Interaction between a plasma jet and a chemically active ionized gas target

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A theoretical study is reported of the linear stage of turbulence development during the interaction between a beam of ions and a chemically active weakly ionized target. It is shown that when inhomogeneities in the beam density are taken into account, this leads to the appearance of a growth threshold that depends on the chemical reaction and charge exchange cross sections, a fact that can be used to determine the cross sections as functions of the beam velocity.

The ion beam method[1] has long been used to investigate the kinetics of ion-molecule reactions, i.e., reactions of the form

$$B^{+}+A_2\rightarrow AB^{+}+A, \qquad (1)$$

By examining the various characteristics of the particles leaving the target one can then determine the cross sections for the elementary processes involved in this system. The main disadvantage of this method is that all the modern sources of ion beams have very small saturation currents ($\lesssim 10^{-6}$ A for U ~ 10 v). This, of course, introduces considerable difficulties for the detection of the chemical reaction products. Moreover, this method cannot be used to produce the above type of chemical reaction on an industrial scale.

It is clear that the effective ion current can be increased by replacing the ion beam with a quasineutral jet of plasma in which the electrons and ions move with directed velocities. It is also important to note that, since in many cases the cross section for the reaction given by (1) is found to increase when the molecules are in an excited state, the target temperature must be sufficiently high, i.e., the target may be a weakly ionized plasma in which the electron density n_e is much less than the neutral-particle density n_n. Under these conditions we have the usual (for plasma systems) problem of the stability of the interaction. It is clear that the appearance of instability may lead to a substantial distortion of the charged-particle distribution function[2] and, consequently, may affect the rate at which the chemical product is produced. In this paper we investigate the case where the ion density is such that $n_B \ll n_e \ll n_n$. Moreover, we shall assume that the directed velocity of the ions and electrons in the beam is $u \ll v_{Te}$. The electrons in the plasma jet cannot then lead to additional instability and may be eliminated from the analysis.

Let us, therefore suppose that the beam of cold ions of density n_B and mass M_B is travelling along the x axis with velocity u. The weakly ionized target consists of molecules A_2 of density n_{A2}, cold ions A^+, and electrons with $n_e = n_a$ at temperature T_e. The target occupies the layer $0 \le x \le L$. It is important to note that in addition to the molecules A_2 the target may also contain the atoms A. They, however, do not participate either in the chemical interaction or in the interaction with the charged particles, and can also be eliminated from the analysis. As already noted, we are assuming that $n_{A_2} \gg n_e \gg n_B$. We shall also assume that $M_B \gg M_{A_2}$ and that collisions of ions B^+ with particles A_2 do not lead to an appreciable change in the directed velocity, i.e., that the velocity of the reaction products is $u_{AB^+} = uM_B/(M_B + M_A) \approx u$.

The stability problem will be solved in the linear approximation, using the equations of multicomponent hydrodynamics. Before we write out these equations, let us consider in greater detail the main processes which distinguish this problem from the well known beam problem in high-temperature plasma.[2] These processes include various types of collision leading, firstly, to the relaxation of the particle momentum and, secondly, to the appearance of density inhomogeneities. For the A^+ particles the main type of collision is resonance charge transfer on atoms A ($A^+ + A \rightarrow A + A^+$), which leads to a substantial relaxation of their momentum (it is well known that the cross section for this process is greater than or of the order of the gas-kinetic cross section), but does not lead to a reduction in the number of the A^+ ions (we are assuming that the target particles are in thermodynamic equilibrium). Consequently, resonance charge transfer is allowed for in the equation of motion for A^+ but not in the continuity equation.

For the B^+ ions, the main processes are collisions leading to the formation of the chemical product AB^+ [see (1)] and charge transfer on neutral atoms and molecules in the target (the cross section for this process may be either greater or smaller than the gas-kinetic cross section depending on the extent to which charge transfer approaches resonance). Both processes lead to a reduction in the density of the B^+ ions and, therefore, enter the equation of continuity. As already noted, we are assuming that $M_B \gg M_A$ and, therefore, all elastic cross sections leading to the relaxation of the momentum of the B^+ ions can be neglected (in a single collision and the change in momentum is $\Delta p/p \approx M_A/M_B \ll 1$).

Finally, for the AB^+ ions, the main processes are chemical interaction, leading to an increase in their density, and charge transfer on neutral particles, leading to a reduction in the number of the AB^+ ions. As in the case of B^+, we are neglecting elastic scattering for the particles AB^+. We shall also suppose that the unstable oscillations are potential and have low enough frequency [$\omega \ll \omega_{pe} = (4\pi e^2 n_e/m)^{1/2}$], so that electrons in the field of this wave succeed in reaching the Boltzmann distribution.

