Editors' Suggestion

Runaway electrons during a coil quench in stellarators

Pavel Aleynikov,* Per Helander, and Håkan M. Smith
Max-Planck-Institut für Plasmaphysik, 17491 Greifswald, Germany

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It is shown that avalanches of runaway electrons can arise in stellarators, even if there is no net toroidal current in the plasma or the magnetic-field coils, if the current in the latter varies rapidly enough, e.g., due to a superconductor quench. In present-day devices such as W7-X, significant runaway generation is theoretically possible only at very low gas or plasma density in the vacuum vessel, e.g., between discharges, if a seed population of free electrons is present. In reactor-scale stellarators, more dangerous runaway generation may occur both between discharges and during low-density plasma operation. Since a radiation-induced seed population is necessarily present in an activated device, an accidental coil ramp-down could convert substantial magnetic energy into wall-damaging runaway currents. There is, however, much more time to mitigate such events than in tokamak disruptions.

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I. INTRODUCTION

Tokamak disruptions are often accompanied by the generation of "runaway" electrons, which are accelerated to relativistic energies by the toroidal electric field that is induced when the plasma current decays [1,2]. Runaway electrons have the potential to cause great damage and remain a serious safety concern for large tokamaks [3]. They have occasionally been observed in stellarators as a result of Ohmic heating [4] and during ramp-up or ramp-down of the helical coil currents [5–7]. Furthermore, a runaway-electron-suppression system had to be implemented at the ATF torsatron to avoid unacceptable levels of hard x-ray emission [8,9].

In a superconducting stellarator like W7-X, the coil currents are normally varied very slowly, but an exception occurs if the quench-detection system shuts down the coil currents because of a quench or for other reasons. In this case, the current is ramped down on a time scale of a few seconds [10,11]. Although no net toroidal current is then interrupted (since none of the coils extend around the torus in the toroidal direction), a toroidal electric field is nevertheless induced, because the coils create a poloidal magnetic flux, which decays during the quench. It is thus possible that runaway electrons could be created, which is the subject of the present work, where we discuss the generation, cross-field drift, and danger associated with such electrons.

II. MAGNETIC AND ELECTRIC FIELDS

The magnetic field can be written in Boozer coordinates as [12]

$$\boldsymbol{B} = \nabla\psi \times \nabla\theta + \nabla\varphi \times \nabla\chi,$$

where θ and φ are the poloidal and toroidal angle coordinates, $2\pi\psi$ is the toroidal magnetic flux, and $2\pi\chi$ is the poloidal flux. These fluxes are functions of space and time, but if the currents in all coils quench at the same rate and the plasma current is neglected, the geometry of the magnetic field remains constant, so that

$$\psi(\boldsymbol{r}, t) = f(t)\psi_0(\boldsymbol{r}),$$
$$\chi(\boldsymbol{r}, t) = f(t)\chi_0(\boldsymbol{r}),$$

and the Boozer angles are then time independent. For simplicity, we restrict our attention to the case of a low-β plasma without net toroidal current, so that most of the magnetic field is produced by the coils. A suitable, single-valued, vector potential is

$$\boldsymbol{A} = \psi\nabla\theta - \chi\nabla\varphi,$$

and the electric field induced by the current quench is

$$\boldsymbol{E} = -\partial\boldsymbol{A}/\partial t = -\dot{f}(t)(\psi_0\nabla\theta - \chi_0\nabla\varphi).$$

It is important to note that the gauge matters: $\boldsymbol{E}$ is not invariant under the transformation $\chi \to \chi + \text{const}$. To determine the correct value of the additive constant, we

*Contact author: pavel.aleynikov@ipp.mpg.de

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consider the somewhat idealized case where the coil currents are replaced by surface currents on some toroidal surface S surrounding the plasma. According to Faraday's law, the loop voltage

$$\oint_C \boldsymbol{E} \cdot \mathrm{d}\boldsymbol{r}$$

around a closed contour $C \subset S$ is equal to minus the time derivative of the linked magnetic flux. If we choose C so as to encircle the torus once toroidally, but not poloidally, then

$$\oint_C \boldsymbol{E} \cdot \mathrm{d}\boldsymbol{r} = 2\pi \dot{f}\,(t)\chi_0 = 0,$$

and we thus conclude that the additive constant should be chosen so that χ vanishes on S. The function $\chi(\psi)$ then equals the poloidal flux between the surface S and the flux surface labeled by ψ.

