Two types of radiative disruptions in tokamaks

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The nature of a radiative disruption in a tokamak is studied as a function of the mass of the radiating impurity. The thermoelectric power exerted on the impurity particles by the plasma causes a symmetric disruption in the case of a heavy impurity and a MARFE disruption in the case of a light impurity.

1. The radiative disruption is a familiar effect in tokamak physics. It imposes an upper limit on the plasma density which can be attained. A qualitative explanation of this disruption runs as follows:[1–3] When the plasma density reaches a critical value, the power radiated by impurities near the tokamak wall becomes so high that the energy arriving from the central region of the plasma is not sufficient to replenish the energy lost by radiation, and the plasma column begins to cool off. Actually, the picture of the disruption is more complex: The disruption can be either poloidally symmetric (in the case of a hard disruption) or asymmetric (in the case of a MARFE disruption).[2,4] Drake[2] has suggested that a MARFE disruption results from a radiation-condensation instability, while a hard disruption results from a simple radiation instability.

The growth rate of the simple radiation instability is[5]

$$\gamma = -(2/3)\partial L/\partial T, \tag{1}$$

where L is the radiation power per plasma particle, and T is the plasma temperature. Expression (1) is positive, since we assume that we are working at temperatures at which the radiation intensifies with decreasing T. Situations of this sort are known in astrophysics and tokamak physics. The radiation-condensation instability has a growth rate[6]

$$\gamma = (2/5)(-\partial L/\partial T + L/T). \tag{2}$$

This instability occurs because even in the case $\partial L/\partial T = 0$ a local decrease in the temperature leads to an increase in radiation, because of an increase in the density (if the pressure can be assumed constant).

According to Drake,[2] the threshold for the radiation-condensation instability in a tokamak is always lower than the threshold for the simple radiation instability, and a hard disruption must always be preceded by MARFE. However, it was recently observed[7] that the disruption passes through a MARFE stage in the JET tokamak with beryllium walls, while in a tokamak with graphite walls the hard disruption occurs immediately. In this letter we show that the nature of the disruption may be changed by the thermoelectric power exerted on the radiating impurity by the bulk plasma.

2. A qualitative explanation of the role of the thermoelectric power runs as follows. The longitudinal thermoelectric power exerted on heavy impurities by the plasma is directed along the temperature gradient. The impurities are displaced from the cold radiating regions, so a radiation-condensation instability is prevented. This effect is weaker in the case of a relatively light impurity. The thermoelectric power depends on the impurity mass because of its ion part.[8] One might therefore expect that the MARFE mode would be suppressed by the thermoelectric power in the case of a heavy impurity.

For quantitative calculations, we use a simple system of equations for the temperature T, the plasma density n, the impurity density N, the longitudinal plasma velocity v, and the longitudinal impurity velocity V:

$$(3/2)\partial_t T = \nabla \chi \nabla T - T\nabla_{\parallel} v - L, \tag{3a}$$

$$\partial_t n + \nabla_{\parallel}(nv) = 0, \tag{3b}$$

$$Mn\partial_t v = -\nabla_{\parallel}(nT), \tag{3c}$$

$$\partial_t N + \nabla_{\|}(NV) = 0, \tag{4a}$$

$$AM\partial_t V = -\nabla_{\|}(NT)/N + \sigma\nabla_{\|}T + \nu AM(v - V). \tag{4b}$$

Here χ is the thermal diffusivity, ν is the impurity-plasma friction, σ is the thermoelectric power, and A is the atomic number of the impurity.

In the case $k_{\|} = 0$ (a hard disruption) the instability growth rate is determined by Eq. (3a) alone:

$$\gamma = (2/3)(-\partial L/\partial T - \chi_{\perp}k_x^2). \tag{5}$$

Here k_x is the radial component of the wave vector. This component is determined by the characteristic width of the radiating zone.[2] The result in (5) does not depend on whether the thermoelectric power is taken into consideration.

