particle energy.

At $|B_z| \gg |F(E)|$, and $T_{av} \approx T$ the formula is reduced to

$$v_{dr} = \frac{qE_{\varphi}}{B_0} = -\frac{(R - R_0)E_{\varphi}}{RB_z}, \quad (6.4)$$

obtained by Guan et al. (2010) and Qin et al. (2011) for the circular orbits. Here $q = (R - R_0)B_0 / B_z R$ is the safety factor of magnetic field. As seen from Fig. 11 the formulas (6.2) and (6.4) give the correct dependence of $v_{dr}$ on the plasma current $I_p$, $v_{dr} \propto I_p^{-1}$ because $B_z \propto I_p$.

However, the formula (6.4) does not describe the situation when the GC orbits take an oval form with increasing the energy similar to the ones shown in Figs. 11 (a) and (b). From the latter it follows that the average outward velocity $v_{dr}$ of the innermost part of the orbit is approximately equal 0.6 m/s and 8 m/s of the outermost part of the orbit.
Drift velocities of innermost $v_{df}(R_i)$ (curves 1, 2, and 3) and outermost $v_{df}(R_o)$ (curves 1', 2', and 3') points of orbits for the different plasma current: curves 1 and 1' correspond to the plasma current $I_p = 100$ kA, curves 2 and 2' correspond to $I_p = 200$ kA, and curves 3 and 3' correspond to $I_p = 300$ kA. Curve 4 describes the increase of energy $E$ (right hand axis).

The toroidal field $B_t = 2.5$ T, major radius $R_0 = 175$ cm, minor radius $a = 46$ cm, the loop voltage $V_{loop} = 5$ V. Note that $v_{dr}$ is multiplied to the proportionality factor $f$.

6.2. RE current decay

The rate $dI_p/dt$ of the runaway current loss due to described outward drift of orbits can be roughly estimated as follow. This loss mechanism is mainly caused by the shrinkage of the beam radius $a$. The rate of such a shrinkage $da/dt$ is of order of the average outward velocity $v_{dr}$. Since $I_p \propto a^2$, we have

$$\frac{dI_p}{dt} = \frac{dI_p}{da} \frac{da}{dt} \sim \frac{2I_p}{a} v_{dr}. \quad (6.5)$$

For the typical values of $I_p \approx 0.2$ MA, $a \approx 0.2$ m, and $v_{dr} \sim 1$ m/s one has $dI_p/dt \approx 4$ MA/s. This estimation is of order of the experimentally measured average decay rate of the runaway current listed in Table 1.

Since the safety factor $q$, as well as $q_{eff}$ of RE beams is about unity, $q_{eff} \sim 1$ then the outward drift slows down for the higher values of the toroidal magnetic field $B_0$. It means that the decay time of RE currents in large tokamaks, like ITER, will be much longer than in smaller tokamaks.

Beside of outward orbit drifts the RE current losses is also caused by the internal MHD mode which be discussed in the next section. The collisions of REs with neutral particles may also contribute to the RE losses.

7. Effect of magnetic perturbations

The effect of the magnetic perturbations on electrons in the post–disruption current beam strongly depends on its safety factor profile $q(\rho)$, the spectrum of magnetic perturbations, and the electron energy. To explain this effect we consider the simplified version of GC motion equations in the presence of magnetic perturbations. (The rigorous consideration of this problem is given in Sec. 6 of Supplementary Part).

The particle drift motion in the presence perturbations can be presented by Hamiltonian equations similar to the equations for magnetic field lines,

$$\frac{d\theta}{d\varphi} = \frac{\partial K}{\partial J}, \quad \frac{dJ}{d\varphi} = -\frac{\partial K}{\partial \theta}, \quad (7.1)$$
Mechanisms of plasma disruption and runaway electron losses in tokamaks

Figure 12. (a) Spectrum of perturbations $K_{mn}$ and (b) corresponding RE orbits with different energies $E$. Curves 1–7 correspond to RE energies 10 keV, 20 MeV, 30 MeV, 40 MeV, 42 MeV, 42.5 MeV, and 42.7 MeV, respectively. Curve 8 corresponds to the separatrix with the critical energy $E_{cr} = 42.646$ MeV. The plasma current $I_p = 150$ kA, the toroidal field $B_0 = 2.5$ T. The toroidal mode number $n = 1$.

with the Hamiltonian $K = K(\theta, J, E, \varphi)$ with the canonical variables $(\theta, J)$, and the toroidal $\varphi$ as the time–like variable. In the absence of perturbations GC orbits wound the drift surfaces $J = \text{const}$ and the poloidal angle $\theta$ is a linear function $\varphi = \varphi/q_{eff}(J, E) + \theta$. In the presence of perturbations Hamiltonian $K$ can be presented as a sum

$$K = \int \frac{dJ}{q_{eff}(J, E)} + \epsilon K_1(\theta_z, J, E, \varphi).$$

(7.2)

