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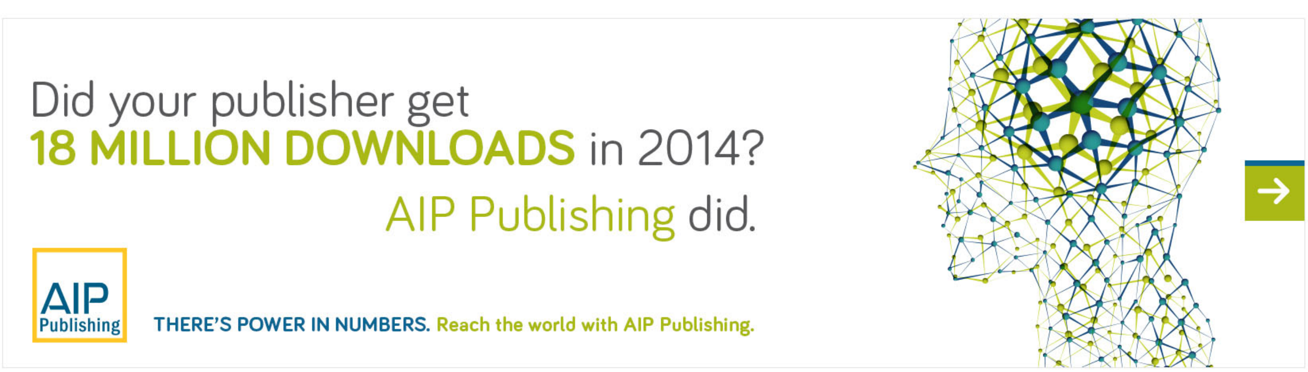

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Mechanism of runaway electron beam formation during plasma disruptions in tokamaks

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A new physical mechanism of formation of runaway electron beams during plasma disruptions in tokamaks is proposed. The plasma disruption is caused by a strong stochastic magnetic field formed due to nonlinearly excited low-mode number magnetohydrodynamic (MHD) modes. It is conjectured 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 magnetic surfaces [$q = 5/4$ or $q = 4/3,\ldots$]. It results in that runaway electron beam current has a helical nature with a predominant $m/n = 1/1$ component. The thermal quench and current quench times are estimated using the collisional models for electron diffusion and ambipolar particle transport in a stochastic magnetic field, respectively. Possible mechanisms for the decay of the runaway electron current owing to an outward drift electron orbits and resonance interaction of high–energy electrons with the $m/n = 1/1$ MHD mode are discussed. [http://dx.doi.org/10.1063/1.4919253]

The runaway electrons (REs) generated during the disruptions of tokamak plasmas may reach a 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.[1–5]

The mitigation of REs by massive gas injections (MGI) and externally applied resonant magnetic perturbations (RMPs) have been extensively discussed in literature (see, e.g., Refs. 6–8 and references therein). However, no regular strategy to solve this problem has been developed because up to now the physical mechanisms of the formation of REs during plasma disruptions are not well understood. In spite of the numerous dedicated experiments to study the problem of runaway current generation during plasma disruptions in different tokamaks (see, e.g., Refs. 7–15), no clear dependence of RE formation on plasma parameters has been established. These numerous experiments show the complex nature of plasma disruption process especially the formation of RE beams.

One of the important features of the formation of RE beams is the irregularity and variability of the beam parameters from one discharge to another one. This is an indication of the sensitivity of RE beam formations on initial conditions which is the characteristic feature of nonlinear processes, particularly, the chaotic system. Therefore, one expects that *ab initio* numerical simulations of the RE formation process may not be quite productive to explain it because of complexity of computer simulations of nonlinear processes.[16] The problems of numerical simulations of plasma disruptions are comprehensively discussed in Ref. 17.

In this work, we propose a new physical mechanism of formation of RE beams during plasma disruptions in tokamaks. It is based on the analysis of numerous experimental results, mainly obtained in the TEXTOR tokamak and the ideas of magnetic field stochasticity.[18] The mechanism explains many features of plasma disruptions accompanied by RE generations.