On the magnetic axis, where the enclosed toroidal flux vanishes, $\psi = 0$, the field becomes $\boldsymbol{E} = \dot{\chi}\nabla\varphi$ and the loop voltage

$$U = \oint \boldsymbol{E} \cdot \mathrm{d}\boldsymbol{r} = 2\pi\,\dot{\chi}.$$

As already noted, the right-hand side represents the time derivative of the poloidal flux, which can also be calculated from the mutual inductance M_i between the coils ($i = 1, \ldots, 5$) and the magnetic axis

$$U = \sum_i M_i N \dot{I}_i,$$

where I_i denotes the current in the i'th coil and N is the number of windings. In W7-X, the total poloidal flux is typically about

$$2\pi\,\chi = \sum_i M_i N I_i \simeq 2\ \mathrm{Tm}^2$$

at full field. If the ramp-down of coil currents occurs on the time scale of a few seconds, we thus expect a loop voltage of a few volts and an electric field,

$$E = \frac{U}{2\pi R} = \frac{\dot{\chi}}{R}, \tag{1}$$

of the order of 0.02 V/m. This field will quickly diffuse through the vacuum vessel, whose skin time is only a few milliseconds. For the same reason, very little of the magnetic energy will be dissipated Ohmically in the vacuum vessel, making this energy instead available to accelerate electrons.

III. ACCELERATION AND DECONFINEMENT OF ELECTRONS

Such electrons would not be of any great concern if they were sufficiently well confined by the magnetic field. Since runaway electrons have such small pitch angles that they are not trapped in local magnetic wells, their orbits deviate very little from the magnetic field, if this is kept constant [13]; however, in a decaying magnetic field, the electrons will move toward the wall because of the $\boldsymbol{E} \times \boldsymbol{B}$ drift. In a vacuum magnetic field, $\boldsymbol{B} = G\nabla\varphi$,

$$\frac{\boldsymbol{E} \times \boldsymbol{B}}{B^2} \cdot \nabla\psi = \frac{G(\nabla\varphi \times \nabla\psi) \cdot \boldsymbol{E}}{B^2} = -\frac{\partial\psi}{\partial t},$$

implying that ψ is a constant of the motion,

$$\frac{\partial\psi}{\partial t} + \left(\frac{v_\parallel \boldsymbol{B}}{B} + \frac{\boldsymbol{E} \times \boldsymbol{B}}{B^2}\right) \cdot \nabla\psi = 0.$$

As the field decays, the electrons on each magnetic surface will thus move radially outward in such a way that the enclosed toroidal flux remains constant. Roughly speaking, if the field strength drops by a factor of 2, their distance from the magnetic axis will grow by a factor $\sqrt{2}$. All free electrons will thus eventually leave the confinement region as $B \to 0$.

This analysis is only strictly valid in the limit of small plasma pressure and current, and it will cease to hold if so many electrons are accelerated that the resulting plasma current substantially modifies the magnetic field. Qualitatively, however, the $\mathbf{E} \times \mathbf{B}$ drift induced by the decaying coil currents will still convect particles outward.

The characteristic time required for accelerating an electron to relativistic energies is

$$t_\mathrm{c} = \frac{m_e c}{eE} \sim 100\ \mathrm{ms} \tag{2}$$

for $E = 0.02$ V/m. This is much shorter than the quench time, so we conclude that practically all runaway electrons will reach MeV energies. The key question is whether such freely accelerating electrons will multiply through impact ionization and thus set off an avalanche.

IV. RUNAWAY AVALANCHE

A. Avalanche in a neutral gas

In both a plasma and a neutral gas, the friction force F_s experienced by a charged particle decreases over a broad energy range. Consequently, electrons run away if their energy exceeds a threshold at which the accelerating force exceeds friction, $eE > F_s$.

The friction on an electron in a neutral gas was derived by Bethe [14],

$$F_s = \frac{4\pi r_0^2 n_e m_e c^2}{\beta^2}\left(\ln\frac{m_e c^2(\gamma-1)\sqrt{\gamma+1}}{\sqrt{2}I} - \left(\frac{2}{\gamma}-\frac{1}{\gamma^2}\right)\ln 2 + \frac{1}{\gamma^2}+\frac{(\gamma-1)^2}{8\gamma^2}\right), \quad (3)$$

where r_0 is the classical electron radius, I is the characteristic ionization potential, n_e is the total (bound) electron density, $\beta \equiv v/c$, and $\gamma = (1-\beta^2)^{-1/2}$. Equation (3) has a maximum near $\gamma_{\max} \approx 2.72I/(m_e c^2) + 1$, where the friction force is

$$F_s^{\max} = \frac{2\pi r_0^2 n_e m_e^2 c^4}{2.72I}.$$

Between discharges in W7-X, the pressure is about 10^{-5} Pa; thus, the hydrogen molecule density $n_n \sim 2.4 \times 10^{15}$ m^{-3} ($n_e = 2n_n$). Accordingly, all free electrons will run away if the electric field exceeds $E_{\max} \approx 1.5$ mV/m, which is much smaller than the field induced during a quench. The force given by Eq. (3) is shown in Fig. 1.

