ROLE OF THE SECONDARY ELECTRON EMISSION IN INTERACTION OF INTENSIVE ION BEAMS WITH METAL OR DENSE PLASMA TARGETS

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Résumé - Le problème de l'interaction de faisceaux ioniques et intenses neutralisés, avec un métal isolé ou des plasmas denses est examiné à l'aide de l'eq. de Boltzmann avec des conditions aux limites prenant en compte l'émission d'électrons secondaires. Dans le cadre de la théorie cinétique, l'émission d'électrons secondaires et la rétrodiffusion des électrons accompagnant les ions peuvent provoquer un excès important de charges sur la cible isolée. De tels effets peuvent jouer un rôle important dans le cas d'une microsphère située dans une chambre de réactions remplie de gaz dense et irradiée par des faisceaux d'ions légers.

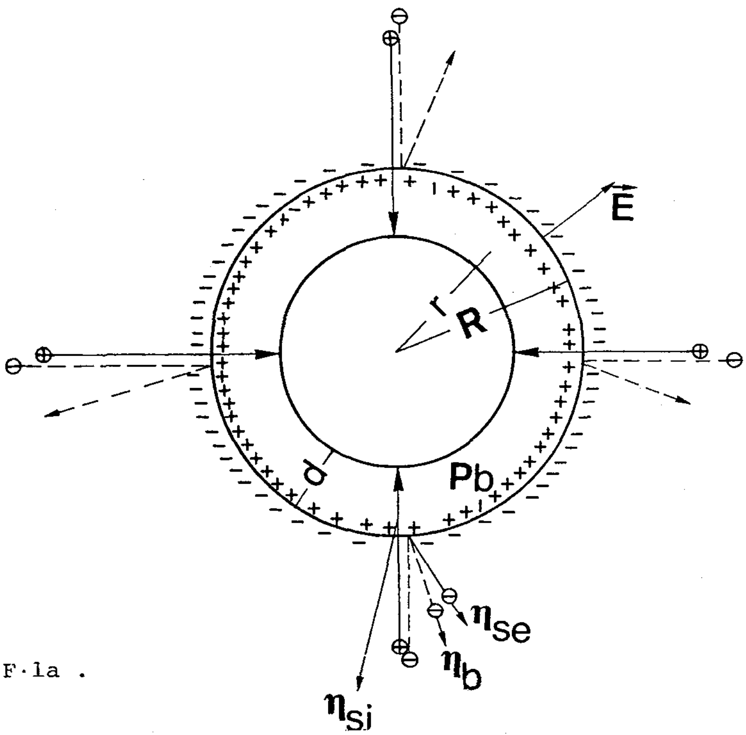
Abstract

A general problem of interaction of intensive neutralized ion beams with isolated metal or dense plasma targets is investigated with the help of Boltzmann equation involving appropriate boundary conditions taking into account emission of secondary electrons provoked by these beams. It is shown in the framework of this kinetic theory that the emission of the secondary electrons and back - scattering of the electrons moving with the ion beams can lead to an appreciable loading of an isolated target. These loading effects can be particularly important for the case of an ICF pellet situated in a reaction chamber filled by a dense gas and irradiated by beams of light ions.

1. Introduction

According to the modern light ion ICF-concept (see e.g. Mankofsky, Sudan, /1/) the ions must enter the pellet situated in a reaction chamber filled by a dense gas through a number of Z-pinch channels. The ion beams are known to induce return currents of electrons in plasma in the Z-pinch channels and in the reaction chamber, i.e. they can be considered as current- neutralized. The current density of the ion beam $j_b=n_b v_b$ is equal, hence, to a good accuracy to density $j_p=n_p v_d$ of the plasma electrons (see e.g. Mankofsky, Sudan, /1/). Here n_b, n_p are the densities of the beam ions and the plasma electrons and v_b, v_d are resp. the beam ion velocity and the drift velocity of the plasma electrons, connected, with $j_b=n_b v_b$ by the relation $v_d=n_b v_b/n_p$.

The equalty of the current of the beam ions and of plasma electron in the gas surrounding pellet does not mean, however, that the currents of positive and negative charges entering the pellet are always equal. The beam ions entering the pellet can provoke a substantial emission of secondary electrons. According e.g. to (Little, /2/) an ion of MeV energy can eject a few (η_{si}) of secondary electrons while striking a metal surface. A plasma electron entering the target to-



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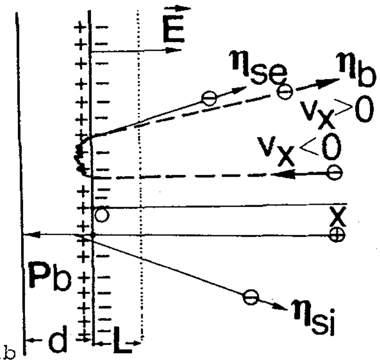

F.1b

Fig. 1a . Schematic representation of situation arising around of the pellet as a result of emission of secondary electrons.

