TRITA-EPP-86-10

ON THE ROLE OF THE IONIZATION FREQUENCY TO GYROFREQUENCY RATIO IN THE CRITICAL IONIZATION VELOCITY INTERACTION

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October 1986

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Abstract

The role of the parameter ν_i/ω_{ci} for critical ionization velocity (CIV) interaction is discussed. This parameter, which can be seen as a combined condition on the neutral gas density and the magnetic field strength, is important from two different aspects. First, the value $\nu_i/\omega_{ci} = 1$ marks the limit between the regions of applicability of the homogeneous and the "ionizing front" (or inhomogeneous) models for the interaction. Second, it determines the energy transfer efficiency in the interaction. This energy transfer efficiency also depends strongly on whether the experiment is steady-state or transient in nature. In steady-state situations, the interaction has earlier been estimated to become gradually weaker when ν_i/ω_{ci} decreases below unity. It is found here that efficient CIV interation is possible at far lower values of ν_i/ω_{ci} in transient experiments. The corresponding limit is approximately $\nu_i/\omega_{ci} = 2(m_e/m_i)^{3/2}\,\kappa^3$. Unfortunately, κ is a rather uncertain parameter: the number of instability growth times required for the modified two-stream instability to tap the energy of a thin ion stream in a plasma.

1. Introduction

The critical ionization velocity (CIV) interaction is a phenomenon which can occur when a plasma and a neutral gas are in a state of relative motion across a magnetic field. It was postulated by Alfvén (1954) that a strong interaction with rapid ionization of the neutrals should be expected when the relative velocity exceeds the critical ionization velocity, v_c, which is given by

$$v_c = \left(\frac{2e\ U_i}{m_n} \right)^{1/2} \qquad (1)$$

where m_n and U_i are the mass and ionization potential of the neutrals.

There has been a wealth of experimental observations of the v_c-interaction, both in laboratory experiment of very different geometries (e.g. Danielsson and Brenning, 1975; Axnäs, 1977; Möbius et al, 1979) and in space (Haerendel, 1982). A computer simulation by Machida et al (1984) also shows a strong increase in the ionization rate at the threshold velocity given by Eq.(1).

As illustrated in Fig.1, CIV interaction has been proposed to operate under very different conditions. In some cases, the interaction is more or less constant in time. Examples of this are the shuttle glow (Papadopoulos, 1983), the production of the Io plasma torus, and some laboratory experiments. Other cases are of a transient nature: the plasma can be a short pulse as in some impact experiments, or the neutral gas can be transient as in the ionospheric release experiments. The geometries are also very different. The magnetic field lines can close within the volume, connect to walls that are stationary with respect to the neutral component, or connect to a surrounding plasma which moves with respect to the neutrals.

In view of these difference one can hardly expect one single model to cover all the situations. However, there are some parameters which are of common interest for most cases: in addition to some minimum value of the velocity, the magnetic field must not be too weak (Brenning, 1985) and the neutral gas density and extension must be above some limits (Haerendel, 1982; Brenning, 1982 a). This paper will discuss another such parameter, ν_i/ω_{ci}, which can be seen as a combined condition on neutral gas density and magnetic field strength (in this case an upper limit). As pointed out by Haerendel (1982), ν_i/ω_{ci} separates the regions of applicability of the "ionizing front" and the "homogeneous" models for electron heating. It also is an important parameter for the efficiency of the the electron heating in the interaction, as discussed by e.g. Galeev (1981), Formisano et al (1982) and Galeev and Sagdeev (1983).

2. The electron heating

The theoretical work on CIV interaction has been focussed on the heating of the electrons, which is necessary for the rapid ionization seen in the experiments. The first-order approach to the electron energy balance is the following: the plasma and the neutral gas initially move with a relative velocity v_o. At each ionization of a neutral particle, the energy $m_n v_o^2/2$ becomes available in the plasma frame. If a fraction η of this energy is transferred to the electrons, on a timescale much faster than the ionization time, the average electron energy will increase according to

$$d\langle n_e W_e\rangle/dt = \nu_i n_e(\eta\, m_n v_o^2/2 - eU_i), \qquad (2)$$

where ν_i is the ionization rate per electron.

The requirement of a positive energy balance for the electrons becomes (Galeev, 1981)

$$v_o > \eta^{-1/2}(2eU_i/m_n)^{1/2}, \qquad (3)$$

which becomes identical to Eq.(1) when $\eta = 1$.

Eq.(2) can give a misleading picture of the electron energy balance in the interaction for several reasons. It neglects the production of ionized neutrals through charge exchange collisions, which sometimes can dominate over electron impact ionization (Axnäs, 1977). It also neglects electron energy loss through line excitation (Newell and Torbert, 1985) as well as convection of energy, both along the flow direction (which can be important when the neutral gas is limited in that direction) and through exchange with colder ambient electrons along the magnetic field. Also, one should treat the electron energy balance together with the momentum balance. This introduces a very strong dependence on the type of situation (see Fig. 1) which is under discussion.

However, the simple picture given by Eq.(2) has the advantage that it clearly shows the importance of the efficiency η of energy transfer to the electrons.

The two strongest candidates for this energy transfer are the "homogeneous model" and the "ionizing front" model. Both invoke the modified two-stream instability (MTSI) which can heat electrons efficiently when the angle θ between the wave vector and the magnetic field is close to 90°. The growth rate at these angles is of the order of the lower hybrid frequency, which is fast enough to explain the experimental observations; according to theoretical estimates, the energy transfer efficiency can be as high as 0.67 (Raadu, 1978; Galeev, 1981). Computer simulations by McBride et al. (1972) give $\eta = 0.28$. Experimental determinations of η puts it in the range 0.2-0.7 for an impact experiment (Brenning, 1982 b), and around 0.1 in a discharge experiment (Axnäs, 1986).

