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COLLISIONAL AVALANCHE EXPONENTIATION OF RUN-AWAY ELECTRONS IN
ELECTRIFIED PLASMAS

BY

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Collisional Avalanche Exponentiation of Run-Away Electrons in Electrified Plasmas

by

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Abstract

In contrast to earlier expectations, it is estimated that generation of runaway electrons from close collisions of existing runaways with cold plasma electrons can be significant even for small electric fields, whenever runaways can gain energies of about 20 MeV or more. In that case, the runaway population will grow exponentially with the energy spectrum showing an exponential decrease towards higher energies. Energy gains of the required magnitude may occur in large Tokamak devices as well as in cosmic-ray generation.

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1. Introduction

Super-thermal "runaway" electrons extracted from the thermal velocity distribution of a plasma by an electric field, E , and then accelerated to energies far exceeding the thermal level have been observed in a wide variety of experiments (Ref. 1). Well known examples include essentially all Tokamak (Ref. 2) and a number of Stellarator (Ref. 3) devices as well as earlier plasma-betatron experiments (Ref. 4). In some of these experiments (Refs. 2,5), runaway electrons were found to carry a large fraction of the total current ranging up to mega-amperes with particle energies up to many MeV's. On the positive side, such runaway "beams" have been credited with improving the confinement quality of some tokamak discharges (Ref. 6), and - in plasma betatrons - they have been expected to strongly expand the current capabilities of normal betatrons (Ref. 7); in addition, the confinement of runaway electrons may provide a good diagnostic for the electromagnetic turbulence spectra existing in such fusion-type plasmas (Ref. 8). On the other hand, these electrons tend to distort the X-ray and synchrotron diagnostic of high-temperature plasmas (Refs. 9, 10) and may enhance their wall-arcing (Ref. 11). Also, due to their sizable total energy content and their concentrated loss to relatively small parts of the vacuum wall, these runaways have been found to cause quite serious damages to some fusion devices (Ref. 12); and, the potential for such damages (Ref. 13) may complicate the design of next-generation high-current fusion machines. Last-but-not-least, the generation of runaways also may be of importance in astrophysical plasmas by providing a primary source for the generation of cosmic rays (Ref. 14).

Existing theoretical treatments of runaway generation (Refs. 15 and 16) were based mainly on collisional-statistical analyses of the electronic velocity distribution in the vicinity of a critical electron energy, W_C , at which the accelerating electric force on the electrons just is balanced by the effective net drag resulting from Coulomb collisions of these electrons with the thermal plasma. Beyond this energy limit, - due to the $1/W^2$ dependence of the Coulomb drag - the runaway electrons very quickly will incur essentially-free-fall acceleration by the applied electric field.

Following the derivations of Parail and Pogutse (Ref. 16) taking account of electron-electron as well as electron-ion collisions, the critical energy for a pure hydrogen plasma (i.e., $Z_{\text{eff}}=1$) is given by

$$W_C/e = (e^2/4\pi\epsilon_0^2) (n_e/E) \ln\Lambda = (KT_e/e)(E_C/E) \quad (1)$$

with electron temperature T_e and the "critical field" E_C defined by

$$E_C = m_e f_{ee} v_t = (e^3/4\pi KT_e \epsilon_0^2) n_e \ln\Lambda \quad (2)$$

with m_e and $v_t = (KT_e/m_e)^{1/2}$ denoting the mass and thermal velocity of the plasma electrons, respectively, and with the usual electron-electron collision frequency $f_{ee} = (e^4/4\pi\epsilon_0^2) (\ln\Lambda) n_e/m_e^2 v_t^3$, containing the plasma density, n_e . In most practical cases, we have $W_C \sim$ several times $KT_e \ll m_e c^2$. As mentioned, electrons exceeding this limit will run away; except, whenever $W_C \sim m_e c^2$ (i.e., for very small fields $E \ll E_C$) where the collisional drag will be larger than the accelerating force eE even for relativistic electrons (Ref. 17).