The particle energy loss described by Eq. (3) arises primarily from ionization processes. Secondary electrons produced in these collisions may themselves be accelerated by the electric field, leading to a runaway avalanche [15,16]. Although the ionization cross section decreases with increasing energy, the total number of secondary electrons generated during the acceleration to nonrelativistic speeds is negligible compared to that produced by a relativistic electron [17], as shown in Appendix B.

In the regime where all particles run away, the avalanche rate is simply given by the ultrarelativistic electron ionization rate. The Bethe cross section for the total ionization in the high-energy limit has the form [18]

$$\sigma_i = \frac{4\pi r_0^2}{\alpha^2\beta^2}\left(M^2\left(\ln(\gamma^2-1)-\beta^2\right)+C\right),$$

where M^2 and C are constants related to the continuum oscillator strength and α is the fine structure constant. For atomic hydrogen, $M^2 = Q/2$ and $C = 2 - Q - M^2\ln\alpha^2$, with $Q = 0.5668$. The resulting avalanche rate when $\beta \to 1$ is

$$\Gamma_n = \frac{4\pi r_0^2 n_e c}{2}\frac{m_e c^2}{I}\frac{Q}{2}\left(\ln\frac{m_e c^2(\gamma^2-1)}{2I}+\frac{4}{Q}-3\right), \quad (4)$$

which represents the maximum possible rate. Taking $n_e = 2n_n \sim 4.8 \times 10^{15}$ m^{-3} and the logarithm to be equal to 15, we obtain $\Gamma_n \approx 15$ s^{-1}, resulting in a net amplification

$$\frac{N_r(t)}{N_r(0)} = \exp\left(\int_0^t \Gamma_n dt\right),$$

which amounts to several tens of orders of magnitude within a few seconds when $eE > F_s^{\max}$. Significant runaway growth via avalanching is thus plausible in W7-X.

When the electric field decreases below $E_{\max}$, only electrons born with sufficiently high energy can run away. The minimum friction force, which occurs near $\gamma \approx 10$ [see Eq. (3)], defines the critical electric field $E_c = \min(F_s)/e$ below which no acceleration is possible [19],

$$eE_c = 4\pi r_0^2 n_e m_e c^2 \ln\Lambda,$$

where $\ln\Lambda$ is the first logarithmic term in Eq. (3). This definition is equivalent to the classical threshold electric field for "runaway breakdown" in a fully ionized plasma, provided that the Coulomb logarithm $\ln\Lambda$ is adapted to describe free-free collisions. To do so, the ionization potential I in Eq. (3) needs to be replaced by $\hbar\omega_p\gamma$, where ω_p is the plasma frequency [1].

For electric fields such that $E_{\max} > E > E_c$, only particles born with energy above the critical energy

$$\varepsilon_c \approx \frac{1}{2}\frac{E_c}{E}m_e c^2$$

can run away. The rate at which runaway electrons are produced in a tokamak in this regime was calculated by Rosenbluth and Putvinski [20]. As shown in Appendix A, their calculation extends to the stellarator case with unimportant modifications and results in the

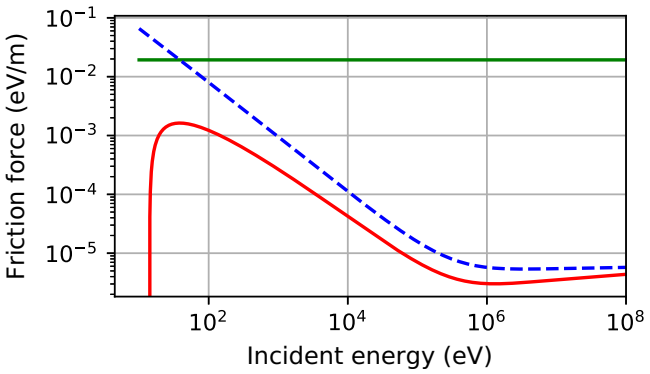


FIG. 1. Friction force [Eq. (3)] acting on an electron in neutral hydrogen with atomic density $n_e \sim 4.8 \times 10^{15}$ m^{-3} (red) and in a fully ionized plasma (dashed) with the same density. The green horizontal line indicates the maximum electric field expected during W7-X coil quenches.

growth rate

$$\Gamma \simeq \sqrt{\frac{\epsilon\pi}{3(Z+5)}}\frac{E/E_c - 1}{\tau \ln\Lambda}, \tag{5}$$

where the relativistic collision time is $\tau = m_e c/(eE_c) = (4\pi r_0^2 n_e c \ln\Lambda)^{-1}$, the factor $\epsilon^{1/2}$ accounts for the fraction of trapped particles, and Z is the effective ion charge experienced by relativistic electrons.