It is believed that the plasma disruption starts with the excitation of magnetohydrodynamic (MHD) modes with low poloidal m and toroidal n numbers, ($m/n = 1/1$, 2/1, 3/2, 5/2,…) that lead to a large–scale magnetic stochasticity (see, e.g., Refs. 19–22 and references therein). The heat and particle transports in the strongly chaotic magnetic field cause the fast temperature drop and cease the plasma current. This process depends on the structure of the stochastic magnetic field which depends on the spectra of magnetic perturbations and on the safety factor profile $q(\rho)$ (ρ is the minor radius of the magnetic surface). At certain conditions, the stochastic magnetic field may not extend up to the central plasma region due to the creation of the outermost intact magnetic surface ρ_c. The electrons confined by this magnetic surface are accelerated by the toroidal electric field induced by the current decay from the outer plasma region, which leads to the formation of the RE beam. The initial RE current $I_p^{(RE)}$ is mainly determined by the pre-disruption plasma current distribution $I_p(\rho)$ confined by the outermost intact magnetic surface ρ_c, i.e., $I_p^{(RE)} \approx I_p(\rho_c)$.

The lifetime of the RE beam mainly depends on two effects: the outward drift of RE orbits induced by the toroidal electric field E_φ[32,33] and the resonant interactions of REs with helical magnetic perturbations. The first one is responsible for the smooth decay of the RE current, while the second one cause the sudden RE losses. The outward drift velocity v_{dr} is determined by E_φ and the RE current, $v_{dr} \propto E_\varphi / I_p^{(RE)} \propto E_\varphi / \rho_c^2$.[32,33] The most stable of the RE beams is expected to form when the corresponding drift velocity is lowest and the low–order rational surfaces within the RE beam are absent or one.

Consider, for example, the pre-disruption plasma with a monotonic safety factor profile $q(\rho)$ with $q(0) < 1$. Then the most stable RE beam can be formed when the outermost intact magnetic surface is located between magnetic surface $q = 1$ and the nearest low–order rational surfaces $q = 5/4$ [or $q = 4/3, \ldots$]. It occurs at the sufficiently small amplitude of the $m/n = 1/1$ mode. There is only one rational magnetic surface $q = 1$ within the RE beam that is resonant to the large–scale magnetic perturbations, particularly, to the RMPs. Such RE beams are relatively stable, since low–energetic REs (up to 10–15 MeV) are not destabilized due to absence of a large scale stochasticity. The loss of REs mainly occurs due to the outward drift of RE orbits and the stochastic instability of high–energetic REs due to the interactions of high–mode harmonics of the $m/n = 1/1$ mode of magnetic perturbations.

In the case of plasma disruptions with $q(0) > 1$, the intact magnetic surface ρ_c would be smaller while the toroidal electric field E_φ would be larger than in the ones with $q(0) < 1$. Due to the large outward drift velocity v_{dr}, such RE beams would cease faster.

The two possible distinct generic structures of a stochastic magnetic field before the current quench (CQ) with the RE-free discharge and with the RE discharge are shown in Figs. 1(a) and 1(b) by the Poincaré sections of magnetic field lines. It is assumed that the perturbation magnetic field contains several low–mode number m/n MHD modes with equal amplitudes B_{mn}: (a) the amplitude B_{11} of the $m/n = 1/1$ mode is equal to others; (b) B_{11} is four times smaller than the amplitudes of other modes. As seen from Fig. 1(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. 1(b), the stochastic magnetic field does not reach the $q = 1$ magnetic surface. The last intact drift surface (red dots) is located between the resonant surfaces $q = 1$ and $q = 5/4$ (blue curves).

The existence of an intact magnetic surface and its location depends on the radial profile of the safety factor and on the spectrum of magnetic perturbations. The latter sensitively depend on the plasma disruption conditions and vary unpredictably 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.