Although both the homogeneous model and the ionizing front model invoke the same instability, they differ in what populations drive it, and the direction of the two-streaming.

2.1. The ionizing front model

The ionizing front model was proposed by Piel, Möbius and Himmel (1980) in connection with a discharge experiment in a homopolar machine. Fig.2 shows the model in the rest frame of the plasma: the ionized neutrals are produced in an ionizing front, start with a velocity $-\underline{v}_o$ with respect to the plasma, and build up a space charge sheath with a potential $U = m_n v_o^2/(2e)$. The electrons get a secondary drift $\underline{v}_D = \underline{E}_s \times \underline{B}/B^2$ in the electric field $\underline{E}_s$ of the sheath.

The instability is driven by this secondary electron drift with respect to all the plasma ions. The highest attainable electron energy can be estimated as follows (a related argument was given by Haerendel, 1982):

The secondary electron current $\underline{i}'$ in the sheath provides the macroscopic $\underline{i} \times \underline{B}$ force that changes the ionized neutral velocity from $-\underline{v}_o$ to zero. The region of ionization should in a real plasma be distributed over the whole sheath (Piel <u>et al</u>, 1980), and not located at the edge of the sheath as in the idealized model of Fig. 2. Momentum balance within the sheath gives

$$\underline{i}' \times \underline{B} = n_e \nu_i m_n \underline{v}_o . \tag{4}$$

The macroscopic electron drift velocity transverse to $\underline{v}_o$ follows from $\underline{i}' = \underline{v}_D n_e e$:

$$v_D = v_o \nu_i m_n/eB = v_O \nu_i/\omega_{ci} . \tag{5}$$

According to Papadopoulos (1984), the instability will produce a high-energy tail with electron velocities up to

$$v_{e,max} = 2.5(m_i/m_e)^{1/2}v_D. \tag{6}$$

The highest electron energy that can be reached in the ionizing front follows from Eq.s (1), (5) and (6):

$$W_{e,max} = 6.25\ eU_i \left(\frac{v_o}{v_c} \cdot \frac{\nu_i}{\omega_{ci}}\right)^2 \tag{7}$$

Thus, electron heating to above the ionization energy is only possible in the ionizing front model if

$$\nu_i/\omega_{ci} > 0.4v_c/v_o \tag{8}$$

2.2. <u>The homogeneous model</u>

The homogeneous model was proposed by Sherman (1969) and further developed by Raadu (1978, 1982). In this model the interaction involves only the electrons and the newly ionized neutrals, which in the rest frame of the plasma have a velocity close to $-v_O$.

The limit of applicability of this model can be estimated in two ways. In a plasma flow with ionization, the transverse scale length is typically $a \approx v_o/\nu_i$. Mikhailovskii (1974) gives a limit of the applicability for the MTSI in a plasma with a density gradient:

$$v_o < \omega_{ci}a \approx \frac{\omega_{ci}v_o}{\nu_i} \tag{9}$$

which gives a limit for the homogeneous model:

$$\nu_i / \omega_{ci} < 1. \tag{10}$$

This condition can also be derived directly without introducing the transverse scale length: According to Eq. (5), the transverse electron drift velocity v_D associated with the pick-up of new ions becomes greater than the original velocity v_o when $\nu_i / \omega_{ci} > 1$. In this situation, the homogeneous model (which disregards this secondary drift) is obviously unsuitable.

3. A lower limit on ν_i/ω_{ci} for efficient electron heating

High values of the energy transfer factor η are possible only if the ionized neutrals are prevented from forming a gyrotropic distribution in velocity space. If such a distribution is formed, the efficiency becomes only 2.5% according to quasilinear calculations (Galeev, 1981; Formisano *et al.*, 1982). According to Eq. (2) the CIV interaction would then require a threshold velocity of $\eta^{-1/2} v_c = 6.3\ v_c$.

The condition for a beam distribution of ionized neutrals used by the authors above was

$$\nu_i / \omega_{ci} > 1 \tag{11}$$

They also give the condition

$$\nu_i/\omega_{ci} < (m_e/m_i)^{1/2} \tag{12}$$

for the formation of a gyrotropic distribution in velocity space, and expect η to vary smoothly from 0.67 to 0.025 as ν_i / ω_{ci} decreases from 1 to $(m_e/m_i)^{1/2}$.

However, these results apply only to the steady state situation, in which the new ions have to compete with an accumulated hot (gyrotropic) background of earlier ionized neutrals. In the initial phase of a transient situation, this background will be small, and the energy transfer can be more efficient. We will here treat the case where ionization starts at t=0, and the initial background of hot ions is zero.

The instability is for this case that of a thin beam in a plasma, with a growth rate that increases with the beam density; it is of the order $\omega_{LH}(n_{beam}/n_e)^{1/3}$. With ionization starting at t=0, the question is whether the instability growth time ever becomes much faster then the gyrotime. In this case, a ring-type ion distribution would never have time to develop.

The growth rate for a thin beam in a plasma depends strongly on the angle between the wave vector and the magnetic field. When the electrons are hot enough to ionize, only the modes with almost perpendicular propagation remain, with a growth rate (see the Appendix):

$$\gamma = \frac{3^{1/2}}{2^{4/3}} \omega_{LH} \left(\frac{\beta m_{ip}}{m_{is}} \right)^{1/3} \tag{13}$$

where $\beta = (n_s/n_e)$ is the fractional beam density, and

$$\omega_{LH} = \omega_{pi}(1+ \omega_{pe}^{2} / \omega_{ce}^{2})^{-1/2} \tag{14}$$

is the lower hybrid frequency of the plasma, calculated with the "main plasma" ion mass (i.e. not the mass of the ionized neutrals). m_{ip} and m_{is} are the ion masses of the background plasme and the beam, respectively.