Fig.1b. Slab- model of the surface processes occuring nearby the pellet

gether with beam ions can also provoke a number η_{se} of secondary electrons. And, finally, the plasma electrons drifting together with beam ions can be partly backscattered from the metal surface, with back- scaattering coefienent η_b (ratio of the reflected electron current to the incoming electron current) reaching about 0,5 (Little, 2). Thus, with the start of the ion pulse the total particle flux of the positive charges entering the pellet $j^+=n_b v_b(1+\eta_{si})+n_p v_d \eta_{se}$ exceeds essentialy the flux of negative charges $j^-=n_p v_d(1-\eta_b)$ entering there. This leads to appearance of the positive charge in the pellet and to accumulation of the electrons in the narrow plasma layer surrounding the pellet, i.e. to appearance of the double layer within which strong electric fields can be formed. A crude estimate of the charges and fields E arising in this way can be performed as follows. The electron density at the pellet surface will increase due to the surface effects reaching a value n_p^* corresponding to dynamical equilibrium $j^+=j^-$, i.e. $n_b v_b(i+\eta_{si})=n_p^* v_d(1-\eta_{se}-\eta_b)$.(We neglect here the flux of plasma ions which is much smaller than $n_b v_b$). It gives following estimate for stationary additional electron density $n_o=n_p^*-n_p$ at the pellet surface due these surface effects

$$n_o=n_p\alpha=n_p\,\frac{\eta_{si}+\eta_{se}+\eta_b}{1-\eta_{si}-\eta_{se}-\eta_b}=\frac{j_b}{v_d}\,\frac{\eta_{si}+\eta_{se}+\eta_b}{1-\eta_{si}-\eta_{se}-\eta_b} \qquad (1)$$

Total number N of additional electrons per cm^2 can be expressed by the formula $N=n_o r_d$ where n_o is given by (1) and r_d is the thickness of the double layer which is equal approximately to the plasma Debye- radius:

$$r_d=(u_e/\omega_p\sqrt{2})=(T/4\pi n_p e^2)^{1/2},\quad u_e=(2T/m_e)^{1/2},\quad \omega_p=(4\pi n_p e^2)/m_e)^{1/2}$$

Here T is the plasma temperature and m_e is the electron mass. Note, that we have assumed above that the beam ions are not disturbed by the pellet. This is connected with the fact that the scattering angle of the beam ions in the metal shell of the pellet is very small and, hence, the ions do not experience any apreciable back scattering. In the equilibrium the number of positive charges in the pellet is equal to N as well. The field E arising at the pellet surface is, hence, given by $E=4\pi e n_o r_d$. Assuming for the proton beam and Pb-tamper shell $\eta_{se}=1$, $\eta_{se}=0,4$, $\eta_b=0,2$ (see below), and the proton energy W_b=5MeV we get $E\approx 10^8$V/cm for the plasma temperature T=50eV and $n_p=10^{18}cm^{-3}$. (This example is taken from (Mankofsky, Sudan /1/)).

These results were formulated for an idealised case of homogeneous irradiation of a target by an ion current of density $j_b=n_b v_b$. To extend them for the realistic case we will use in the following amodel of quasihomogeneous irradiation of a pellet of radius R meaning under $j_b=n_b v_b$ total current J due to a number ot ion beams, incoming onto the pellet, divided by the pellet surface area $4\pi R^2$: $j_b=\frac{J}{4\pi R^2}$. Accordingly we will use the model of a spherically symmetric double layer surrounding the pellet. This model seems to be not too far from reality because the electric field arising at the highly conducting pellet must be necessarily normal to its surface and constant along it. According to this model the total charge Q accumulated in the pellet of radius R can be represented by the formula

$$Q = R^2E = 4\pi R^2 e n_o r_d = e\frac{J}{\omega_p}\frac{u_e}{v_d\sqrt{2}}\cdot\frac{\eta_{si}+\eta_{se}+\eta_e}{1-\eta_{si}-\eta_{se}-\eta_e} \qquad (2)$$

which gives a new insight into the phenomenon of the pellet loading provoked by a neutralized ion beam: the equilibrium state characterised by equality of fluxes of positive and negative charges entering the pellet arises not immidiately after start of the ion pulse but within a plasma time $1/\omega_p$, (for $u_e\approx v_d$).