This conjecture on the mechanism of RE beam formation agrees with the important features of the experimental observations in the TEXTOR tokamak. In the experiments, the plasma disruptions were triggered by gas injections (see, e.g., Refs. 9–11,23): the disruptions with REs were triggered by argon (Ar) injection and the RE–free disruptions with Ne injection. 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 Ar gas injection.

The plasma current decay in the CQ and the RE plateau regimes for all discharges is well approximated by the linear function of time $I_p = I_{p0} + bt$, with the average CQ rate $b = \langle dI_p/dt \rangle$ as shown in Fig. 2(a). The current decay rates $|\langle dI_p/dt \rangle|$ in the CQ stage and the RE plateau stage versus the initial RE current $I_p^{(RE)}$ for a number discharges are plotted Fig. 2(b). The plausible radial profiles of $I_p(\rho)$ and the corresponding safety factor $q(\rho)$ are plotted for the two values of $q(0)$ in Fig. 3.

Since ρ_c is located between the magnetic surfaces ρ_1 and ρ_3 corresponding to $q(\rho_1) = 1$ and $q(\rho_3) = 4/3$, the RE current $I_p^{(RE)}$ should take values in the finite interval. This expectation is supported by the experimental values of the plasma current $I_p^{(RE)}$ as seen from Figs. 2(a) and 2(b). These values of $I_p^{(RE)}$ also lie in the region between the resonance magnetic surfaces $q(\rho_1) = 1$ and $q(\rho_3) = 4/3$ [or $q(\rho_2) = 3/2$] as shown in Fig. 3 where the radial profile of the pre–disruption equilibrium plasma current $I_p(\rho)$ (curve 1) and the corresponding safety factor profile $q(\rho)$ (curve 2) are plotted.

The average values of $\langle |dI_p/dt| \rangle$ 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 CQ stage. The values of $I_p^{(RE)}$ are in the range between 170 kA and 260 kA (see Fig. 2(b)). These values of $\langle |dI_p/dt| \rangle$ and $I_p^{(RE)}$ are close to the ones observed in the similar experiments in the DIII-D tokamak (see, e.g., Ref. 15).

As seen from Fig. 2(a), there are untypical discharges with the highest and lowest values of $I_p^{(RE)}$ that correspond to

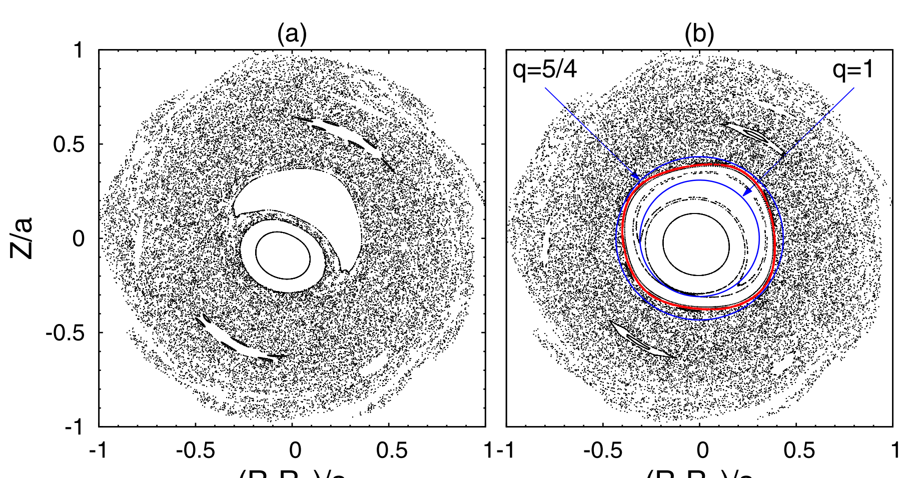

FIG. 1. Poincaré sections of magnetic field lines in a pre–disruption plasma caused by several m/n MHD modes, ($n = 1, 2, 3$; $m = 1, \ldots 8$): (a) all mode amplitudes B_{mn} are equal; (b) the amplitude B_{11} of the $m/n = 1/1$ mode is four times smaller than B_{mn}. The safety factor at the magnetic axis $q(0) = 0.8$ and at the plasma edge $q(a) = 4.7$.