In the case $k_{\|} \neq 0$ (MARFE), under the assumption $\nu, c_s k_{\|} \gg \gamma$ [where $c_s = (T/M)^{1/2}$ is the sound velocity], we find

$$-T\delta N/(N\delta T) = [\gamma - c_s^2 k_{\|}^2(\sigma - 1)/(A\nu)][\gamma + c_s^2 k_{\|}^2/(A\nu)]^{-1}. \tag{6}$$

If we ignore the thermoelectric power ($\sigma = 0$), Eq. (6) reduces to the ordinary relation for the radiation-condensation instability: $\nabla_{\|}(NT) = 0$. In this case we find the growth rate to be

$$\gamma = (2/5)(-\partial L/\partial T + L/T - \chi_{\perp}k_x^2 - \chi_{\|}k_{\|}^2). \tag{7}$$

In tokamaks, as we pointed out earlier, MARFE threshold (7) is lower than the threshold for a hard disruption, (5). Under the assumption $L \simeq T^{-1}$, for example, we find the MARFE threshold to be $L_m = (1/2)\chi_{\perp} k_x^2$, while the threshold for a hard disruption is found to be $L_c \simeq \chi_{\perp} k_x^2$. (Here we have used[2] $\chi_{\|} k_{\|}^2 \ll \chi_{\perp} k_x^2$, but other situations are also possible.[9])

Now abandoning the assumption $\sigma = 0$, we find, in place of (7),

$$(5/2)\gamma = -\partial L/\partial T - \chi_{\perp}k_x^2 - \chi_{\|}k_{\|}^2 - (L/T)[T\delta N/(N\delta T)], \tag{8}$$

where $T\delta N/(N\delta T)$ is given by (6). Normalizing the growth rate γ by division by $\gamma_0 = \chi_{\perp} k_x^2$, we find the dispersion relation

$$(5/2)\gamma = \lambda - 1 + \lambda[\gamma - \gamma_1(\sigma - 1)][\gamma + \gamma_1]^{-1}, \tag{9}$$

where $\gamma_1 = c_s^2 k_{\|}^2/(A\nu\gamma_0)$ and $\lambda = L/\gamma_0$. (In this new notation, the MARFE threshold in the case $\sigma = 0$ is $\lambda_m = 1/2$, while the threshold for a hard disruption is $\lambda_c = 1$.)

From (9) we can find the MARFE threshold λ_m as a function of σ and γ_1. To do this, we need to solve the equation $\mathrm{Re}\gamma = 0$. As a result, we find

$$\lambda_m = \begin{cases} 1/2 + 5\gamma_1/4, & \gamma_1 < 2\sigma/(10-5\sigma) \\ 1/(2-\sigma), & \gamma_1 > 2\sigma/(10-5\sigma) \end{cases}. \quad (10)$$

As expected, we have $\lambda_m > \lambda_m(\sigma - 0) = 1/2$. If the MARFE is to be observed with increasing λ, we must also require that the threshold λ_m be less than one. Otherwise, a hard disruption will occur first. Using the relation $\sigma > 1$, which holds for real radiating impurities, we rewrite the condition $\lambda_m < 1$ as $1/2 + 5\gamma_1/4 < 1$. Hence $\gamma_{1c} = 2/5$. In the case $\gamma_1 < \gamma_{1c}$ we have a MARFE disruption, while in the case $\gamma_1 > \gamma_{1c}$ we have a hard disruption.

Estimating $\chi_\perp$ and k_x by the method of Ref. 2, we find that γ_1 is proportional to $T_L^{9/2}$, where T_L is a characteristic temperature of the impurity radiation. The impurities in the peripheral zone of a tokamak are not at equilibrium, and the corona model gives a poor description of the radiation. According to Ref. 10, a switch from beryllium to graphite results in a threefold increase in T_L, so γ_1 increases by two orders of magnitude, and the MARFE may be impossible.

We note in conclusion that dispersion relation (9) is a model relation, and the numbers $\gamma_{1c} = 2/5$, etc., should not be taken seriously. The reason is that the perpendicular thermal conductivity $\chi_\perp$ depends on the plasma temperature and density, and just how is quite uncertain. We have assumed that $\chi_\perp$ is a constant. Clearly, if we allow $\chi_\perp$ to vary with n and T, then the critical value γ_1 will change by an amount, in general, on the order of $\max(\partial \ln \chi_\perp/\partial \ln T, \partial \ln \chi_\perp/\partial \ln n)$, i.e., by an amount on the order of one. However, the qualitative effects (the stabilizing role of the thermoelectric power and the dependence of the magnitude of the effect on the nature of the radiating impurity) will be present in any model.

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