During this time an accumulation of the positive charges in the pellet due to the surface effects take place, which proves, to be, naturally, of the order of $Q\approx eJ/\omega_p$. For very high current J the charges Q in the pellet can be appreciable, notwithstanding the smallness of $1/\omega_p$. The aim of this paper is to give estimates of the charges and fields arising as a result of these effects in isolated targets and to analyse their implications for generally adopted ICF concept based on the idea of irradiation of a freely falling pellet by beams of light ions.

2. Electric fields outside of the pellet

According to model outlined above we can consider the problem of calculation of distribution of charges and electric fields concentrated in a narrow layer of thickness r_d outside of the pellet as a one dimensional one. The electron distribution function $f_e(\vec{v},x,t)$ for the electron component outside of the pellet must depend, hence, only on the coordinate x counted perpendicular to the pellet surface. In the stationary state $f_e(\vec{v},x,t)$ must obey the Boltzmann equation

$$\frac{\partial f_e}{\partial t} + v_x \frac{\partial f_e}{\partial x} - \frac{eE}{m_e}\frac{\partial f_e}{\partial v_x} = St(f_e) \tag{3}$$

where $St(f_e)$ is a collsion intergral and $E=-\partial\varphi/\partial x$ is a static electric field arising outside of the pellet.
We will start our analysis with case of small coefficients η_{si}, η_{se}, η_e, corresponding to an assumption that the surface effects influence only slightly the electron distribution function outside of the pellet. It means also that, in the first approximation, the ion beam remains current-neutralized both in the reaction chamber and in the immediate vicinity of the pellet. Thus, we assume that the electron distribution function is, in the first approximation, Maxwellian everywhere outside of the pellet

$$f_e \approx f_{eo} = \frac{n_e}{\pi^{3/2}u_e^3}\exp\left[-\frac{(v_x-v_d)^2+v_\perp^2}{u_e^2}\right], \quad u_e^2 = \frac{2T}{m_e} \tag{4}$$

where $v_\perp$ is velocity of the plasma electrons directed perpendicular to the ion beam and v_x is velocity of the electrons directed parallel to the ion beam. As a result of the surface effects the distribution (4) will be slightly disturbed nearby the pellet and can be, hence, represented in the form $f_e=f_{eo}+f$ where $f \ll f_{eo}$ is a small perturbation satisfying the Boltzmann equation

$$\frac{\partial f}{\partial t} + v_x \frac{\partial f}{\partial x} - \frac{eE}{m_e}\frac{\partial f_{eo}}{\partial v_x} = -\nu f, \tag{5}$$

ν being an electron collision frequency. The beam ion distribution function f_b can be assumed to be undisturbed by the target as explained in the Introduction.

Accordingly we can use the Poisson equation in the form

$$\operatorname{div} \vec{E} = \frac{\partial^2 \varphi}{\partial x^2} = 4\pi e(n_b + n_i - \int f e \, d\vec{v}) - E(o)\ \delta^+(x) \tag{6}$$

Here n_i is the plasma ion density and the term $E(o)\ \delta^+(x)$ $(\int_0^\infty \delta^+(x)dx=1)$ takes into account the presence of the positive charge at the pellet surface. The plasma ions are repelled from the positively charged pellet and, hence, we can assume that the perturbation of the ion density due to the electric field appearing at the pellet obeyes the Boltzmannian law, i.e. one can write $n_i=n_{io}+n_{io}\frac{e\varphi}{T}$ where n_{io} is an unperturbed plasma ion density. (A more defaited analysis shows that this assumption is valid if $u_i=\sqrt{T/m_i}<v_d$, m_i being the plasma ion mass). Using this relation, the definition$f_e=f_{eo}+f$ and the condition of the plasma quasineutrality $n_b+n_{io}= n_p=\int f_{eo}\, d\vec{v}$ we reduce (6) to the form

$$\operatorname{div} \vec{E} = \frac{\partial^2 \varphi}{\partial x^2} = -4\pi \int f d\vec{v} - E(o)\ \delta^+(x) + \chi_i^2 \varphi \tag{7}$$

where

$$\chi_i^2 = 4\pi n_{io} e^2/T. \tag{8}$$

To solve the system (5),(7) we need boundary conditions at the pellet surface taking into account the emission of the secondary electrons and backscattering of the electrons drifting with the ion beam to the surface. Similar boundary conditions were formulated in (Alpert, et al, /3/) in connection with theory of disturbances provoked in the ionosphere plasma by an Earth' satellite. In our case the boundary conditions can be represented in the form