In view of the foregoing, the equations for the problem can be written as follows:

$$\frac{\partial^2\varphi}{\partial x^2} = -4\pi e(n_A' + n_B' + n_{AB}' - n_e'), \qquad n_e' = \frac{e\varphi}{T_e},$$

$$\frac{\partial v_A'}{\partial t} = -\frac{e}{M_A}\frac{\partial\varphi}{\partial x} - \nu_0 v_A', \qquad \frac{\partial n_A'}{\partial t} + n_A\frac{\partial v_A'}{\partial x} = 0,$$

$$\frac{\partial v_B'}{\partial t} + u\frac{\partial v_B'}{\partial x} = -\frac{e}{M_B}\frac{\partial\varphi}{\partial x},$$

$$\frac{\partial n_B'}{\partial t} + u\frac{\partial n_B'}{\partial x} + \frac{\partial}{\partial x} n_B v_B' = -(\nu_x + \nu_1) n_B',$$

$$\frac{\partial v_{AB}'}{\partial t} + u\frac{\partial v_{AB}'}{\partial x} = -\frac{e}{M_{AB}}\frac{\partial\varphi}{\partial x}$$

$$\frac{\partial n_{AB}'}{\partial t}+u\frac{\partial n_{AB}'}{\partial x}+\frac{\partial}{\partial x}n_{AB}v_{AB}'=\nu_x n_B'-\nu_2 n_{AB}'. \quad (2)$$

In these expressions φ is the perturbation of the electric potential of the wave, all the perturbed quantities are primed, $\nu_0\langle\sigma_0 v_{Ta}\rangle n_A$ is the frequency of resonance charge transfer between ions and atoms of type A, $\nu_1=\langle\sigma_1 v_{TB}\rangle n_n$ is the charge transfer frequency between ions of type B^+ and neutrals in the target, $\nu_2=\langle\sigma_2 v_{TAB}\rangle n_n$ is the charge transfer frequency between ions of type AB^+ and neutrals in the target, and $\nu_x=\langle\sigma_x v_{TB}\rangle n_{AB}$ is the frequency of collisions between ions of type B^+ and molecules A_2, which lead to the formation of the ions AB^+.

In the zero-order approximation $\varphi=0$, and it follows from (2) that

$$n_A=n_0=\text{const},\quad n_B=n_{0B}\exp\left(-\frac{\nu_1+\nu_x}{u}x\right),$$
$$n_{AB}=n_{0B}\frac{\nu_x}{\nu_1+\nu_x-\nu_2}\left[\exp\left(-\frac{\nu_2}{u}x\right)-\exp\left(-\frac{\nu_1+\nu_x}{u}x\right)\right]. \quad (3)$$

Therefore, the set of equations given by (2) for the perturbed quantities contains coefficients which are explicit functions of x and, therefore, these perturbations can be sought in the form $\varphi\sim\varphi(x)e^{-i\omega t}$. The derivation of the dispersion relation in the general form is quite laborious and, therefore, we shall give it in detail only for the case when there are no chemical interactions, i.e., when $\nu_2=\nu_x=0$. The equations given by (2) then reduce to the form

$$\left[1-\frac{\omega_{pa}^2}{\omega(\omega+i\nu_0)}\right]\frac{\partial^2\varphi}{\partial x^2}-\frac{\omega_{pa}^2}{c_s^2}\varphi=-4\pi e n_B', \quad (4)$$

$$-i\omega v_B'+u\frac{\partial v_B'}{\partial x}=-\frac{e}{M_B}\frac{\partial\varphi}{\partial x}, \quad (5)$$

$$-i\omega n_B'+u\frac{\partial n_B'}{\partial x}+\frac{\partial}{\partial x}n_B v_B'=-\nu_1 n_B'. \quad (6)$$

In these expressions $\omega_{pa}^2 4\pi e^2 n_0/M_A$ and $c_s^2=T_e/M_A$.