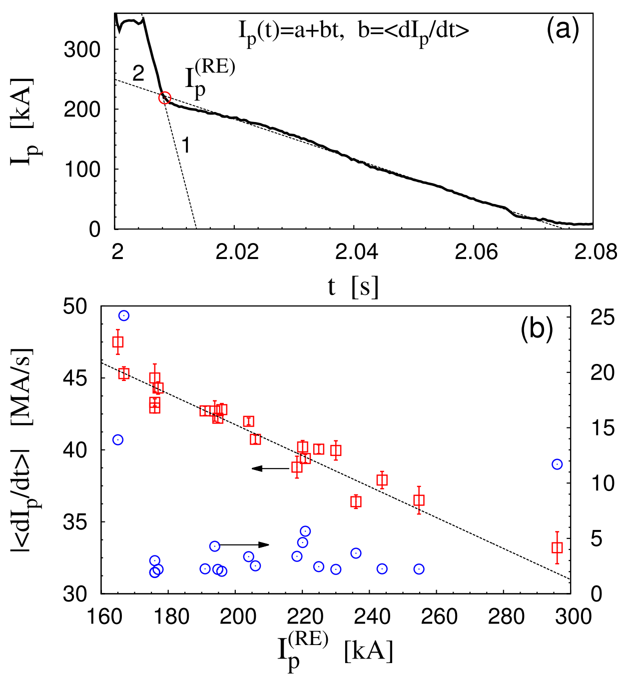

FIG. 2. (a) Typical time evolution of the plasma current with RE current. The average current decay rates $\langle dI_p/dt\rangle$ at the CQ and the RE plateau stages are determined by fitting with a linear function $I_p(t) = a + bt$. Symbol $\odot$ corresponds to the plasma current $I_p^{(RE)}$ at the initial stage of the RE plateau. (b) The decay rates $|\langle dI_p/dt\rangle|$ versus $I_p^{(RE)}$. Symbols $\boxdot$ (red) correspond to the CQ rate (lhs axis), and $\odot$ (blue)—the RE plateau (rhs axis).

ρ_c at the borders of region $\rho_1 < \rho < \rho_3$. For these discharges, the CQ rates $\langle|dI_p/dt|\rangle$ take highest or lowest values. The RE current decay rates of these discharges take the highest values. They have the shortest duration time of RE currents. One expects that 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 for the discharge with the highest $I_p^{(RE)}$ 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 a stochastic zone at the edge of the RE beam.

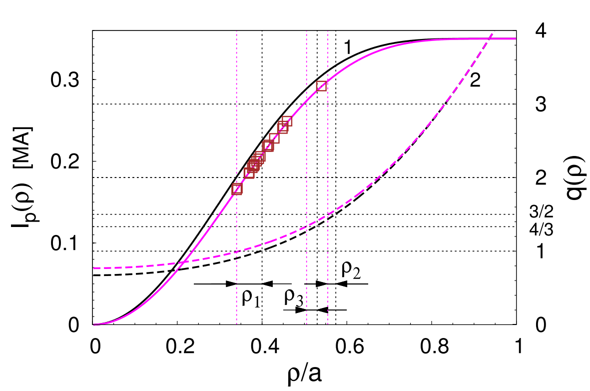

FIG. 3. Radial profile of the plasma current $I_p(\rho)$ (solid curves 1 on lhs axis) and the corresponding safety factor profile $q(\rho)$ (dashed curves 2 on rhs axis). The rectangular (red) dots correspond to the experimentally measured values of $I_p^{(RE)}$ for several TEXTOR discharges. The plasma parameters are $I_p = 350$ kA, $B_0 = 2.4$ T, $R_0 = 1.75$ m, and $a = 0.46$ m. The values of $q(0)$ are 0.7 (solid black curves) and 0.8 (dashed magenta curves), respectively. The radii ρ_1, ρ_2, and ρ_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.