$$v_x f(\vec{v},o)=\int w_b(\vec{v},\vec{v}')v_x' f_{eo}(v')d\vec{v}'+\int w_{se}(\vec{v},\vec{v}')v_x' f_{eo}(v')d\vec{v}'+\int w_{si}(\vec{v},\vec{v}')v_x' f_b(v')d\vec{v}' \tag{9}$$

where f_{eo} is given by (4) and w_b is the probability for an electron (having the velocity $\vec{v}'$ and moving with the ion beam) to acquire after the back-scattering the velocity $\vec{v}$; w_{se} is the probability that this electron will emit a secondary electron having the velocity $\vec{v}$ and w_{si} is the probability for an beam ion having the velocity $\vec{v}'$ to emit a secondary electron having the velocity $\vec{v}$. The probability w_{si} (and, similarly, w_{se}, w_b) has the property $\int w_{si}(\vec{v}, \vec{v}')\, d\vec{v} = \eta_{si}(\vec{v}')$. The function $f(\vec{v},o)$ is defined by (9) for $v_x>0$ i.e. for electrons moving from the target to the right (Fig.1). It is clear that in the stationary state the flux of electrons moving to the right (from the target) must be compensated by the flux of electrons moving to the left and entering the target, i.e. the condition $\int_{-\infty}^{\infty} v_x\, f(\vec{v},o)\, dv_x = 0$ must hold. (Note that $f(\vec{v},0)$ in (9) is time independant and, hence, this condition is valid for any moment). This requirement will be satisfied if we continue the function $f(\vec{v},o)$ defined by (9) for $v_x>0$ to the negative electron velocities $v_x<0$, assuming, hence, that $f(\vec{v},o)$ is an even function of v_x. Thus the boundary condition at x=0 will be formulated as $v_x f(\vec{v},o)=q(\vec{v})$ where $q(\vec{v})$ is an odd function of v_x defined explicitely for $v_x>o$ by (10). Additionally we will assume that the functions E(x), f(x) diminish sufficiently fast with $x\to\infty$. The set of equations (5),(7) satisfying these boundary conditions can be solved with the help of the Laplace- transformation. Multiplying these equations by e^{-ikx}, integrating them in the interval $0\leqslant x<\infty$ and using the equality $v_x\int_0^\infty \frac{\partial f}{\partial x} e^{-ikx}dx = -v_x f(\vec{v},o)+ikv_x f_k = -q(\vec{v})+ikv_x f_k$ we arrive at the equations for the Laplace- transforms $f_k=\int_0^\infty fe^{-ikx}dx$, $E_k=\int_0^\infty Ee^{-ikx}dx=ik\,\varphi_k$

$$\frac{\partial f_k}{\partial t} + (\nu + ikv_x\, f_k) - \frac{e\, E_k}{m_e}\frac{\partial f_{eo}}{\partial v_x} = q \qquad (10)$$

$$-k^2\varphi_k = \chi_i^2\, \varphi_k - 4\pi e \int f_k\, d\vec{v} \qquad (11)$$

We now apply to (10), (11) a Laplace time transform

$$\varphi_{kp} = \int_0^\infty \varphi_k(t)e^{-pt}dt,\quad \varphi_{kp} = \int_0^\infty f_k(t)e^{-pt}dt \qquad (12)$$

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$$\varphi_{kp} = \frac{4\pi e}{k^2\, \varepsilon_\ell(p,k)} \int_{-\infty}^{\infty} \frac{\bar{q}(v_x)dv_x}{p + ikv_x+\nu} \qquad (13)$$

where $\bar{q}(v_x) = \int q(\vec{v})\, d\vec{v}_\perp$ and

$$\varepsilon_\ell(p,k) = 1 - \frac{4\pi e^2 i}{m_e k} \int_{-\infty}^{\infty} \frac{\partial f_{eo}/\partial v_x}{p + \nu + ikv_x}\, dv_x + \frac{\chi_i^2}{k^2} \qquad (14)$$

We can now represent the distribution of potential $\varphi(x,t)$ outside of the pellet, using (13), with the help of the inverse Laplace- transformation as follows

$$\varphi(x,t) = \frac{e}{\pi} \int_{-\infty+i\mu}^{\infty+i\mu} \frac{e^{ikx}}{k^2}\, dk \left(\int_{\sigma-i\infty}^{\sigma+i\infty} \frac{e^{pt}dp}{p\, \varepsilon_\ell(p,k)} \int_{-\infty}^{\infty} \frac{\bar{q}(v_x)dv_x}{p + \nu + i\ \ v_x} \right) \qquad (15)$$

where $\mu=\mathrm{Im}k<0$ and integration in the p- plane is carried aout along the straight line $\mathrm{Re}p=\sigma$ lying to the right of all the singularities of the function φ_{kp}. (The analysis leading to (15) is similart to that used in the theory of longitudinal plasma waves, see e.g. (Akhieser et al, /4 /, p.7).