Since the perturbation are periodic in poloidal and toroidal angles and in time it can be presented by a Fourier series

$$K_1(\theta_z, J, E, \varphi) = \sum_{mn} K_{mn}(J, E) \exp[i(m\theta - n\varphi)].$$

(7.3)

The strongest influence of perturbation on particles takes place on the $(m, n)$ resonant drift surfaces, i.e.,

$$m = n q_{eff}(J, E),$$

(7.4)

originating from the $(m, n)$ term in (7.3) with the amplitude $K_{mn}(J, E)$. They are determined by the magnetic perturbation spectrum $b_{mn}$,

$$K_{mn}(J, E) \propto \sum_{m'} b_{mn} \int_0^{2\pi} d\theta \exp[i(m\theta - m'\varphi_M)],$$

(7.5)

where $\varphi_M$ the poloidal angle associated with magnetic field lines is a function of $\varphi$ as well as particle energy $E$.

For low–energy electrons (up to 5 MeV) the spectrum of amplitudes $K_{mn}(J, E)$ weakly depends on energy $E$ and close to the spectrum of magnetic perturbations $b_{mn}$ of $(m, n)$–th modes. With increasing the energy the spectrum of perturbations $K_{mn}(J, E)$ deviates from $b_{mn}$ and acquires more higher poloidal harmonics $m$. The example of the poloidal spectra of perturbation $K_{mn}(J, E)$ for different particle energies is shown in Fig. 12 (a). The corresponding unperturbed orbits are plotted in Fig. 12 (b). It is assumed that the magnetic perturbation contains a single $(m = 1, n = 1)$ mode. For the low energy electrons with $E < 10$ MeV the spectrum $K_{mn}$ contains the predominant $m = 1$ mode.

With increasing the energy the amplitudes $K_{mn}$ of higher $m$ also grow and the width
of the poloidal spectrum $K_{mn}$ in $m$ becomes wider as shown in Fig. 12 (b). For the spectrum $K_{mn}$ one can obtain the following asymptotical formula for the orbits close to the separatrix (see Sec. 3.4 in [Abdullaev 2014])

$$K_{mn} \propto \frac{1}{q_{eff}} \exp \left( -\frac{mC}{q_{eff}} \right),$$

(7.6)

where $C$ is a finite constant, and the effective safety factor $q_{eff}$ diverges as $[0,1]$. As was shown in Sec. 5 (see also Fig. 8) the typical values of $q(\rho)$ varies between $q(0) \approx 0.7 \div 0.8$ at the magnetic axis and $q(a) < 1.5$ at the plasma edge. Therefore, the strongest effect of the RMPs on electron orbits may expect if its spectrum $b_{mn}$ contains a sufficient number of ($m,n$) components that are resonant to the magnetic surfaces with $q$ in the interval $q(0) < q = m/n < q(a)$ that would create a stochastic zone of magnetic field lines. The electrons from this stochastic layer would be then radially transported to wall.

Below we discuss the influence of magnetic perturbations on RE orbits for the two specific cases. First we consider the effect of internal single helical magnetic field, and then we analyze the effect of the external RMPs, namely, the TEXTOR-DED on the confinement of REs.

7.1. Effect of a single helical magnetic field

Assume the magnetic perturbation (3.2) contains the single ($m = 1, n = 1$) MHD mode as was proposed in the model of the post–disruption current beam described in Sec. 5. For the low–energy electrons it creates a single island structure since the deviations of their GC orbits from the magnetic surfaces is small. Such a system is stable because the single MHD mode does not create stochasticity of magnetic field lines. The example of this case is shown in Fig. 13 (a) by the Poincaré sections of RE orbits (red dots) and magnetic field lines (blue dots).

With increasing the energy of electrons and decreasing the beam current the electron’s GC orbits strongly deviate from the magnetic field lines. The effective safety factor $q_{eff}$ of the GC orbit increases as the electron energy grows as was shown in Fig. 10. At certain time instants the value of $q_{eff}$ reaches the integer value so that the resonant condition may be satisfies for the higher harmonics ($m > 1, n > 1$) of the GC orbits with the ($m = 1, n = 1$) magnetic perturbation. This generates a number of island chains of GC orbits. The interaction of several such island structures may even lead to the formation of the stochastic layer near the separatrices [see Figs. 9 (a) and (b)].