A crucial parameter is the time (the "energy transfer time" t_{ET}) that is required to smear out the beam and transfer the energy to the electrons. It is an unknown factor κ larger than the growth time γ^{-1} of the instability:

$$t_{ET} = \kappa/\gamma . \tag{15}$$

The condition for preventing a ring-type distribution becomes

$$t_{ET} \ll \omega_{cin}^{-1}, \tag{16}$$

where ω_{cin} is the gyrofrequency of the ionized neutrals. Since the ion beam is smeared out on the time t_{ET}, the fractional beam density is given by the number of ionizations during this time:

$$\beta = \nu_i t_{ET}. \tag{17}$$

The condition (16) can then be rewritten into a condition on ν_i/ω_{cin}:

$$\frac{\nu_i}{\omega_{cin}} \gg \left(\frac{m_e}{m_n}\right)^{3/2} \cdot \left(\frac{m_{ip}}{m_n}\right)^{1/2} \cdot \left(\frac{\omega_{ce}^2}{\omega_{pe}^2} + 1\right)^{3/2} \cdot \frac{2^4\kappa^3}{3^{3/2}} \tag{18}$$

Most experiments and proposed applications of the critical velocity phenomenon are in a parameter range where $\omega_{pe}^2/\omega_{ce}^2 > 1$ (Brenning, 1982 b). When this is true, Eq. (18) can for practical purposes be written:

$$\frac{\nu_i}{\omega_{ci}} \gg \left(\frac{m_e}{m_n}\right)^{3/2} \kappa^3 \tag{19}$$

(This condition was independently derived by Haerendel, 1986).

It is of some interest to discuss more specifically what the $\gg$ sign in (18) means. Fig. 3 illustrates that the ions will retain much of their beam character in the direction of $\underline{v}_o$, even if the energy transfer time is as long as 0.5 ω_{ci}^{-1}. With this as the limit, Eq. (16) becomes $t_{ET} < 0.5\ \omega_{ci}^{-1}$, and (19) becomes

$$\frac{\nu_i}{\omega_{ci}} > 2 \left(\frac{m_e}{m_n} \right)^{3/2} \kappa^3 \qquad (20)$$

This is the condition for efficient electron heating in transient situations. The number κ of growth times needed for relaxation of the beam ion distribution is a very uncertain parameter. An experimental determination from a discharge-type critical velocity experiment (Axnäs, 1986) only gives an upper limit: $\kappa < 20$. Computer simulations by McBride et al. (1972) gave $\kappa \approx 6$, while Tanaka and Papadopoulos (1983) in similar simulations obtain $t_H = 25 \cdot \omega_{LH}^{-1} = 12.5\ \gamma^{-1}$, i.e., $\kappa = 12.5$. (These simulations were made for the case where all ions drift with respect the the electrons, which gives $\gamma \approx \omega_{LH}/2$ for the modes with $\cos\theta = (m_e/m_i)^{1/2}$.)

In spite of this uncertainty of κ, it is clear that efficient electron heating, during the initial phase of a transient interaction, is possible for values of ν_i/ω_{ci} far below those derived by e.g. Galeev (1981) for steady-state interaction.

4. Summary and Discussion

The value of ν_i/ω_{ci}, which can also be seen as a magnetic field strength/neutral density relation, is important for CIV interaction. In this context, ν_i is the total rate of production of quasistationary ions. It is immaterial both for the momentum balance, (which was the key to the discussion of Section 2), and to the instability growth (Section 3) if this production is due to electron impact ionization or some other process, e.g. charge exchange or photoionization. We will here assume that all these process are included in ν_i. In a "transient" critical velocity experiment, ν_i/ω_{ci} splits up the interaction into different regions, as scematically illustrated in Fig. 4.

<u>High ionization rate</u>; $\nu_i/\omega_{ci} > 1$: The "homogeneous model" of Sherman (1969) and Raadu (1978) is not applicable. The "ionizing front model" by Piel <u>et al</u>. (1980) can give efficient electron heating. Alternatively, the lower hybrid drift instability could act as energy-transferring mechanism.

<u>Intermediate ionization rate</u>; $(m_e/m_n)^{3/2}\kappa^3 < \nu_i/\omega_{ci} < 1$: The ionizing front model gives too poor electron heating to be a candidate for electron heating. The "homogeneous model" is applicable. (1): In a transient situation, the MTSI can give efficient energy transfer ($\eta \approx 0.67$) down to approximately $\nu_i/\omega_{ci} = 2(m_e/m_n)^{3/2}\kappa^3$. For lower ionization rates, the value should gradually decrease towards the constant low value in the region of low ionization. (2): In a steady state situation, the η value begins to decrease immediately when ν_i/ω_{ci} decreases below unity. The region of low ionization rate is reached at $\nu_i/\omega_{ci} < (m_e/m_i)^{1/2}$ (Galeev, 1981).

<u>Low ionization rate</u>; $\nu_i/\omega_{ci} << (m_e/m_n)^{3/2}\kappa^3$: The MTSI has too slow growth to prevent the formation of a gyrotropic distribution of the ionized neutrals, and the energy transfer factor becomes low, $\eta = 0.025$, even in a transient situation.

This division into three regions contains the very uncertain parameter κ , the number of growth times required for the energy transfer from the ionized neutrals to the electrons. Computer simulations indicate that κ lies in the range 6 - 12.5.