The existence of the intact magnetic surface ρ_c between the $q = 1$ and $q = 4/3$ rational magnetic surfaces and its location depends on the level magnetic perturbation ϵ_{MHD} (more exactly on the spectrum B_{mn}). With increase of ϵ_{MHD}, the radius ρ_c shrinks and it can be broken at the certain critical perturbation level ϵ_{cr}. It leads to the total destruction of confinement of electrons and ions. This is in agreement with experimental observations of the existence of critical magnetic perturbations from which on runaway beams are not generated.[10]

The shrinkage of ρ_c with increasing the magnetic perturbation ϵ_{MHD} leads to the decrease of the RE current $I_p^{(RE)}$ since $I_p^{(RE)} \approx I_p(\rho_c)$. On the other hand, if one assumes that the plasma current decay is caused by the radial transport of particles in the stochastic magnetic field, the CQ rate dI_p/dt should be proportional to the square of the magnetic perturbation level ϵ_{MHD}, $|\langle dI_p/dt\rangle| \propto |\epsilon_{MHD}|^2$. Therefore, one expects that to the higher values of $|\langle dI_p/dt\rangle|$ correspond the lower values of the RE current $I_p^{(RE)}$. This expectation is in agreement with the experimental measured values of these quantities presented in Fig. 2(b).

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. One observes that the radiation is localized within a finite radial extent in the central plasma region.

The strong radial transport along the stochastic magnetic field lines causes the losses of heat and plasma particles from the stochastic zone. The TQ can be explained by the fact that the anomalously large heat transport in a stochastic magnetic field is mainly determined by the electron diffusion. The CQ is determined by the particle transport in a stochastic magnetic field and has an ambipolar nature. Using the collisional test particle transport model in a stochastic magnetic field,[24] we estimated the heat conductivity $\chi_r(\rho)$ and the ambipolar diffusion coefficient D_p of particles.

For typical magnetic perturbations and pre-disruption plasma temperatures (0.5–1.0 keV), the magnitude of $\chi_r(\rho)$ has the order of several 10^2 m^2/s. The characteristic heat diffusion time $\tau_H = a^2/2\chi_r$ is of the order of 10^{-4} s that agrees with the experimentally observed times. The quantitative analysis based on the numerical solution of the heat diffusion equation also gives similar values for τ_H.

The ambipolar particle transport in a stochastic magnetic field is strongly collisional due to the low plasma temperature (from 5 eV to 50 eV) after the TQ. At these plasma temperatures, the corresponding diffusion time $\tau_p = a^2/D_p$ of particles changes from 1 s to 0.3 s. Since the diffusion coefficient $D_p \propto B_{mn}^2$ and therefore $\tau_p \propto B_{mn}^{-2}$, then τ_p can be reduced to one order smaller value for a three times larger perturbation than in Fig. 1. This timescale is still much longer than the experimental values. However, the collisional model does not take into account the effect of the toroidal electric field. One expects that the acceleration of electrons and ions by the toroidal electric field increases the radial transport of particles. To include this effect in the collisional model, one can assume that the effective temperature of the plasma is