Applying the operator $(-i\omega+u\partial/\partial x+\nu_1)$ to both sides of (4), and using (3) and (6), we obtain

$$e^{\nu_1 x/u}\left(-i\omega+u\frac{\partial}{\partial x}+\nu_1\right)\left[\left(1-\frac{\omega_{pa}^2}{\omega(\omega+i\nu_0)}\right)\frac{\partial^2\varphi}{\partial x^2}-\frac{\omega_{pa}^2}{c_s^2}\varphi\right]=-4\pi e n_{0B}\left(\frac{\partial}{\partial x}-\frac{\nu_1}{u}\right)v_B'. \quad (7)$$

Applying the operator $(-i\omega+u\partial/\partial x)$ to both sides of these equations, and using (5), we finally obtain

$$\left(-i\omega+u\frac{\partial}{\partial x}+\nu_1\right)^2\left[\left(1-\frac{\omega_{pa}^2}{\omega(\omega+i\nu_0)}\right)\frac{\partial^2\varphi}{\partial x^2}-\frac{\omega_{pa}^2}{c_s^2}\varphi\right]=\omega_{pb}^2 e^{-\nu_1 x/u}\left(\frac{\partial}{\partial x}-\frac{\nu_1}{u}\right)\frac{\partial\varphi}{\partial x}. \quad (8)$$

In this expression $\omega_{pb}^2=4\pi e^2 n_{0B}/M_B$.

In the general case, a similar procedure leads to the following dispersion relation:

$$\left(-i\omega+u\frac{\partial}{\partial x}+\nu_1+\nu_x\right)^3\left(-i\omega+u\frac{\partial}{\partial x}+\nu_2\right)^2\left[\left(1-\frac{\omega_{pa}^2}{\omega(\omega+i\nu_0)}\right)\frac{\partial^2\varphi}{\partial x^2}-\frac{\omega_{pa}^2}{c_s^2}\varphi\right]=-\omega_{pb}^2\exp\left(-\frac{\nu_1+\nu_x}{u}x\right)\left(-i\omega+u\frac{\partial}{\partial x}+\nu_2-\nu_1-\nu_x\right)$$
$$\times\left(-i\omega+u\frac{\partial}{\partial x}\right)\left(-i\omega+u\frac{\partial}{\partial x}+\nu_2-\nu_1\right)\left(\frac{\partial}{\partial x}-\frac{\nu_1+\nu_x}{u}\right)\frac{\partial\varphi}{\partial x}$$
$$-\omega_{px}^2\exp\left(-\frac{\nu_1+\nu_x}{u}x\right)\left(-i\omega+u\frac{\partial}{\partial x}+\nu_2-\nu_1-\nu_x\right)\left(-i\omega+u\frac{\partial}{\partial x}\right)^2$$
$$\times\left(\frac{\partial}{\partial x}-\frac{\nu_1+\nu_x}{u}\right)\frac{\partial\varphi}{\partial x}+\omega_{px}^2\exp\left(-\frac{\nu_2}{u}x\right)\left(-i\omega+u\frac{\partial}{\partial x}+\nu_1+\nu_x-\nu_2\right)^3\left(\frac{\partial}{\partial x}-\frac{\nu_2}{u}\right)\frac{\partial\varphi}{\partial x},$$

$$\omega_{px}^2=\frac{4\pi e^2 n_{0B}}{M_{AB}}\frac{\nu_x}{\nu_1+\nu_x-\nu_2}. \quad (9)$$

Equation (9) cannot be solved in a general form, and we shall therefore examine only some special cases which are most readily realized in practice. However, before we proceed to the particular analysis of these situations, we must consider a number of further points. Analysis of the instability of oscillations produced in plasma by an ion beam is meaningful only when $\omega\gg\nu_{max}$. When this condition is not satisfied, the oscillations will of course be damped. Since in beam instabilities the growth rate is a maximum for oscillations with phase velocity approaching the beam velocity, the inequality $\omega\gg\nu_{max}$ can be rewritten in the form

$$\left|\frac{\partial\ln\varphi}{\partial x}\right|\gg\frac{\nu_{max}}{u}. \quad (10)$$

We shall therefore assume throughout that the inequality given by (10) is satisfied. Finally, we recall that the inequality $n_B\ll n_A$ was assumed right from the outset. Since the problem contains small parameters, we can use perturbation theory to solve (9). We shall describe this in detail for the special case where there are no chemical interactions in the system.