The main conclusion here is that efficient energy transfer is possible in transient situations, even if ν_i/ω_{ci} is far below unity. This result strictly applies only to a fraction of the first gyrotime after ionization has begun. At later times, the ion distribution consists of a mixture of the newly ionized neutrals, which form a "beam" and can give efficient electron heating, and the ions that have had time to form a isotropic distribution perpendicular to $\underline{B}$.

In this situation, the results above will be modified. However, it is not immediately obvious in what direction the modification will go. On one hand, the beam will be formed in an already large fluctuating electric field; the number κ of growth times required to transfer the energy to the electrons can then be substantially reduced. This has a large influence in Eq:s (18)-(20), where κ appears to the third power.

On the other hand, the efficiency η and the growth rate of the instability will be reduced by the gyrotropic background of ions with velocities comparable to the ion beam velocity. This would be increasingly important as time goes on, and the background density grows. It is difficult to estimate where the transition to the steady-state case will come. In this context, it would be desirable with computer simulations of the MTSI for the case where the ion beam density grows from zero.

Acknowledgements: I want to thank Professor C. -G. Fälthammar, and Drs I. Axnäs and M. A. Raadu for many stimulating discussions. This work has been financed by the Swedish Natural Science Research Council.

Appendix

The linear growth rate of the MTSI with a transverse ion stream in a cold, homogeneous plasma

The dispersion relation for the electrostatic modes is obtained by putting the plasma dielectric constant equal to zero. For the case of an ion stream perpendicular to $\underline{B}$ in a uniform, collisionless plasma with non-magnetized ions, $\omega \gg \omega_{ci}$, this gives

$$1 - \frac{\omega_{pe}^2 \cos^2\theta}{\omega^2} - \frac{\omega_{pe}^2 \sin^2\theta}{\omega^2 - \omega_{ce}^2} - \frac{(1-\beta)\omega_{pi}^2}{\omega^2} - \frac{\beta(m_{ip}/m_{is})\omega_{pi}^2}{(\omega - \underline{k}\cdot\underline{v})^2} = 0, \quad (I)$$

where θ is the angle between the wave vector and the magnetic field, m_{ip} is the plasma ion mass, and m_{is} is the ion mass in the stream. The stream has velocity v and fractional density $\beta = n_{is}/n_e$. Here, ω_{pi} is the ion plasma frequency in the absence of a stream: $\omega_{pi} = ((n_e e^2)/(\varepsilon_o m_{pi}))^{1/2}$.

The electrostatic oscillations of the plasma at rest are found by omitting the last term in Eq.(I). They are

$$\omega_{1,2}^2 = \omega_{pi}^2 \frac{m_{ip}}{2m_e}\left(\frac{1}{\alpha} + 1 + \frac{m_e}{m_{ip}}(1-\beta) \pm \left\{\left(\frac{1}{\alpha} + 1 + \frac{m_e}{m_{ip}}(1-\beta)\right)^2 - \frac{4\cos^2\theta}{\alpha} - 4\frac{m_e}{m_{ip}}\cdot\left(\frac{1-\beta}{\alpha}\right)\right\}^{1/2}\right), \quad (II)$$

where $\alpha = \omega_{pe}^2/\omega_{ce}^2$.

The higher of these frequencies, ω_1, lies above $\max(\omega_{pe}, \omega_{ce})$, and the lower, ω_2, below $\min(\omega_{pe}, \omega_{ce})$. For almost perpendicular propagation, ω_1 approaches the upper hybrid frequency $\omega_{UH} = (\omega_{pe}^2 + \omega_{ce}^2)^{1/2}$, while ω_2 approaches the lower hybrid frequency (Eq. 14). We are here only interested in the low frequency branch, which corresponds to the minus sign in Eq.(II).

For a thin beam ($\beta \ll 1$) the phase velocity of the waves will be close to the beam velocity v for all values of Θ. The resonant wavenumber is then given by

$$\omega_2 = k_\perp v. \tag{III}$$

The corresponding growth rate is derived from Eq:s (I)-(III) in the usual way for a beam-plasma instability. It is shown in Figs (5) and (6). An important parameter is the "equal effective mass angle" θ_o, given by $\cos\theta_o = (m_e/m_{ip})^{1/2}$.

For $\theta \ll \theta_o$, the plasma ions take no part in the interaction, and the growth rate therefore is independent of m_{ip}. The instability for $\theta \ll \theta_o$ is a mixture of the modified two-stream instability and the beam cyclotron instability, which plays a role for $\theta < 60^o$, but only when $\omega_{pe}/\omega_{ge} > 1$.

An analytical solution of the dispersion relation can be obtained under the assumption $\omega \ll \omega_{ce}$. This gives the real frequency in the ion stream frame

$$\mathrm{Re}\,\omega_{(stream)} = \omega_{pe} K \tag{IV}$$

and the growth rate

$$\gamma = 3^{1/2} \omega_{pe} K \tag{V}$$

where

$$K = \frac{(\cos^2\theta + m_e/m_{ip})^{1/6} (\beta m_e/m_{is})^{1/3}}{2^{4/3}(1 + \sin^2\theta(\omega_{pe}/\omega_{ge})^2)^{1/2}} \tag{VI}$$

The approximate growth rate from Eq. (V) is shown with dotted lines in Fig (5) for comparison with the exact results from the solution of Eq:s I -III. It is found to be an acceptable approximation (within 10%) for $\theta > 70^{\circ}$. Since the modes for smaller angles are damped at rather low electron temperatures, Eq. (V) is a good enough approximation for a discussion of the critical velocity interaction, where electron energies of the order of the ionization energy eU_i are needed.