For fully-ionized plasmas, the various calculations - with quite good agreement - predict a production rate density of runaways.

$$dn_r/dt = n_e f_{ee} S(\epsilon, Z_{\text{eff}}) \quad (3)$$

where the normalized source strength $S(\epsilon, Z_{\text{eff}})$ - denoting the percentage of plasma electron running away per electronic collision time - depends only on the normalized electric field, $\epsilon = E/E_C$, and the effective ionic charge, Z_{eff} , of the plasma. Although some uncertainties still may exist concerning the scaling of S with Z_{eff} (Ref. 18), the source function S nowadays generally is approximated in the form (Ref. 16)

$$S(\epsilon, Z_{\text{eff}}) = A(Z_{\text{eff}}) \epsilon^{-3(Z_{\text{eff}}+1)/16} \exp\{-(1/4\epsilon) - [(Z_{\text{eff}}+1)/\epsilon]^{1/2}\} \quad (4)$$

with the various calculations differing only in the less important functional factor in front of the exponential, and with $A(Z_{\text{eff}})$ varying between .43-.14 over the range $Z_{\text{eff}} = 1-10$.

Actually, S falls exceedingly steeply towards small ϵ -values, so that - for $Z_{\text{eff}} = 1$ - runaway generation in most cases of interest will be significant only for $\epsilon > .02-.025$; for $Z_{\text{eff}} > 1$, the S -curve and the lower limit for effective runaway generation are shifted to somewhat larger ϵ -values roughly proportional with $Z_{\text{eff}}^{1/5}$ but the generation rate still will be negligible below a corresponding limit $\epsilon_{\text{min}} \sim (.02-.025)Z_{\text{eff}}^{1/5}$.

In contrast, the present paper points out that potentially significant additional runaway generation also may occur even at $\epsilon < .02$, namely resulting from close collisions of the thermal plasma electrons with fast particles, in particular with already existing runaway electrons: A plasma electron recoiling from such collisions with an energy $W_s > W_c$ will have a high probability to become a new runaway. This process had briefly been considered earlier by Harrison, but judged to be "very small", and consequently these contributions still are generally considered unimportant (Ref. 19). In the present paper, following a more simplified analysis, we estimate that this process actually may give rise to an avalanche-like exponential growth of the runaway population whenever runaway electrons can acquire from the electric field E energies larger than about 10-20 MeV (except when $W_c \sim m_e c^2$). Quite importantly, in that case, not only the final density but also the energy distribution of the runaway electrons and its apparent time dependence will differ strongly from those based on Eqs. (3) and (4). Acceleration energies in excess of that value may be available in various situations including in next-generation tokamak devices (in particular during disruptions) as well as in some cosmic-ray-generating astrophysical plasmas.

II. Basic Treatment :

For an assessment of the significance of this process we will consider first the energy and angular distribution of the plasma electrons recoiling from a collision with the fast charged particles: In the general case of an unshielded Coulomb collision of a cold plasma electron with a fast charged particle (having a velocity $v_f \gg v_t$ and a charge Z_f), the energy distribution, dn_s/dW_s , of non-relativistic "knock-on" plasma electrons ("secondaries") is given by the energy-differential cross section (Ref. 20)

$$d\sigma/dW_s = Z_f^2 e^4 / (8\pi \epsilon_0^2 m_e v_f^2 W_s^2). \quad (5)$$

The velocity of these knock-ons always lies in the forward half-cone around the direction of the fast particle, with low-energy knock-ons, i.e., those with energies $W_s \ll (m_e v_f^2/2)$, moving nearly perpendicular to the fast particle: for secondaries with $W_s / (m_e v_f^2/2) < .1$, the angle with the fast-particle direction will be about $\theta_s > 54$ deg.

Still considering the case with $Z_{eff} = 1$, the probability, $p_r(W_s, \theta_s)$, for such secondaries to become new runaways obviously will depend on their initial energy, W_s , and their direction, θ_E , relative to the accelerating E-field, with the latter obviously depending also on the direction of the initial fast particle relative to the E-field. The actual probability, $p_r(W_s)$, for an "average" secondary with initial recoil energy W_s obviously will have to be determined by statistical methods including averages over the respective angles. In the case of secondary generation by already existing runaways - which is the case of main interest here - , the direction of the fast particles generally are very nearly anti-parallel to the accelerating E-field. Thus, the initial direction angle, $\theta_E = \pi - \theta_s(W_s)$, of the low energy secondaries all are nearly perpendicular to the accelerating E-field with a slight component in the direction of acceleration. In this case, p_r becomes essentially a function, $p(W_s)$, of W_s only.