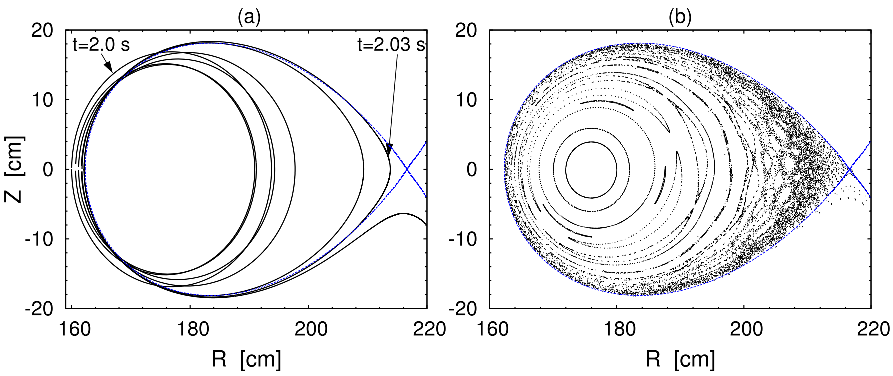

FIG. 4. (a) Evolution of a RE orbit in the (R, Z)-plane for the TEXTOR discharge No. 117527. (b) Poincaré section of RE orbit of energy 11.7 MeV in the (R, Z)-plane. The plasma current $I_p = 50$ kA.

higher than the measured one. The particle diffusion time τ_p at the effective temperature 2 keV is about 8×10^{-3} s. This timescale gives the average current decay rate $dI_p/dt \approx I_p/\tau_p = 0.35/(8.0 \times 10^{-3}) \approx 44.0$ MA/s which is order of the experimental measured one given in Fig. 2(b).

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 likes the ones in Fig. 1 leads to poloidally and toroidally localized heat and particle deposition patterns on the wall (see, e.g., Ref. 21) similar to those in ergodic divertor tokamaks (see, e.g., Ref. 18).

From the described scenario of plasma disruption, it follows that a typical runaway beam 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 ρ but also vary along the poloidal θ and the toroidal φ angles due to the presence of the $(m/n = 1/1)$ magnetic island. This agrees with the analysis of numerous disruptions in the JET tokamak.[25] One can assume that the radial profiles of the RE current density averaged along poloidal and toroidal angles are almost uniform. This gives 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.[26–31]

The toroidal electric field accelerates electrons to higher energies. With increasing electrons energy, their orbits drift outwardly[32,33] and eventually hit the wall. It is illustrated in Fig. 4(a). This effect may be one of mechanisms of slow RE current decay. Calculations show that the outward drift velocity v_{dr} is of the order of a few m/s for typical discharges in TEXTOR. The RE current decay rate dI_p/dt due to outward drift RE orbits can be roughly estimated as follows. This loss mechanism is mainly caused by the shrinkage of the beam radius a. The rate of such a shrinkage da/dt is of the order of the average outward velocity v_{dr}. Since $I_p \propto a^2$, we have $dI_p/dt \propto (2I_p/a)da/dt = (2I_p/a)v_{dr}$. 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 in the order of the experimentally measured average decay rate of the runaway current plotted in Fig. 2(b).

The effect of magnetic perturbation on RE beams depends on their safety factor profile q. The latter varies in the interval $[q(0) < 1,\, q(a)]$ with its edge value $q(a)$ less than 3/2 [or 4/3, 5/3]. Such a RE beam is relatively stable to the effect of magnetic perturbations. The single $m/n = 1/1$ mode does not create the stochastic layer at the beam edge for REs with energies up to several MeVs, since their drift surfaces are close to magnetic surfaces. With increasing the energy of electrons, the drift surfaces strongly deviate from magnetic ones and thus create the perturbation harmonics with higher mode numbers $m > 1$. The interactions of several resonance modes of perturbations may form the stochastic zone at the beam edge, which leads to fast RE losses as illustrated in Fig. 4(b). This process may explain the sudden RE current drop accompanied by magnetic activity and RE bursts observed in experiments (see, e.g., Refs. 4 and 9).

Based on the analysis of numerous experimental data obtained in the TEXTOR tokamak, we have proposed the mechanism of RE beam formation during the plasma disruption. The plasma disruption starts due to a large–scale magnetic stochasticity caused by nonlinearly excited of MHD modes with low (m, n) numbers ($m/n = 1/1$, 2/1, 3/2, 5/2,…). The RE beam is formed in the central plasma region confined by the intact magnetic surface. Its location depends on the safety factor profile $q(\rho)$ and the spectrum of MHD modes. In the cases of plasmas with the monotonic profile of $q(\rho)$ and at sufficiently small amplitude of the $m/n = 1/1$ mode, the most stable RE beams are formed by the intact magnetic surface located between the magnetic surface $q = 1$ and the closest low–order rational surface $q = m/n > 1$ ($q = 5/4$, $q = 4/3$, or $q = 3/2$).