The lowest growth rate is found at $\theta = 90^{\circ}$. It is

$$\gamma_{min} = \frac{3^{1/2}}{2^{4/3}} \omega_{LH} \left(\beta \frac{m_{ip}}{m_{is}} \right)^{1/3} \qquad \text{(VII)}$$

Figure captions

Fig.1. The geometrical and temporal differences between some proposed cases of CIV interaction.

Fig.2. The ionizing front model according to Piel et al (1980).

Fig.3. Illustration that the new ions retain their beam character in the direction of $\underline{v}_o$, even if the instability takes the time 0.5 ω_{ci}^{-1} to tap their energy.

Fig.4. A schematic illustration of how the energy transfer efficiency η can be expected to vary with ν_i / ω_{ce}. $_i$ is here the total production of quasistationary neutrals, including electron impact ionization, charge exchange and photoionization. κ is the number of growth times it takes the instability to tap the ion energy (the quasilinear relaxation time in units of growth times). Upper curve: the initial phase of a transient interaction. Lower curve: homogeneous steady-state interaction (schematically from Galeev, 1981).

Fig.5. Solid line: growth rate of the (cold) lower hybrid instabilities with an ion stream of fractional density $\beta = n_{is}/n_e$ transverse to the magnetic field. m_{is} = stream ion mass; m_{ip} = plasma ion mass. The curves are calculated for a hydrogen plasma. For angles $\theta > 85^o$, there is a weak additional dependence on the the ion mass m_{ip}, which is shown in the next figure. Dotted line: The approximate growth rate of Eq. (V), which is a good analytical approximation when $\theta > 70^o$.

Fig.6. The growth rate for almost perpendicular propagation, where it is influenced by the plasma ion mass m_{ip}. The growth rate at $\theta = 90^o$ is expressed by Eq. (VII).

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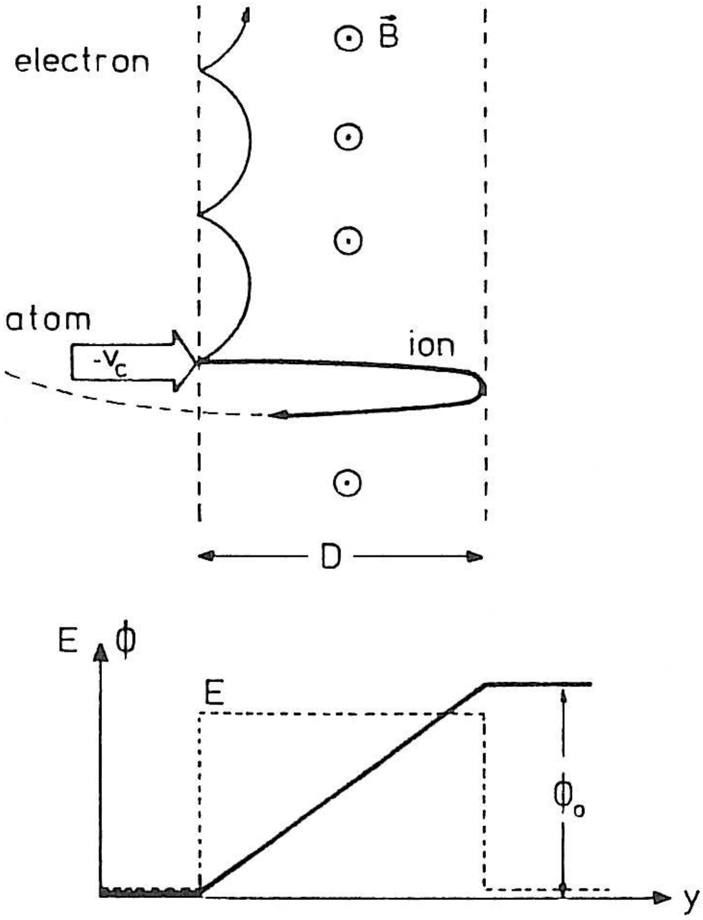
Time duration → / Spatial extension ↓	A few ion transit times (or ion gyro times)	More than ten ion transit times (or ion gyro times)	Steady state
Limited plasma, limited neutral gas.		Laboratory impact experiments.	Laboratory ExB discharge experiments.
Limited neutral gas, larger plasma.	Ionospheric experiments with shaped charge explosions.	Proposed Ionospheric releases from the Space shuttle.	Io plasma torus; Space shuttle glow; Comet interaction with solar wind.
Large plasma and neutral gas.			Alfvéns solar system theory; Solar wind - interstellar gas interaction.

Fig.1. The geometrical and temporal differences between some proposed cases of CIV interaction.