Correspondingly, the total rate of secondary runaway generation density by already existing runaways of density n_r and velocity v_r then becomes

$$\begin{aligned} dn_{sr}/dt &= n_r n_e v_r \int_0^{W_{\max}} (d\sigma/dW_s) p_r(W_s) dW_s \quad (6) \\ &= n_r n_e (e^4/8\pi\epsilon_0^2 m_e v_r) \int_0^{W_{\max}} (p_r(W_s)/W_s^2) dW_s \end{aligned}$$

where $W_{\max}$ (either $\sim m_e c^2$ or $\sim m_e v_f^2/2$) denotes the effective maximum energy of the secondaries.

As mentioned, the exact shape of $p_r(W_s)$ will have to be determined by a more careful statistical analysis. At present, we only note that electrons with $W_s \gg W_C$ mostly will become runaways, while most electrons with $W_s \ll W_C$ will rejoin the thermal distribution, i.e., we have $P(W_s \gg W_C) \sim 1$ and $P(W_s \ll W_C) \sim 0$, with a continuous transition around $W_s \sim W_C$.

Thus, using an effective threshold energy $W_{ce} = a(Z_{\text{eff}})W_C$, with a Z_{eff} -dependent factor $a = a(z_{\text{eff}})$, we can write

$$\begin{aligned} dn_{sr}/dt &= n_r n_e (e^4/8\pi\epsilon_0^2 m_e v_r) [1/a(Z_{\text{eff}})W_C] \\ &\sim n_r (eE/2m_e \ln \Lambda) (1/a)(1/v_r). \end{aligned} \quad (7)$$

Recognizing that all secondaries are directed very nearly perpendicular to the E-field and considering in some detail the acceleration and scattering processes, we presently estimate that, for $Z_{\text{eff}} = 1$, this effective critical threshold energy is quite close to W_C , i.e. $a(Z_{\text{eff}}=1) \sim 1$ (probably within about 20-30%). For $Z_{\text{eff}} > 1$, the situation is more complex and harder to estimate: The standard critical energy W_C (at which the average collisional drag force exerted by the plasma electrons and ions becomes equal to the electric force eE) is proportional to $(1+Z_{\text{eff}})$. However, only collisions with the plasma electrons will directly diminish the energy of

the secondaries, while collisions with the ions only will change their directions, with energy changes only occurring subsequently through the influence of E . Correspondingly, secondaries starting initially below the official W_C will tend to heat up, so that a sizable fraction of them can be expected to end up as runaways rather than be dragged back into the thermal population. Obviously, this process also is reflected in the earlier mentioned, relatively slow ("about- $Z_{eff}^{1/5}$ ") dependence of runaway generation rate $S(\epsilon, Z_{eff})$. Clearly, again, a full statistical analysis of this process will be required to calculate the effective threshold energy $W_{Ce} = a(Z_{eff} > 1) W_C$ for this case. However, considering the various processes we presently estimate that $a(Z_{eff})$ is not larger than about 2-4 even for $Z_{eff} = 10$.

Furthermore, it appears that the assumed simplifications of our basic model will not be deleterious for most applications of present interest: (i) The Debye shielding length generally is large compared with the impact parameters needed for scattering thermal electrons to energies $W_S > W_C$; and (ii) the actual thermal spread of the plasma electrons only will lead to a small smearing of the velocity distribution of the scattered electrons near W_C , but - due to $W_C \sim$ about (10-20) KT, we estimate that the thermal spread does not seriously influence the overall rate of runaway production in Eq. (7).