This mechanism reproduces well the essential features of the measurements. Particularly, the TQ and the CQ are determined by the strong electron diffusion and ambipolar transport of particles in a stochastic magnetic field, respectively. The slow decay of the RE current is due to the outward drift of RE orbits induced by a toroidal electric field, and the spiky quick decay of REs is due to resonant interaction of high–energy REs with the $m/n = 1/1$ MHD mode. The effect of external resonant magnetic perturbations on low-energy electrons (up to 5–10 MeV) is weak and does not cause their loss. This is in agreement with the recent experiments in the TEXTOR tokamak.[34] The detailed description of the mechanism of RE formation and the evolution of RE current based on the analyses of experimental observations will be given in a separate publication.[35]

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[1]J. Wesson, R. Gill, M. Hugon, F. Schüller, J. Snipes, D. Ward, D. Bartlett, D. Campbell, P. Duperrex, A. Edwards *et al.*, Nucl. Fusion **29**, 641 (1989).
[2]R. D. Gill, Nucl. Fusion **33**, 1613 (1993).
[3]F. Schüller, Plasma Phys. Controlled Fusion **37**, A135 (1995).
[4]R. D. Gill, B. Alper, A. W. Edwards, L. C. Ingesson, M. F. Johnson, and D. J. Ward, Nucl. Fusion **40**, 163 (2000).
[5]G. Papp, T. Fülöp, T. Fehér, P. C. de Vries, V. Riccardo, C. Reux, M. Lehnen, V. Kiptily, V. V. Plyusnin, B. Alper *et al.*, Nucl. Fusion **53**, 123017 (2013).
[6]T. C. Hender, J. C. Wesley, J. Bialek, A. Bondeson, A. H. Boozer, R. J. Buttery, A. Garofalo, T. P. Goodman, R. S. Granetz, Y. Gribov *et al.*, Nucl. Fusion **47**, S128 (2007).
[7]E. Hollmann, N. Commaux, N. W. Eidietis, T. E. Evans, D. A. Humphreys, A. N. James, T. C. Jernigan, P. B. Parks, E. J. Strait, J. C. Wesley *et al.*, Phys. Plasmas **17**, 056117 (2010).
[8]M. Lehnen, A. Alonso, G. Arnoux, N. Baumgarten, S. A. Bozhenkov, S. Brezinsek, M. Brix, T. Eich, S. N. Gerasimov, A. Huber *et al.*, Nucl. Fusion **51**, 123010 (2011).
[9]M. Forster, K. Finken, T. Kudyakov, M. Lehnen, O. Willi, Y. Xu, L. Zeng, and the TEXTOR Team, Phys. Plasmas **19**, 092513 (2012).
[10]L. Zeng, H. R. Koslowski, Y. Liang, A. Lvovskiy, M. Lehnen, D. Nicolai, J. Pearson, M. Rack, H. Jaegers, K. H. Finken, K. Wongrach, and Y. Xu, Phys. Rev. Lett. **110**, 235003 (2013).
[11]K. Wongrach, K. Finken, S. S. Abdullaev, R. Koslowski, O. Willi, L. Zeng, and the TEXTOR Team, Nucl. Fusion **54**, 043011 (2014).
[12]Z. Y. Chen, W. Kim, Y. Yu, A. England, J. Yoo, S. Hahn, S. Yoon, Y. Lee, Y. K. Oh, J. Kwak, and M. Kwon, Plasma Phys. Controlled Fusion **55**, 035007 (2013).