1. INTERACTION OF AN ION BEAM WITH WEAKLY IONIZED TARGET IN THE ABSENCE OF CHEMICAL REACTIONS

This situation is encountered when dense, low-energy beams of neutral particles are produced by charge transfer from a plasma jet to a gas target. In this case $\nu_x=\nu_2=0$, and (9) takes the form given by (8). Assuming that $\omega=\omega_0+\delta$ and $\omega_0\gg\delta$, we can rewrite (8) in the form

$$\left(-i\omega+u\frac{\partial}{\partial x}+\nu_1\right)^2\left[\left(1-\frac{\omega_{pa}^2}{\omega_0^2}\right)\frac{\partial^2\varphi}{\partial x^2}-\frac{\omega_{pa}^2}{c_s^2}\varphi\right]=-\omega_{pb}^2 e^{-\nu_1 x/u}\frac{\partial^2\varphi}{\partial x^2}\left(-i\omega+u\frac{\partial}{\partial x}+\nu_1\right)^2\frac{\omega_{pa}^2}{\omega_0^3}(2\delta+i\nu_0)\frac{\partial^2\varphi}{\partial x^2}. \quad (11)$$

The solution of (11) will be sought in the form of the series $\varphi=\varphi_0+\varphi_1+\ldots$ where $n_{0B}/n_0\ll 1$ and $\varphi_1\ll\varphi_0$. In the zero-order approximation, we have from (11)

$$\left(1-\frac{\omega_{pa}^2}{\omega_0^2}\right)\frac{\partial^2\varphi_0}{\partial x^2}=\frac{\omega_{pa}^2}{c_s^2}\varphi. \quad (12)$$

It follows from this equation that $\varphi=\varphi_0 e^{ikx}$. The relationship between the oscillation frequency and the wave vector k is given by

$$\omega_0=\frac{\omega_{pc}}{(1+\omega_{pa}^2/k^2c_s^2)^{1/2}}. \quad (13)$$

In the first approximation we have from (11), using (13),

$$\left(-i\omega+u\frac{\partial}{\partial x}+\nu_1\right)^2\frac{\omega_{pa}^2}{k^2c_s^2}\left[\frac{\partial^2\varphi_1}{\partial x^2}+k^2\varphi_1\right]=k^2\omega_{pb}^2e^{-\nu_1 x/u}\varphi_0 e^{ikx}$$
$$-(\omega-ku+i\nu_1)^2\frac{\omega_{pa}^2}{\omega_0^3}(2\delta+i\nu_0)k^2\varphi_0 e^{ikx}. \quad (14)$$

We now multiply (14) from the left by $\varphi_0^* e^{-ikx}$ and integrate it with respect to x across the target $0\le x\le L$. Successive integration of the left-hand side of the resulting equation by parts, using the fact that the operator $\partial^2/\partial x^2-k^2$ is self adjoint, and the fact that the potential and all its derivatives are zero on the boundaries of the target, we obtain

$$(\omega-ku+i\nu_1)^2\frac{\omega_{pa}^2}{\omega_0^3}(2\delta+i\nu_0)=\omega_{pb}^2\frac{u}{\nu_1 L}(1-e^{-\nu_1 L/u}). \quad (15)$$

Before we proceed to investigate this equation, we must introduce the following remark. We recall that, firstly, (15) is valid only at the initial stage of development of the instability, when the oscillation amplitude is still small enough and the reaction of the waves on the distribution of ions in the beam in plasma can be neglected. Secondly, it was assumed in the derivation of (15) that the wave potential was zero on both boundaries of the target. At the point of entry of the beam into the target this is always valid and, since at x = 0 the oscillations increase from the thermal-noise level, their amplitude can be neglected. The condition that the potential at x = L is zero imposes certain definite restrictions on the length of the target and the growth rate. In fact, since we have included in our analysis only waves travelling in the direction of propagation of the beam (such oscillations have the maximum growth rate), the amplitude at x = L is not in general zero. After reflection from the plasma boundary and transformation in the thin transition layer, these oscillations leave the system (we recall that we are taking into account only oscillations traveling in the direction of propagation of the beam), and this leads to a reduction in the wave energy within the plasma.

The influence of this effect on the growth rate can be estimated quite readily by considering the integral equation for the energy balance:

$$\frac{\partial}{\partial t}\bar{\varphi}=\bar{\gamma}\bar{\varphi}-V_{gr}\frac{\bar{\varphi}}{L}. \tag{16}$$

In this expression $\bar{\varphi}$ is the mean amplitude of the potential within the target, $\bar{\gamma}$ is the mean growth rate given by (15), and v_{gr} is the group velocity of the oscillations. It follows from (16) that the escape of the oscillations from the system can be neglected when

$$\bar{\gamma}>V_{gr}/L. \tag{17}$$

Specific estimates based on this formula will be given below for each particular case.