For analyzing the actual significance of this avalanche process, we will first estimate separately the numbers of secondaries produced (i) while an existing runaway electron is accelerated from the critical energy W_C to relativistic speeds, and (ii) during the runaway's further acceleration at relativistic speeds, and subsequently will look at the consequences for the density and energy spectrum of the final runaway population.

(i) Concerning the number of next-generation secondaries, N_{Sn} , generated by each runaway while it is being accelerated from the critical velocity, $v_C = (2W_C/m_e)^{1/2}$, to near-relativistic speeds, we consider

$$dn_{Sr}/dv_r \sim (dn_{Sr}/dt)(m_e/eE) = n_r(1/a \ln \Lambda)(1/v_r) \quad (8)$$

and get

$$\begin{aligned}
 N_{sn} &\sim (1/n_r) v_c \int_c^{\infty} (dn_{sr}/dv_r) dv_r = (1/a \ln \Lambda) \ln(c/v_c) \\
 &\sim (1/2a \ln \Lambda) \ln(m_e c^2/2W_c).
 \end{aligned}
 \tag{9}$$

Although still dependent on the value of $\ln \Lambda$, this number generally will be negligible in most applications.

(ii) For considering the generation of secondaries by a relativistic runaway, we can replace v_f in Eq. (7) with c , and obtain:

$$\begin{aligned}
 dn_{sr}/dW_r &= (dn_{sr}/dt)(1/ecE) = \\
 &= n_r / 2m_e c^2 a(Z_{eff}) \ln \Lambda
 \end{aligned}
 \tag{10}$$

where W_r denotes the kinetic energy of the relativistic runaway. Thus, taking $\ln \Lambda \sim 10-20$, we find that each runaway in the average will give rise to one next-generation runaway approximately for each $\delta W_r \sim a(Z_{eff}) \times (10-20 \text{ MeV})$ it receives from the electric field E . Very interestingly, this energy value δW_r is widely independent of the plasma parameters, in particular of density and temperature, except through $\ln \Lambda$ and $a(Z_{eff})$; as apparent from our earlier remarks, this result will always apply, except in the unlikely cases where the Debye length is smaller than the impact parameter required for transferring to the recoiling secondary an energy equal to the effective critical energy aW_c . Similarly, δW_r is independent also of the electric field E , except for very small E when aW_c becomes relativistic, i.e., when even relativistic electrons effectively will no longer be accelerated.

Neglecting the generation of secondaries by non-relativistic runaways, voiding the differentiation between various runaway generations and

assuming no losses of the generated runaways, Eq. (7) indicates an avalanche-like exponential growth of the relativistic runaway population

$$n_r = n_{r0} \exp (t/t_0) \quad (11)$$

with

$$t_0 = 2m_e c a (Z_{\text{eff}}) (\ln \Lambda) / e E \quad (12)$$

i.e., with an exponentiation time equal to the time needed for a relativistic runaway to gain the above mentioned energy of $a(Z_{\text{eff}})(10-20 \text{ MeV})$ from the E-field; obviously, this time will be increased - although not materially - by the mentioned secondary generation from nonrelativistic runaways.

Since the most recently born runaways will have gained the least energies, the energy distribution of the runaways also will tend to be exponential: Neglecting again secondary generation by nonrelativistic runaways and any losses of runaways, the energy spectrum of the resulting runaways will be

$$(1/n_r) dn_r / dW_r \sim (1/W_r) \exp[-W_r / \delta W_r]. \quad (13)$$

In case the runaways are lost at a certain maximum energy $W_{r,\text{max}}$ the above spectrum obviously will be cut off at that energy.

Overall, such runaway avalanching clearly will be important only in cases where the runaways will gain from the electric field energies in excess of $am_e c^2 \ln \Lambda \sim a (10-20 \text{ MeV})$. This limit excludes most of the earlier plasma experiments; however, it can easily be broached in the presently existing and planned large tokamak devices - in particular during disruptions where inductive potentials of up to $LI/2\pi R \sim 50-100 \text{ MeV}$ are available for each 1 MA of discharge current I (Ref. 21). Equally, it appears that such multiplication could also be significant during the acceleration of cosmic rays whenever such acceleration occurs at least in part parallel to the magnetic field.

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