Figure 13 (b) illustrates the typical structure of high–energy electrons in the presence of the internal helical magnetic field with a single ($m = 1, n = 1$) mode. Such a structure leads to the widening area of lost electrons and decreasing the critical energy $E_{cr}$. The characteristic escape time of REs from the stochastic layer is of order of 10 $\mu$s. Sudden RE bursts in many discharges is probably related with the loss REs from the stochastic layer. Occurrence of the MHD mode signals accompanied these events will be discussed in the next Sec. 7.3.

As was discussed above in Secs. 2 and 5 (also Figs. 1 (b) and 8) there are some exceptional discharges (for example #117859) with the highest RE current and several low–order rational surfaces within plasma beam. Such a beam can be easily destabilized by the magnetic perturbations containing several MHD modes with low–order ($m,n$) numbers. Such a magnetic perturbation may affect strongly on electrons creating the chaotic zone at the beam edge open to wall. Such an effect probably explains the sudden lost of REs at certain times seen in Fig. 11 (b).
Mechanisms of plasma disruption and runaway electron losses in tokamaks

7.2. Influence of the TEXTOR–DED

The coil configuration of the TEXTOR–DED is designed to have the poloidal spectra of magnetic perturbations localized near the magnetic surface $q = 3$ of the flat–top plasma discharges (see Sec. 5.2 of Supplementary part). Therefore, these perturbations do not contain a necessary number of resonant components to create a stochastic zone of magnetic field lines in the post–disruption current beam with the safety factor $q$ lying between $q(0) < 1$ and $q(a) < 1.5$.

In the so–called 3/1 operational mode with the predominant toroidal mode $n = 1$ toroidal there is only one ($m = 1, n = 1$) component resonant to the magnetic surface $q = 1$. The similar situation takes place in the 6/2 mode ($n = 2$) with the resonant component ($m = 2, n = 2$). [There are no magnetic surfaces in the plasma region that are resonant to the components ($m = 1, n = 2$) and ($m = 3, n = 2$)]. On the other hand this resonant component of the DED field is weak since it is located away from the maximum of the spectrum. Therefore, the effect of the DED on the RE beam does not create the stochastic zone of magnetic field lines from which electrons would escape to wall as in the case of the stochastic zone in a flat–top plasma operation. The only $m/n = 1/1$ component of the DED perturbations may create an island structure near the $q = 1$ magnetic surface similar to one shown in Fig. 13 (a).

With increasing the energy of REs and decreasing the plasma current the DED perturbation starts to affect on REs because of appearance of high–mode resonances $q_{ef}f = m/n$ similar to the case discussed in Sec. 7.1. It generates the structures with islands and a stochastic layer. Figures 14 (a) and (b) show the typical Poincaré sections of GC orbits of energetic electrons affected by the TEXTOR–DED: (a) corresponds to the 3/1 mode with the DED current $I_{ded} = 3$ kA; (b) corresponds to the 6/2 mode with $I_{ded} = 7$ kA. The particle energy is taken $E = 20$ MeV, the plasma current $I_p = 94$ kA, and the toroidal field $B_0 = 2.5$ T. These structures explain the fast decay of RE current in its final stage accompanied by spikes in the scintalation probe (see Fig. 1 (b)).

The structures of RE orbits shown in Figs. 13 (b) and 14 correspond to the final termination stages of RE current. They have features which are characteristic for the
so-called stable and unstable manifolds created by the splitting of separatrices (see, e.g., Abdullaev 2014). They lead to the toroidally and poloidally localized deposition patterns of REs on wall. Toroidal peaking and spatial–temporal evolution of hard X-ray emission in the final stage of RE current loss observed in DIII-D experiments (James et al. 2012) is consistent the described topology of REs.

The experimental observations in the TEXTOR-DED have indeed showed that the RMPs field which switched on just after the thermal quench does not affect on the radial transport and the loss of low–energy electrons (Koslowski et al. 2014; Wongrach et al. 2014) (see also Table 1). This is mainly because of the mentioned features of the poloidal and toroidal spectra of the DED field.

7.3. Generation of magnetic perturbations by high–energy electrons

The above mentioned in Sec. 2 the occurrence of the MHD activities during the sudden RE bursts can be explained by the nonlinear interaction of high–energy electrons with the $(m = 1, n = 1)$ MHD mode. The MHD magnetic perturbations with mode numbers $(m, n)$—higher than the initial $(m = 1, n = 1)$ mode can be generated during the acceleration process of the REs. At certain energy of REs their orbits strongly deviate from the magnetic surfaces which creates in turn higher $(m, n)$–harmonics, $(m > 1, n > 1)$, of the MHD $(m = 1, n = 1)$ mode (7.5). The resonant interaction of RE orbits with these harmonics leads to the redistribution of corresponding current near these orbits according to the helicity of these modes. Therefore, the current density (5.5) acquires higher $(m, n)$–components $j_{mn}$ which in turn generates the corresponding MHD modes. The bursts of magnetic activities accompanied by accompanied by sudden runaway current drops observed in experiments (see Figs. 1 (a) and (b)) are probably related to the described phenomenon.