In a weakly ionized plasma, Eq. (5) should be used with caution since collisions of secondary particles with neutrals are less efficient than collisions in a fully ionized plasma. The avalanche rate in a neutral gas is therefore expected to be higher. References [21,22] suggested a modification to Eq. (5) that approximately accounts for the effect of reduced collisions. The suggested recipe amounts to using the critical runaway energy ε_c as the characteristic energy when evaluating $\ln\Lambda$ in the case of a neutral gas.

Equations (4) and (5) can be combined into an interpolation formula (assuming $E/E_c \gg 1$) between the two avalanche regimes,

$$\Gamma_c \simeq 4\pi r_0^2 n_e c \left(\frac{1}{a}\frac{E_c}{E} + \frac{1}{q}\frac{I}{m_e c^2}\right)^{-1}, \tag{6}$$

where $\ln\Lambda$ in E_c is the logarithmic term in the parentheses of Eq. (3), with γ evaluated from the force-balance relation $eE = F_s$. Furthermore, $a = \sqrt{\epsilon\pi/(3(Z+5))}$ and

$$q = \frac{Q}{4}\left(\ln\frac{m_e c^2(\gamma^2-1)}{2I} + \frac{4}{Q} - 3\right) \approx 2.5,$$

taking $\gamma \approx 10$ for primary electrons. Figure 2 shows the resulting avalanche rate as a function of the electric field, taking $Z = 1$ and $\epsilon = 1$ [23].

Taking $E(t) = E_0 e^{-t/t_{\rm quench}}$ with $E_0 t_{\rm quench} = \chi_0/R$ and $t_{\rm quench} = 3$ s (see Ref. [11]), we obtain $N_{\rm r}(\infty)/N_{\rm r}(0) \approx 400$, suggesting that a small seed population cannot give rise to a substantial runaway current. It should, however, be emphasized that Eq. (6) is merely an interpolation, with the electric field of interest falling between rigorously described limits (indicated by a thick red line in Fig. 2). This estimate assumes that only a small fraction of secondary electrons become runaways even when $E > eF_s^{\rm max}$ (shadowed region in Fig. 2). Using instead Eq. (4) for $eE > F_s^{\rm max}$ and Eq. (5) for $eE < F_s^{\rm max}$ results in an amplification factor $N_{\rm r}(\infty)/N_{\rm r}(0) \approx 10^{45}$, which represents the most "pessimistic" limit. The possibility of a strong runaway avalanche therefore cannot definitely be ruled out. More accurate kinetic modeling is required to refine the quantitative prediction.

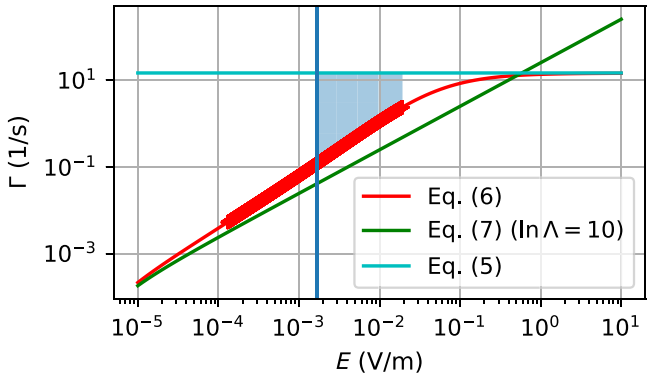


FIG. 2. Avalanche rate as a function of electric field in neutral hydrogen (same density as in Fig. 1). The vertical line indicates $E_{\rm max} \equiv F_s^{\rm max}/e$.

B. Runaway breakdown and plasma operation

Equation (6) is valid in a neutral gas, but as runaway breakdown progresses and the gas is ionized, the secondary electrons experience more efficient collisions with the ionized plasma (see Fig. 1). Therefore, the ionization level needs to be tracked to estimate the avalanche rate more accurately. The seed runaway density $N_r(0)$ then plays a decisive role.

In the limit of a fully ionized plasma, the avalanche rate is given by Eq. (5). Substituting Eq. (1) for the electric field in Eq. (5) and integrating over time, we see that the final number of runaway electrons does not depend on the quench time, nor on any other aspect of the quench history if $E/E_c \gg 1$ [1,20],

$$\frac{N_{\rm r}(\infty)}{N_{\rm r}(0)} = \exp\left(\sqrt{\frac{\epsilon\pi}{3(Z+5)}}\frac{e|\chi_0|}{m_e c R \ln\Lambda}Q\right), \tag{7}$$

where for exponential E-field decay

$$Q = 1 - \frac{E_c}{E_0}\left(1 - \ln\frac{E_c}{E_0}\right) \simeq 1$$

if $E/E_c \gg 1$. This estimate gives $N_{\rm r}(\infty)/N_{\rm r}(0) \sim 1$, which implies that no significant runaway population is generated if a quench occurs during plasma operation in W7-X.