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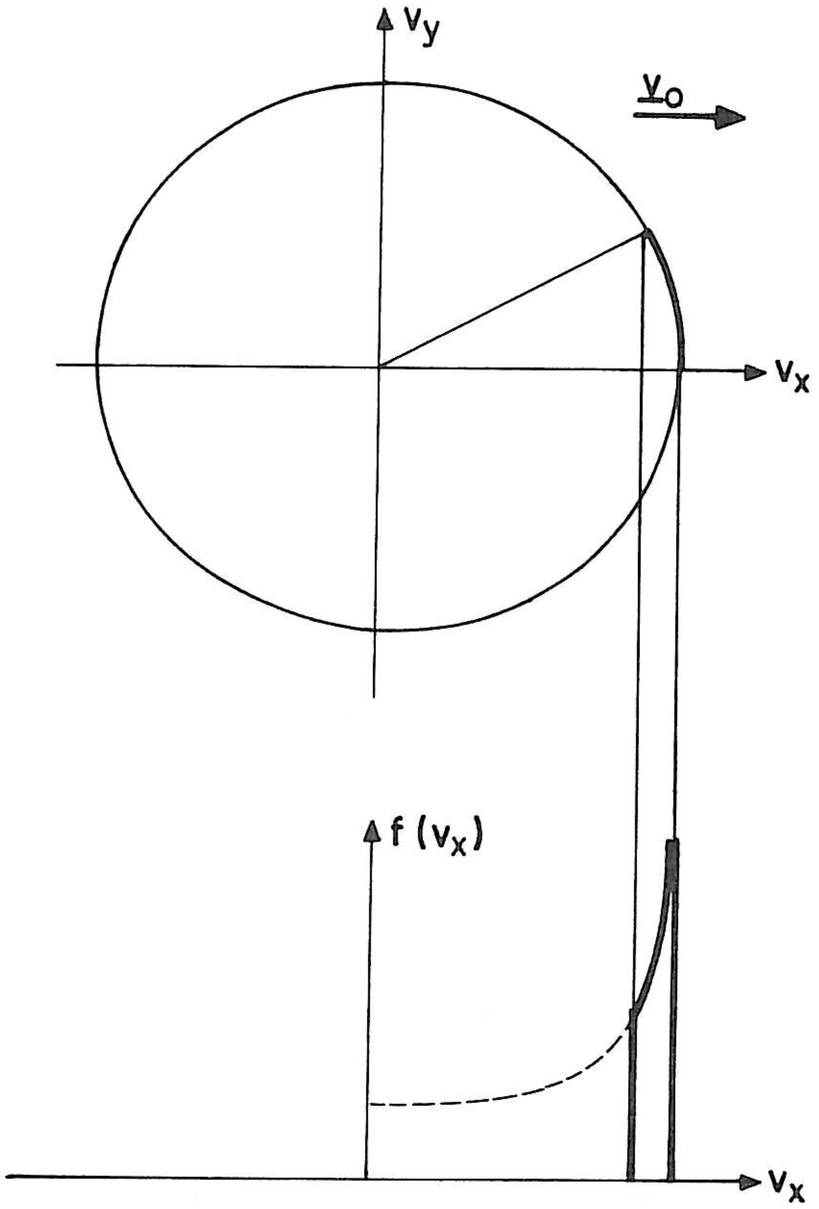

Fig.3. Illustration that the new ions retain their beam character in the direction of $\underline{v}_o$, even if the instability takes the time 0.5 ω_{ci}^{-1} to tap their energy.