8. Summary

Based on the analysis of numerous experimental data obtained in the TEXTOR tokamak we have proposed a possible mechanism of the plasma disruption with the formation of RE beams. The plasma disruption starts due to a large–scale magnetic
Mechanisms of plasma disruption and runaway electron losses in tokamaks

23

stochasticity caused by nonlinearly excited of MHD modes with low \((m,n)\) numbers \(\left((m/n = 1/1, 2/1, 3/2, 5/2, \ldots)\right)\). At the sufficiently small amplitude of the \(m/n = 1/1\) mode there exists an intact magnetic surface located between the magnetic surface \(q = 1\) and the low–order rational surface \(q = m/n\). Depending on the spectrum of magnetic perturbations this rational could be one of these ones: \(q = 4/3, q = 5/4\) or \(q = 3/2\).

Such intact magnetic surface forms the transport barrier for particles in the central plasma region. Electrons in this confined region are accelerated by the inductive toroidal electric field. Such a situation occurs, for instance, in plasma disruptions with runaway beams initiated by the Argon gas injection. Heavy Ar atoms do not penetrate sufficiently deep into the plasma and therefore they do not excite the \(m/n = 1/1\) mode with the amplitude necessary to create the fully chaotic magnetic field. On the other hand the injection of the lighter noble gases neon and helium does not generate runaways since the light gases penetrate deeper into the plasma and excite the large–amplitude \((m/n = 1/1)\) mode.

Based on this scenario we proposed the models of the pre-disruption and post-disruption plasmas with REs to study the processes of thermal and currents quenches, the runaway current losses. The model of magnetic field was proposed to describe the large–scale magnetic stochasticity due to interaction of low–mode–number MHD modes. The radial transport of heat and particles in a stochastic magnetic field are studied using the collisional diffusional models. It was shown that the temperature drop during the fast phase of disruption is caused by the radial heat transport determined by the collisional electron transport in a stochastic magnetic field. We have estimated a current decay time using the ambipolar collisional particle transport model. The dynamics of RE orbits in a post–disruption plasma in the presence of the inductive toroidal electric field is investigated by integrating the equations of guiding center motion. We analyzed the effect of the internal MHD mode and external RMPs on the topology of RE orbits.

The new model reproduces for the first time remarkably well the essential features of the measurements:

(a) The outer part of the plasma is clearly ergodized while the inner section is still intact. This agrees with the observation that the runaways are only seen in the inner half of the torus while they are obviously quickly lost from the outer part.

(b) one observes a short tiny spike during the energy quench; we have described this spike previously; this spike is attributed to the loss of runaways born at the start up of the discharge from the ergodic zone.

(c) The duration of the energy quench is determined by the strong electron diffusion in a stochastic magnetic field. The plasma current decay is mainly due to the ambipolar particle transport in the presence of the inductive toroidal electric field. The estimations of the energy quench and the current decay times based on the model well agree with observations.

(d) The slow decay of the RE current in the plateau phase is explained by the loss of runaways due to two effects: (i) an outward shift of the runaways due to their continuous acceleration and the subsequent loss at the wall; (ii) by the formation of a stochastic layer of high–energy REs at the beam edge in the presence of the \(m/n = 1/1\) MHD mode.

(e) The effect of the external resonant magnetic perturbations on low-energy electrons (up to 5-10 MeV) is weak and does not cause their loss.

The new mechanism explains well the observed disruptions in present day tokamaks. One can expect the following consequences, e.g., for ITER.

(1) The structure of the stochastic zone during the thermal quench allows persistence
of preexisting runaways through this phase such that they act as seeds during the following phase of high loop voltage.

(2) The decay phase of the REs is rather long such that REs can acquire very high energy.

(3) External magnetic perturbations acting on REs seem little promising unless the core of a RE beam can be ergodized.

(4) A means for eliminating the REs completely is the injection of about $10^{25}$ molecules $\text{H}_2$ or $\text{D}_2$ into the discharge [Hender et al. 2007]. This massive gas injection may impose a heavy load on the cryo-pumping system.

REFERENCES


Bécoulet, et al. 2013 Science and technology research and development in support to ITER and the Broader Approach at CEA. Nuclear Fusion 53 (10), 104023.


Mechanisms of plasma disruption and runaway electron losses in tokamaks


Mechanisms of plasma disruption and runaway electron losses in tokamaks