Finally, we note that no avalanche at all is possible if $E_0 < E_{\rm c}$, or

$$E_c t_{\rm quench} > |\chi_0|/R. \tag{8}$$

For W7-X, this condition translates to

$$\frac{n_e}{10^{20}\ {\rm m}^{-3}} \times \frac{t_{\rm quench}}{1\ {\rm s}} > 1.$$

Runaway avalanches are thus in principle impossible in plasmas with high enough density.

V. CONCLUSIONS AND OUTLOOK

The disruption mitigation system is an essential part of future tokamak reactors. In stellarators, the quench detection system aims to prevent coil damage by rapidly ramping down the coil currents. We show, however, that runaway electrons can be generated in a stellarator if the currents in the coils are ramped down quickly, even if there is no net toroidal current in the coils or the plasma. Although particles circulating around the torus are generally well confined in stellarators, during coil-current ramp-down, electrons drift radially outward as the magnetic field decays, ultimately producing hard x-rays upon wall impact and perhaps damaging the wall.

In W7-X, the avalanche amplification factor is of order unity during routine plasma operation, implying negligible secondary multiplication and thus little hazard from runaways in this regime. In contrast, if a quench were to occur between discharges, when the gas pressure in the vacuum vessel were low, the avalanche factor would be very large, so the final runaway population would depend critically on the initial seed $N_r(0)$. In principle, no significant seed is expected in the absence of radiation-related sources between discharges, but the relatively large potential avalanche gain makes it difficult to entirely rule out substantial runaway production. The existence of a "weak" seed—involving only a few fast electrons—remains an open question. A practical mitigation strategy could be to maintain a sufficiently high neutral gas pressure (in W7-X above 2×10^{-6} mbar), such that $eE_0 < F_s$ for electrons with energies below 100 eV. Since the spontaneous appearance of higher-energy electrons is unlikely, this pressure threshold should effectively suppress runaway initiation.

The situation is markedly different in a reactor-scale stellarator, as follows.

(a) The ratio χ_0/R is expected to be an order of magnitude larger than in W7-X [24]. For example, in a reactor-size device [25] with $R \approx 12.5$ m, minor radius $a \approx 1.25$ m, and $B \approx 9$ T, one can expect $\chi_0/R \approx 0.5$ T m. The conservative estimate in Eq. (6) then yields $N_r(\infty)/N_r(0) \sim 10^{25}$ for interdischarge quenches at a gas density of 10^{16} m^{-3} and quench duration t_{quench} in the range 3–10 s.

(b) With higher χ_0/R, a quench during low-density operation produces substantial avalanche amplification. Indeed, Eq. (7) gives $N_r(\infty)/N_r(0) \sim 10^5$ for $n_e = 10^{19}$ m^{-3}. Note that Dreicer primary generation can be efficient in these conditions: it becomes significant when $E/E_D > 0.02$, where the Dreicer field is $E_D \equiv E_c m_e c^2/T_e$. For $t_{\text{quench}} = 10$ s, giving $E_0 \approx 0.05$ V m^{-1}, this criterion is met, e.g., for a $T_e \sim 1$ keV plasma.

(c) In addition to the Dreicer primary generation mechanism, sources such as Compton scattering of γ radiation from activated walls and tritium decay within the plasma provide a seed population during and between discharges.

Under these circumstances, substantial runaway generation becomes likely if the coil currents are rapidly ramped down and the plasma or neutral-gas density is below a critical level. The flux of γ-rays between plasma discharges depends on the activation history and the structural composition of the vessel wall. It is expected to be a few orders of magnitude lower than that during plasma operation, where it can be estimated to be up to the level of the neutron flux, i.e., $\Gamma_\gamma \leq P/(S\Delta E)$, where P and S denote the reactor's power and surface area, respectively, and $\Delta E = 17.6$ MeV. The Compton-scattering cross section depends on the γ-ray energy but can be crudely approximated by the Thomson cross section $\sigma_\gamma = 8\pi r_e^2/3$. The resulting seed-production rate is then $\dot{n}_{\text{seed}} = n_e \Gamma_\gamma \sigma_\gamma \approx 10^6$ m^{-3} s^{-1} for $P/S = 0.5$ MW/m^2 and $n_e = 10^{16}$ m^{-3}. The energies of these Compton-produced seed electrons can reach the MeV range, making them capable of runaway acceleration in any field satisfying $E > E_c$.

Rapid ramp-down of coil current is not the only mechanism for fast-particle generation in a stellarator. Fast electrons have also been observed in W7-X experiments during low-density electron cyclotron resonance heating (ECRH) operation and during ECRH ramp-ups [26]; however, runaway-electron generation in those scenarios is less efficient than in the coil ramp-down case and is not expected to have detrimental effects.