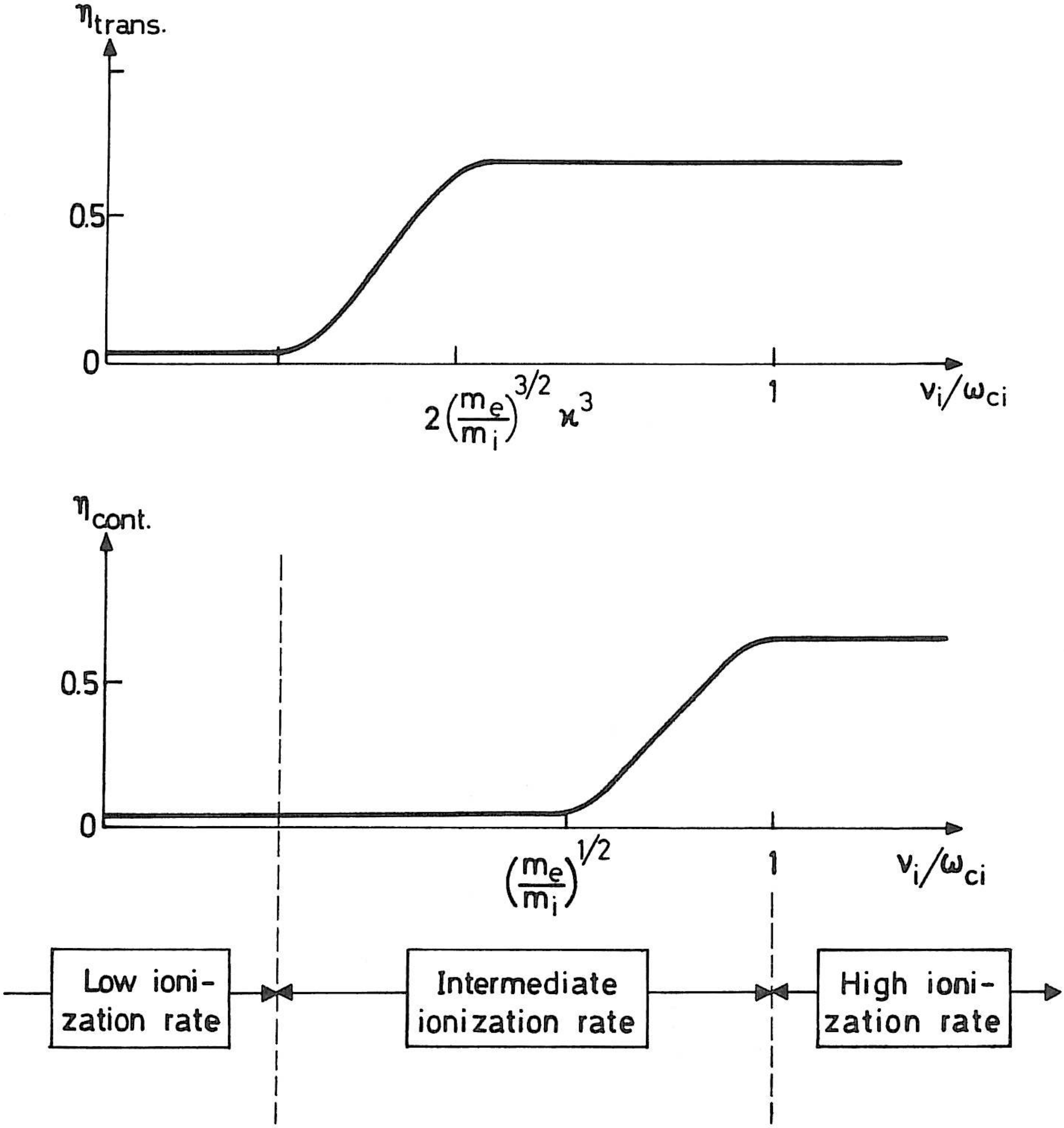

Fig.4. A schematic illustration of how the energy transfer efficiency η can be expected to vary with ν_i/ω_{ce}. $_i$ is here the total production of quasistationary neutrals, including electron impact ionization, charge exchange and photoionization. κ is the number of growth times it takes the instability to tap the ion energy (the quasilinear relaxation time in units of growth times). Upper curve: the initial phase of a transient interaction. Lower curve: homogeneous steady-state interaction (schematically from Galeev, 1981).

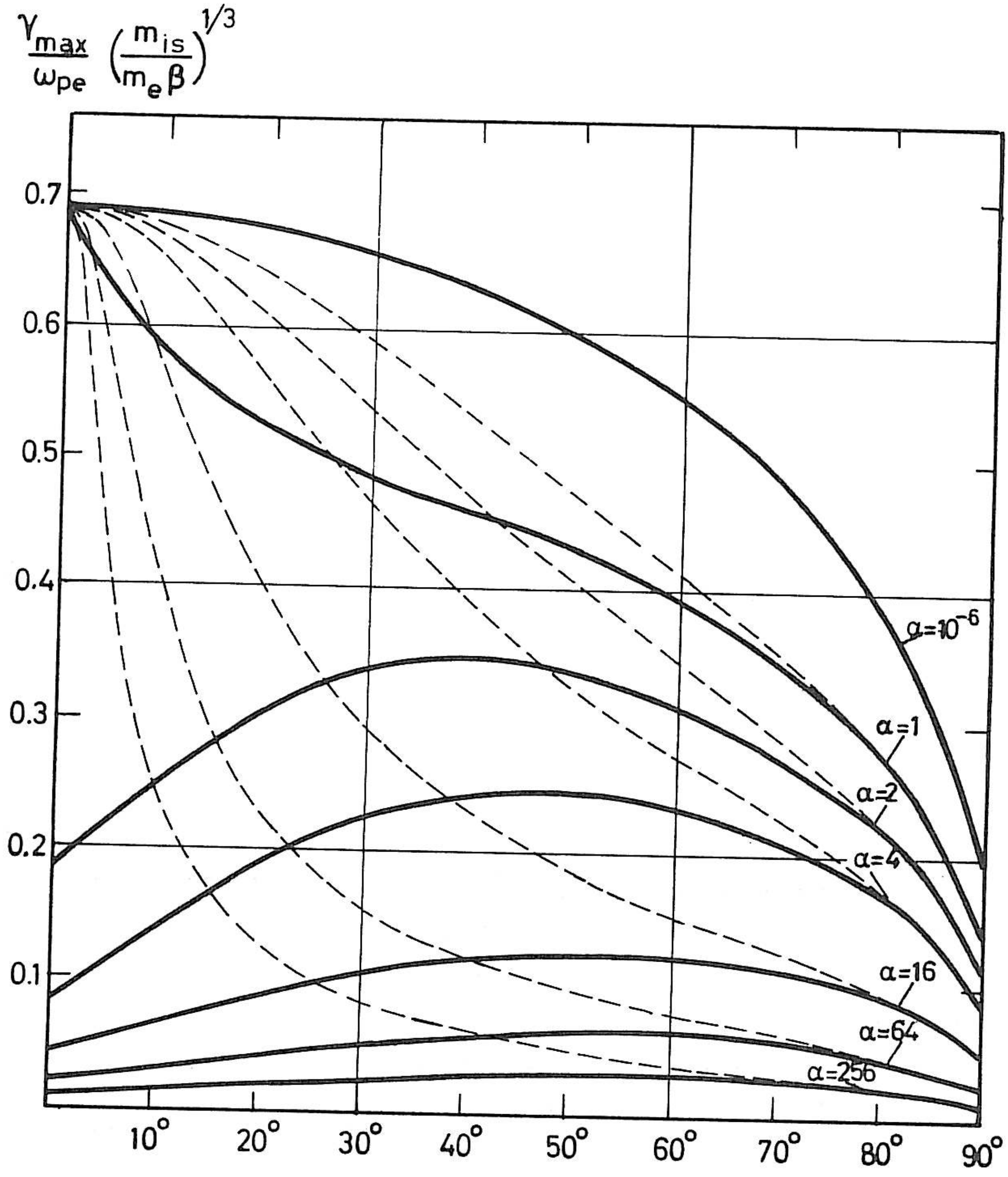

Fig.5. Solid line: growth rate of the (cold) lower hybrid instabilities with an ion stream of fractional density $\beta = n_{is}/n_e$ transverse to the magnetic field. m_{is} = stream ion mass; m_{ip} = plasma ion mass. The curves are calculated for a hydrogen plasma. For angles $\theta > 85^{\circ}$, there is a weak additional dependence on the the ion mass m_{ip}, which is shown in the next figure. Dotted line: The approximate growth rate of Eq. (V), which is a good analytical approximation when $\theta > 70^{\circ}$.