[13]V. V. Plyusnin, V. Riccardo, R. Jaspers, B. Alper, V. Kiptily, J. Mlynar, S. Popovichev, E. de La Luna, F. Andersson, and JET EFDA Contributors, Nucl. Fusion **46**, 277 (2006).
[14]N. Commaux, L. R. Baylor, S. K. Combs, N. W. Eidietis, T. E. Evans, C. R. Foust, E. M. Hollmann, D. A. Humphreys, V. A. Izzo, A. N. James *et al.*, Nucl. Fusion **51**, 103001 (2011).
[15]E. M. Hollmann, M. E. Austin, J. A. Boedo, N. H. Brooks, N. Commaux, N. W. Eidietis, D. A. Humphreys, V. A. Izzo, A. N. James, T. C. Jernigan *et al.*, Nucl. Fusion **53**, 083004 (2013).
[16]L. P. Kadanoff, *Perspectives in Computational Science* (IEEE-CS and AIP, 2004), March/April p. 57.
[17]A. H. Boozer, Phys. Plasmas **19**, 058101 (2012).
[18]S. S. Abdullaev, *Magnetic Stochasticity in Magnetically Confined Fusion Plasmas* (Springer, Cham, 2014).
[19]B. B. Kadomtsev, Plasma Phys. Controlled Fusion **26**, 217 (1984).
[20]J. Wesson, *Tokamaks*, 3rd ed. (Clarendon Press, Oxford, 2004).
[21]S. E. Kruger, D. D. Schnak, and C. R. Sovinec, Phys. Plasmas **12**, 056113 (2005).
[22]V. A. Izzo, D. A. Humphreys, and M. Kornbluth, Plasma Phys. Controlled Fusion **54**, 095002 (2012).
[23]S. A. Bozhenkov, M. Lehnen, K. H. Finken, M. W. Jakubowski, R. C. Wolf, R. Jaspers, M. Kantor, O. V. Marchuk, E. Uzgel, G. VanWassenhove, O. Zimmermann, D. Reiter, and TEXTOR Team, Plasma Phys. Controlled Fusion **50**, 105007 (2008).
[24]S. S. Abdullaev, Phys. Plasmas **20**, 082507 (2013).
[25]S. N. Gerasimov, T. Hender, J. Morris, V. Riccardo, L. Zakharov, and JET EFDA Contributors, Nucl. Fusion **54**, 073009 (2014).
[26]H. Soltwisch, W. Stodiek, J. Manickam, and J. Schlüter, in *Proceedings of the 11th IAEA Conference on Plasma Physics and Controlled Fusion Research, Kyoto, 13–20 November, 1986* (IAEA, Vienna, 1987), Paper No. IAEA-CN-47/A-V-1, Vol. 1, pp. 263–273.
[27]M. Yamada, F. Livinton, N. Pomphrey, R. Budny, J. Manickam, and Y. Nagayama, Phys. Plasmas **1**, 3269 (1994).
[28]H. Soltwisch and H. R. Koslowski, Plasma Phys. Controlled Fusion **37**, 667 (1995).
[29]J. O'Rourke, Plasma Phys. Controlled Fusion **33**, 289 (1991).
[30]H. R. Koslowski, H. Soltwisch, and W. Stodiek, Plasma Phys. Controlled Fusion **38**, 271 (1996).
[31]H. Soltwisch and H. R. Koslowski, Plasma Phys. Controlled Fusion **39**, A341 (1997).
[32]X. Guan, H. Qin, and N. Fisch, Phys. Plasmas **17**, 092502 (2010).
[33]S. S. Abdullaev, Phys. Plasmas **22**, 030702 (2015).
[34]H. R. Koslowski, L. Zeng, M. Lehnen, A. Lvovskiy, K. Wongrach, and TEXTOR Team, in *Proceedings of the 41st EPS Conference on Plasma Physics*, Berlin, June 22–28, 2014, No. P5.028.
[35]S. S. Abdullaev *et al.*, "Mechanisms of plasma disruption and runaway electron losses in tokamaks," J. Plasma Phys. (to be published).