These considerations suggest that it would be wise to take precautions to avoid runaway-electron generation in future large stellarators. A superconducting-coil quench is a few orders of magnitude slower than a tokamak disruption, and the problem of runaway-electron mitigation should be correspondingly simpler. There is, for instance, ample time to inject material into the torus to stop the runaways, but some form of intervention is likely to be necessary.

DATA AVAILABILITY

The data that support the findings of this article are openly available [27].

APPENDIX A: KINETIC THEORY

The drift kinetic equation for relativistic electrons is

$$\frac{\partial f}{\partial t} + \dot{\boldsymbol{r}} \cdot \nabla f + \dot{p}_\parallel \frac{\partial f}{\partial p_\parallel} = C(f) + S,$$

where S encapsulates the effect of collisions at close range and the velocity-space coordinates have been chosen to be

$\mu = p_\perp^2/(2B)$ and $p_\parallel = \gamma v_\parallel$, with

$$\gamma = \frac{1}{\sqrt{1 - v^2/c^2}} = \sqrt{1 + \frac{1}{c^2}\left(2\mu B(\boldsymbol{r}) + p_\parallel^2\right)}.$$

The equations of motion are [28]

$$\dot{\boldsymbol{r}} = \frac{p_\parallel \boldsymbol{B}^*}{\gamma B_\parallel^*} + \frac{\boldsymbol{E}^* \times \boldsymbol{b}}{B_\parallel^*},$$
$$\dot{p}_\parallel = \frac{q\boldsymbol{B}^* \cdot \boldsymbol{E}^*}{mB_\parallel^*},$$
$$\boldsymbol{B}^* = \boldsymbol{B} + \frac{p_\parallel}{q}\nabla \times \boldsymbol{b},$$
$$\boldsymbol{E}^* = \boldsymbol{E} - \frac{m}{q}\left(c^2\nabla\gamma - p_\parallel \frac{\partial \boldsymbol{b}}{\partial t}\right),$$

where $q = -e$ denotes the charge, and the subscript $\parallel$ indicates the component in the direction of $\boldsymbol{b} = \boldsymbol{B}/B$. The kinetic energy $\epsilon = p_\parallel^2/2 + \mu B$ evolves according to

$$\dot{\epsilon} = \frac{qp_\parallel E_\parallel}{m} + \mu\frac{\partial B}{\partial t}.$$

The last term, which describes the conversion of magnetic to kinetic energy, is relatively small,

$$\mu\frac{\partial B}{\partial t}\Big/\frac{qp_\parallel E_\parallel}{m} \sim \frac{\gamma^2 v_\perp^2}{B}\,|\nabla \times \boldsymbol{E}|\Big/\frac{q\gamma v_\parallel E_\parallel}{m} \sim \frac{v_\perp \rho}{v_\parallel R} \ll 1,$$

where $\rho = \gamma m v_\perp/(qB)$ denotes the gyroradius. To make contact with Ref. [20], we change velocity-space variables to (p, λ), where $p^2 = 2\epsilon$ and $\lambda = 2\mu/p^2$. Then, $\dot{p} = \dot{\epsilon}/p$ and $\dot{\lambda} = -2\lambda\dot{\epsilon}/p^2$, and the drift kinetic equation becomes

$$\frac{\partial f}{\partial t} + \dot{\psi}\frac{\partial f}{\partial \psi} + \dot{\theta}\frac{\partial f}{\partial \theta} + \dot{\varphi}\frac{\partial f}{\partial \varphi} + \frac{q\xi E_\parallel}{m}\left[\frac{1}{p^2}\frac{\partial(p^2 f)}{\partial p} - \frac{2}{p}\frac{\partial(\lambda f)}{\partial \lambda}\right] = C(f) + S, \quad \text{(A1)}$$

where $\xi = p_\parallel/p = \pm\sqrt{1 - \lambda B}$.

To solve this equation, we expand the distribution function in the smallness of the gyroradius and the electric field, $f = f_0 + f_1 + \cdots$. To lowest order in $\rho/R \ll 1$, the orbits follow field lines, i.e., $\dot{\psi} = 0$ and

$$\dot{\theta} = \frac{v_\parallel \boldsymbol{B} \cdot \nabla\theta}{B} = \frac{\iota v_\parallel B}{G} = \iota\dot{\varphi},$$

and the distribution function is thus a function of (ψ, p, λ). For circulating particles, the parallel-streaming terms in (A1) are annihilated by the operator

$$\left\langle \frac{B}{\xi}(\cdots)\right\rangle,$$

and so is the next-order radial drift, $\langle B\dot{\psi}/\xi\rangle = 0$, where angular brackets denote the usual flux-surface average [12]. The resulting equation,

$$\left\langle\frac{B}{\xi}\right\rangle\frac{\partial f_0}{\partial t} + \frac{q\langle E_\parallel B\rangle}{m}\left[\frac{1}{p^2}\frac{\partial(p^2 f_0)}{\partial p} - \frac{2}{p}\frac{\partial(\lambda f_0)}{\partial \lambda}\right] = \left\langle\frac{B}{\xi}[C(f_0) + S]\right\rangle,$$

coincides with that solved in Ref. [20]. The distribution function and the growth rate of the avalanche are thus essentially the same as in a tokamak, notwithstanding nonaxisymmetry and the radial drift of the electrons due to the decaying magnetic field.