We now return to (15). We have already noted that the oscillations with ω_0 = ku have the maximum growth rate. In that case, we have from (15)

$$\bar{\gamma}=\operatorname{Im}\omega=\begin{cases}\gamma_{max}=\omega_0\left[\dfrac{n_{0B}}{n_0}\dfrac{u}{\nu_1 L}\left(1-\exp\left(-\dfrac{\nu_1 L}{u}\right)\right)\right]^{1/3}, & \gamma_{max}>\nu_0,\nu_1\\ \left[\dfrac{n_{0B}}{n_0}\dfrac{\omega_0^3}{\nu_0}\dfrac{u}{\nu_1 L}\left(1-\exp\left(-\dfrac{\nu_1 L}{u}\right)\right)\right]^{1/4}, & \nu_1<\gamma_{max}<\nu_0\end{cases} \tag{18}$$

In both cases the oscillations are unstable.

If, on the other hand, we have

$$\left[\frac{n_{0B}}{n_0}\frac{\omega_0^3}{\nu_0}\frac{u}{\nu_1 L}\left(1-\exp\left(-\frac{\nu_1 L}{u}\right)\right)\right]^{1/4}<\nu_1, \tag{19}$$

then the instability cuts off and Im $\simeq -\nu_1$. This can be explained physically as follows. It follows from (2) that the rate at which the beam ions disappear is proportional to their density. Since the beam ions move with velocity approaching the phase velocity of the wave, their density in the wave field is modulated. The instability tends to increase this modulation, and consequently, the opposite effect of density spreading leads to a reduction in the wave energy, i.e., to the suppression of the instability.

In addition to the inequalities given by (18) and (19), there is a further condition which is necessary for the appearance of instability. In fact we have already noted that the maximum growth rate is exhibited by oscillations with phase velocity approaching the ion velocity in the beam. From (13) we then obtain the following condition for the appearance of instability:

$$u\leq c_s. \tag{20}$$

In the opposite case, only the oscillations propagating at an angle to the direction of propagation of the beam can be unstable.

Finally, let us estimate the effect of the escape of oscillations on the growth rate. Since $v_{gr} \lesssim u$ throughout, it follows from (17) that the necessary condition for the development of instability is

$$L>u/\gamma_{max}. \tag{21}$$

It is interesting to note that the threshold for the development of instability depends explicity on the charge transfer cross section of the B^+ ions. This can probably be used as a method for determining the dependence of the cross section on the velocity of the incident ions.

2. INTERACTION OF AN ION BEAM WITH WEAKLY IONIZED TARGET FOR $\nu_1 \sim \nu_2 \gg \nu_x$

This case includes the interaction between a plasma jet and a target in which the original B^+ ions and the product AB^+ ions undergo resonance charge transfers on the target atoms. In fact, since the resonance charge-transfer cross section is much greater than the cross sections for all the other processes, it may be supposed that $\nu_x \ll \nu_1, \nu_2$. If in addition $\nu_1 \approx \nu_2$ and $M_B \approx M_{AB}$ (and, consequently, $u_B \approx v_{AB}$), it follows from (2) that the set of B^+ and AB^+ ions is equivalent in the electrodynamic sense to a beam of B^+ ions interacting with the plasma in the absence of the chemical reactions. This has already been considered above.

3. THE CASE ν_2 = 0 (NO AB^+ CHARGE TRANSFERS)

It is known[3] that, if the ionization potential of the incident ion is less than the ionization potential of the target atom, charge transfer is possible only at very high energies of relative motion of the particles ($\epsilon \sim M\Delta E/m$, where ΔE is the ionization-potential difference between the atom and the ion). As an example, let us consider the interaction between oxygen ions O^+ and a weakly ionized hydrogen target. The ionization potential of oxygen is I_O = 13.61 eV and the ionization potential of atomic hydrogen is I_H = 13.59 eV. Finally, the ionization potential of the radical is I_{OH} = 13.18 eV. When the oxygen ions enter the gaseous target they undergo charge transfers on the hydrogen atoms. At the same time, there is the chemical reaction resulting in the formation of OH^+. These ions cannot undergo charge transfers, and freely pass through the target. In this case, ν_2 = 0 and (9) assumes the form