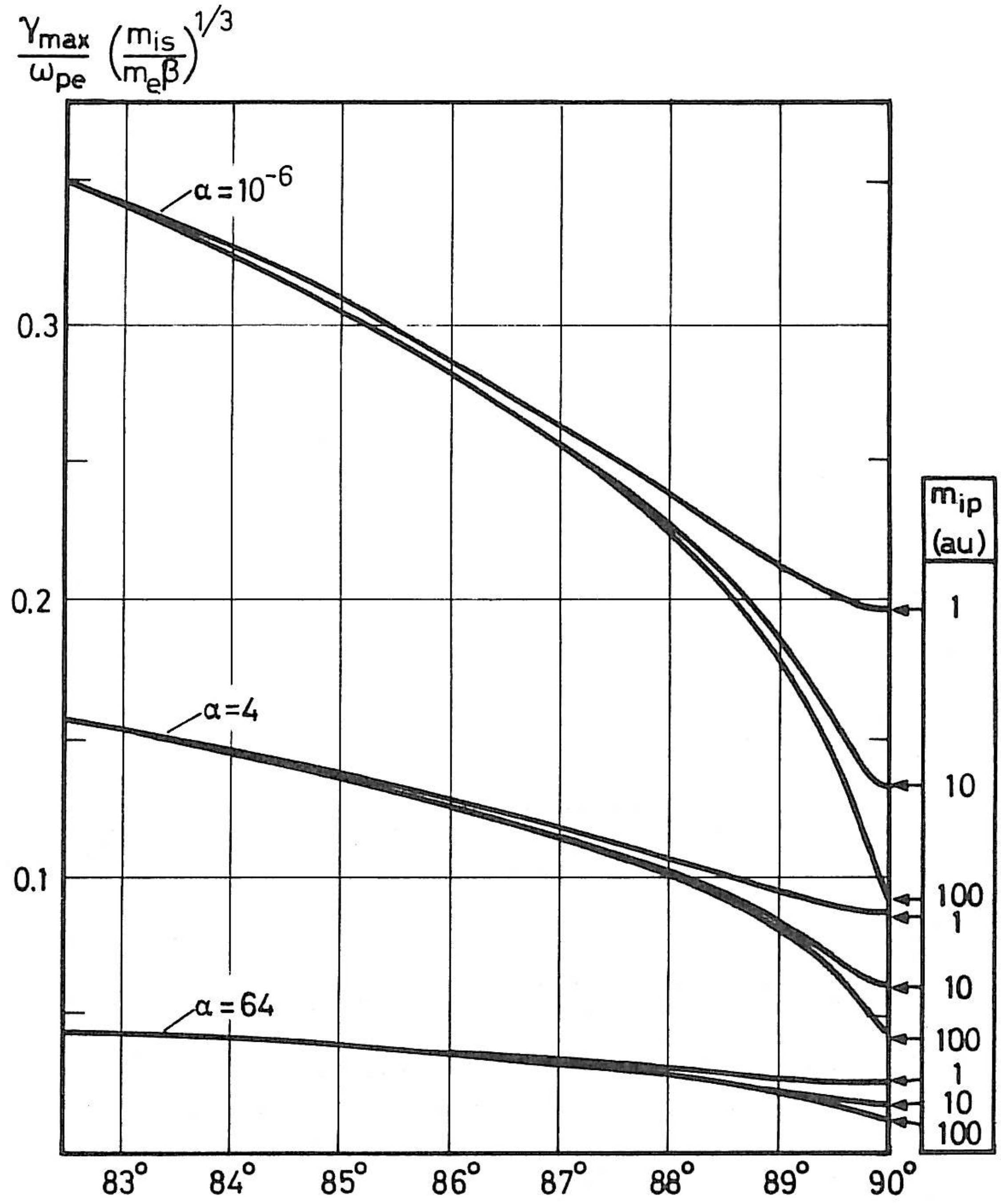

Fig.6. The growth rate for almost perpendicular propagation, where it is influenced by the plasma ion mass m_{ip}. The growth rate at $\theta = 90^{\circ}$ is expressed by Eq. (VII).

TRITA-EPP-86-10 ISSN 0348-7539

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ON THE ROLE OF THE IONIZATION FREQUENCY TO GYROFREQUENCY RATIO IN THE CRITICAL IONIZATION VELOCITY INTERACTION

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October 1986, 25 pp. incl. ill., in English

Abstract

The role of the parameter ν_i/ω_{ci} for critical ionization velocity (CIV) interaction is discussed. This parameter, which can be seen as a combined condition on the neutral gas density and the magnetic field strength, is important from two different aspects. First, the value $\nu_i/\omega_{ci} = 1$ marks the limit between the regions of applicability of the homogeneous and the "ionizing front" (or inhomogeneous) models for the interaction. Second, it determines the energy transfer efficiency in the interaction. This energy transfer efficiency also depends strongly on whether the experiment is steady-state or transient in nature. In steady-state situations, the interaction has been estimated to become gradually weaker when ν_i/ω_{ci} decreases below unity. It is found here that efficient CIV interaction is possible at far lower values of ν_i/ω_{ci} in transient experiments. The corresponding limit is approximately $\nu_i/\omega_{ci} = 2(m_e/m_i)^{3/2}\kappa^3$. Unfortunately, κ is a rather uncertain parameter: the number of instability growth times required for the modified two-stream instability to tap the energy of a thin ion stream in a plasma.

Key words: Critical Velocity, Critical Ionization Velocity, Alfvéns Critical Velocity.