APPENDIX B: IONIZATION DURING INITIAL ACCELERATION

We can estimate the number of secondary electrons during initial acceleration to nonrelativistic energies by noting that the expected number of ionization events caused by an accelerating electron is equal to

$$\alpha = \int \sigma n_{\mathrm{n}} v\,\mathrm{d}t = \int \sigma n_{\mathrm{n}} v\frac{\mathrm{d}w}{\dot{w}},$$

where w denotes the kinetic energy of the electron, n_{n} is the neutral density, and $\sigma(w)$ is the ionization cross section. Since $\dot{w} = eEv$, we thus have

$$\alpha = \frac{n_{\mathrm{n}}}{eE}\int_0^{w_{\max}} \sigma\,\mathrm{d}w, \quad \text{(B1)}$$

where the integral is taken from 0 to the maximum energy the electron reaches during the time period in question. In the nonrelativistic Born approximation,

$$\sigma \sim \frac{\sigma_0 w_0}{w}, \quad w \to \infty,$$

where w_0 is some reference energy. The integral therefore diverges at as $w_{\max} \to \infty$, but only logarithmically. If $w_{\max}/w_0 \gg 1$, the main contribution comes from high energies, and we thus obtain the estimate

$$\alpha \simeq \frac{n_{\mathrm{n}}\sigma_0 w_0}{eE}\ln\frac{w_{\max}}{w_0}.$$

Between discharges in W7-X, the pressure in the plasma vacuum vessel is about 10^{-7} mbar $= 10^{-5}$ Pa, and the hydrogen molecule density is thus $n_{\mathrm{n}} \sim 2.4 \times 10^{15}$ m^{-3}.

The ionization cross section peaks at about 70 eV, where it reaches $\sigma_{\max} \simeq 10^{-20}$ m^2 and has fallen to about a quarter of this value at 1 keV. It thus appears that the nonrelativistic ionization probability (the expected number of ionization events caused by an electron as it is accelerated to subrelativistic energies) is only of the order of 1%–10%. Nonrelativistic electrons are thus not expected to produce breakdown.

[1] B. N. Breizman, P. Aleynikov, E. M. Hollmann, and M. Lehnen, Physics of runaway electrons in tokamaks, Nucl. Fusion **59**, 083001 (2019).

[2] P. Helander, L.-G. Eriksson, and F. Andersson, Runaway acceleration during magnetic reconnection in tokamaks, Plasma Phys. Controlled Fusion **44**, B247 (2002).

[3] M. Salewski, *et al.*, Energetic particle physics: Chapter 7 of the special issue: On the path to tokamak burning plasma operation, Nucl. Fusion **65**, 043002 (2025).

[4] W. Bernstein, F. Chen, M. Heald, and A. Kranz, Runaway electrons and co-operative phenomena in B-1 stellarator discharges, J. Nucl. Energy (1954) **7**, 276 (1958).

[5] A. C. England, G. L. Bell, R. H. Fowler, J. C. Glowienka, J. H. Harris, D. K. Lee, M. Murakami, G. H. Neilson, D. A. Rasmussen, J. A. Rome, M. J. Saltmarsh, and J. B. Wilgen, Runaway electron studies in the ATF torsatron, Phys. Fluids B: Plasma Phys. **3**, 1671 (1991).

[6] L. Rodríguez-Rodrigo, A. López-Fraguas, and A. Gabriel, Runaway electrons control during ramp-up and ramp-down of the magnetic fields in the coils of the TJ-II flexible heliac (abstract), Rev. Sci. Instrum. **70**, 645 (1999).

[7] V. Moiseenko, V. Korovin, I. Tarasov, M. Tarasov, D. Sitnikov, I. Garkusha, N. Zamanov, M. Lytova, R. Pavlichenko, A. Kulaga, *et al.*, The effect of an electrostatic field on runaway electrons in the Uragan-3M stellarator, Tech. Phys. Lett. **40**, 669 (2014).

[8] D. A. Rasmussen, A. C. England, C. C. Eberle, W. R. Devan, J. H. Harris, T. C. Jernigan, R. R. Kindsfather, R. N. Morris, M. Murakami, G. H. Neilson, J. A. Rome, M. J. Saltmarsh, R. B. Wysor, O. Motojima, H. Okada, and K. Uo, Suppression of x-rays generated by runaway electrons in ATF, in *Proceedings of the U.S.–Japan Heliotron-Stellarator Workshop* (Oak Ridge National Lab., TN, USA; Kyoto Univ., Japan, 1987).