$$\begin{aligned}&\left(-i\omega+u\frac{\partial}{\partial x}+\nu_1+\nu_x\right)^3\left(-i\omega+u\frac{\partial}{\partial x}\right)^2\left[\left(1-\frac{\omega_{pa}^2}{\omega(\omega+i\nu_0)}\right)\frac{\partial^2\varphi}{\partial x^2}-\frac{\omega_{pa}^2}{c_s^2}\varphi\right]\\&=-\omega_{pb}^2\exp\left(-\frac{\nu_1+\nu_x}{u}x\right)\left(-i\omega+u\frac{\partial}{\partial x}-\nu_1-\nu_x\right)\left(-i\omega+u\frac{\partial}{\partial x}\right)\\&\times\left(-i\omega+u\frac{\partial}{\partial x}-\nu_1\right)\left(\frac{\partial}{\partial x}-\frac{\nu_1+\nu_x}{u}\right)\frac{\partial\varphi}{\partial x}-\omega_{px}^2\exp\left(-\frac{\nu_1+\nu_x}{u}x\right)\\&\times\left(-i\omega+u\frac{\partial}{\partial x}-\nu_1-\nu_x\right)\left(-i\omega+u\frac{\partial}{\partial x}\right)^2\left(\frac{\partial}{\partial x}-\frac{\nu_1+\nu_x}{u}\right)\frac{\partial\varphi}{\partial x}\\&+\omega_{px}^2\left(-i\omega+u\frac{\partial}{\partial x}+\nu_1+\nu_x\right)^3\frac{\partial^2\varphi}{\partial x^2}.\end{aligned} \tag{22}$$

Using the above method, we find that, as before, the relation between the frequency of the excited oscillations and the wave vector is given by (13). In the first approximation, it follows from (22) that

$$(\delta+i\nu_0)(\delta+i\nu_1+i\nu_x)^3\delta^2\frac{\omega_{pa}^2}{\omega_0^3}=\omega_{pb}^2\frac{u}{(\nu_1+\nu_x)L}\left(1-\exp\left(-\frac{\nu_1+\nu_x}{u}L\right)\right)$$
$$\times\delta(\delta-i\nu_1)(\delta-i\nu_1-i\nu_x)-\omega_{px}^2\frac{u}{(\nu_1+\nu_x)L}\left(1-\exp\left(-\frac{\nu_1+\nu_x}{u}L\right)\right) \quad (23)$$
$$\times\delta^2(\delta-i\nu_x-i\nu_1)+\omega_{px}^2(\delta+i\nu_1+i\nu_x)^3.$$

In these expressions $\omega_{px}^2 = \omega_{pb}^2\nu_x/(\nu_1 + \nu_x)$.

Let us suppose to begin with that $\nu_x \ll \nu_1$. The oscillation growth rate is then very dependent on the length of the target. In fact, when $L < u/\nu_x$ we can neglect in (23) the presence of the beam of AB^+ ions, and the conditions for the appearance of instability are given by (18)–(21). When the target is thick enough and $L > u/\nu_x$, we can neglect in (23) the first two terms on the right-hand side:

$$\delta^2(\delta+i\nu_0)=\omega_0^2\omega_{px}^2/\omega_{pa}^2. \quad (24)$$

It is known[4, 5] that (24) always has solutions with $\mathrm{Im}\,\omega > 0$, i.e., the instability will always develop in a sufficiently thick target. It is readily shown that, when $\nu_x \gg \nu_1$, the oscillations will be unstable for any target thickness and their growth rate will be given by (24) (to within small terms of the order of ν_1/ν_x). It is important to note, however, that when

$$\left[\frac{n_{0B}}{n_0}\frac{\omega_0^3}{\nu_0}\frac{u_x}{\nu_1}\right]^{1/2} < \nu_1$$

the instability in a sufficiently thick target will appear only for $x > u/\nu_x$, i.e., whenever the chemical reaction has practically terminated. This can be of considerable practical interest if we recall that any beam instability during the nonlinear stage leads to an appreciable broadening of the ion distribution function for the beam, and may thus influence the rate of the chemical reaction.[4, 5]

We can now summarize the above results as follows. To avoid the effect of instability on the course of the chemical reactions during the interaction between a monoenergetic ion beam and a chemically active, weakly ionized target we must, in any case, satisfy the inequality

$$\left(\frac{n_{0B}}{n_0}\frac{\omega_0^3}{\nu_0}\right)^{1/2} < \nu_1+\nu_x. \quad (25)$$

If this is so, then the instability, even if it appears, cannot lead to an appreciable change in the distribution function for the incident ion beam. When the opposite inequality is satisfied, the problem must be considered in its nonlinear version.

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14