[9] A. C. England, W. R. DeVan, C. C. Eberle, R. H. Fowler, W. A. Gabbard, J. C. Glowienka, J. H. Harris, G. R. Haste, R. R. Kindsfather, and R. N. Morris, *Runaway studies in the ATF (Advanced Toroidal Facility) torsatron*, Tech. Rep. (Oak Ridge National Lab., TN, USA, 1988).

[10] D. Birus, M. Schneider, T. Rummel, and M. Fricke, The quench detection system of Wendelstein 7-X, Fusion Eng. Des. **86**, 1566 (2011). Proceedings of the 26th Symposium of Fusion Technology (SOFT-26).

[11] K. Risse, T. Rummel, T. Mönnich, F. Füllenbach, and H.-S. Bosch, Updates on protection system for Wendelstein 7-X superconducting magnets, Fusion Eng. Des. **146**, 910 (2019). sI:SOFT-30.

[12] P. Helander, Theory of plasma confinement in non-axisymmetric magnetic fields, Rep. Progr. Phys. **77**, 087001 (2014).

[13] The radial orbit excursion is of the order of $\Delta r \sim \gamma m_e c/(eB)$, which is about 1 cm for a 10-MeV electron in W7-X. Here, γ denotes the Lorentz factor, m_e is the electron rest mass, and c is the speed of light.

[14] H. Bethe, Bremsformel für elektronen relativistischer geschwindigkeit, Z. Phys. **76**, 293 (1932).

[15] Y. A. Sokolov, Multiplication of accelerated electrons in a tokamak, JETP Lett. **29**, 218 (1979).

[16] A. Gurevich, G. Milikh, and R. Roussel-Dupre, Runaway electron mechanism of air breakdown and preconditioning during a thunderstorm, Phys. Lett. A **165**, 463 (1992).

[17] R. Jayakumar, H. Fleischmann, and S. Zweben, Collisional avalanche exponentiation of runaway electrons in electrified plasmas, Phys. Lett. A **172**, 447 (1993).

[18] Y.-K. Kim, J. P. Santos, and F. Parente, Extension of the binary-encounter-dipole model to relativistic incident electrons, Phys. Rev. A **62**, 052710 (2000).

[19] J. Connor and R. Hastie, Relativistic limitations on runaway electrons, Nucl. Fusion **15**, 415 (1975).

[20] M. N. Rosenbluth and S. V. Putvinski, Theory for avalanche of runaway electrons in tokamaks, Nucl. Fusion **37**, 1355 (1997).

[21] J. R. Martín-Solís, A. Loarte, and M. Lehnen, Runaway electron dynamics in tokamak plasmas with high impurity content, Phys. Plasmas **22**, 092512 (2015).

[22] L. Hesslow, O. Embréus, O. Vallhagen, and T. Fülöp, Influence of massive material injection on avalanche runaway generation during tokamak disruptions, Nucl. Fusion **59**, 084004 (2019).

[23] Note that the trapped particle fraction can be significant in stellarators, even on the magnetic axis. In W7-X, the trapped fraction on axis typically ranges from 0.15 to 0.3.

[24] F. Schauer, K. Egorov, and V. Bykov, HELIAS 5-B magnet system structure and maintenance concept, Fusion Eng. Des. **88**, 1619 (2013). Proceedings of the 27th Symposium On Fusion Technology (SOFT-27); Liège, Belgium, September 24-28, 2012.

[25] C. Hegna, D. Anderson, E. Andrew, A. Ayilaran, A. Bader, T. Bohm, K. C. Mata, J. Canik, L. Carbajal, A. Cerfon, *et al.*, The infinity two fusion pilot plant baseline plasma physics design, J. Plasma Phys. **91**, E76 (2025).

[26] C. Killer, P. Aleynikov, C. Biedermann, C. Dhard, P. Drews, Y. Gao, O. Grulke, M. Jakubowski, A. Knieps, G. Kocsis, D. Naujoks, A. Puig Sitjes, M. Spolaore, T. Stange, T. Szepesi, A. Tancetti, and H. Thomsen, Observation of non-thermal electrons outside the sol in the Wendelstein 7-X stellarator, Nucl. Mater. Energy **33**, 101274 (2022).

[27] P. Aleynikov, Data to support Figures 1 and 2 in "Runaway electrons during a coil quench in stellarators", https://doi.org/10.5281/zenodo.18196420 (2026).

[28] J. R. Cary and A. J. Brizard, Hamiltonian theory of guiding-center motion, Rev. Mod. Phys. **81**, 693 (2009).