To evaluate the integral over v_x in (15) we will use an asssumption that the distribution functions feo, fb are sharply peaked with respect to v_x. Conside-ring, hence, f_{eo}, f_b in (9) as delta-functions we get an expression

$$\bar{q} = \pm\, n_b\, v_b\, w_{si}\ (v_x,\ v_b) \pm n_p\, v_d\, w_{se}\ (v_x, v_d) \pm n_p\, v_d\, w_b\ (v_x, v_d) \qquad (16)$$

for $\bar{q}$, where $\pm$ correspondes resp. to $v_x>0$ and $v_x<0$. The distribution w_{si} (v_x) in (16) (and similarly the distributions w_{se}, w_b) possesses pronounced maxima for the secondary electron velocities corresponding approximately to energy about 10eV. This distribution can be hence approximated by the function (see Akhieser, /4 / , p.12)

$$\eta_{si} = \frac{\eta_{si}}{\pi} \frac{v_1}{(v_x-v_0)^2 + v_1^2} \qquad (17)$$

where v_0 (maximum) and v_1 (width) are some functions of v_b.

We can now represent the integral over v_x in (15) as follows

$$\int_{-\infty}^{\infty} \frac{\bar{q}\ (v_x)\ dv_x}{p + \nu + ikv_x} = -\eta_{si} \frac{1}{p + \nu + ikv_0 + kv_1} \qquad (18)$$

(Terms containing w_{se} and w_b give similar contributions).

This function has evidently a singularity for $p=-ikv_0-kv_1-\nu$. The integral over p in (15) possesses, hence, two poles: for $p = 0$ and for $p = -ikv_0-kv_1 - \nu$.

We can now evaluate the integral over p in (15) using these two poles by the residue method thus arriving at the expression

$$\int_{\sigma-i\infty}^{\sigma+i\infty} \frac{e^{pt}dt}{p\ \varepsilon_\ell(p,k)} \left(\int_{-\infty}^{\infty} \frac{\bar{q}(v_x)dv_x}{p + \nu + ikv_x} \right) = \frac{2\pi\ i\ \eta_{si}}{k(iv_0+v_1)+\nu} \left(\frac{1}{\varepsilon_\ell(o,k)} - \frac{e^{-(\nu+k(iv_0+v_1)t}}{\varepsilon_\ell(k(iv_0+v_1)+\nu,k)} \right) \quad (19)$$

To evaluate the remaining integral over k in (15) we will use the assumptions which are usually well fulfilld for ICF conditions, $v_d \ll u_e$ and $v_0 \ll u_e$.
It these case case we can replace the dielectric function $\varepsilon_\ell(p,k)$ by a simple expression

$$\varepsilon_\ell = 1 + \frac{\chi^2}{k^2}, \quad \chi^2 = \chi_e^2 + \chi_i^2 = \frac{8\pi\, n_p\, e^2}{T} \tag{20}$$

The integral over k prossesses, hence, a pole $k = -i\chi$ in this case. Using this pole we obtain finally by the residue method

$$\varphi = 8\pi e \frac{1}{\chi^2} e^{-\chi x} n_p \,\mathrm{Re}\, \eta_{eff} \left(1 - e^{-\chi v_0 t + i\chi v_1 t}\right) \tag{21}$$

where $\eta_{eff} = \dfrac{v_d}{(v_0+iv_1)_{si}} + \dfrac{v_d}{(v_0+iv_1)_{se}} + \dfrac{v_d}{(v_0+iv_1)_b}$

The expression (21) shows that shortly after the start of the ion pulse, for $t > \dfrac{1}{\chi v_0} = \dfrac{1}{\omega_p} \dfrac{u_e}{v_0}$, the potential distribution tends to stationary state described by the formula

$$\varphi = \frac{8\pi e}{\chi^2} e^{-\chi x} n_p \,\mathrm{Re}\, \eta_{eff} = \frac{T}{e} e^{-\chi x} \mathrm{Re}\, \eta_{eff} \tag{22}$$

Corresponding density distribution of the electron accumulated at the surface as a result of the secondary electron emission and of the back-scattering is given for $t \to \infty$ by

$$n(x) = \int f d\vec{v} = \frac{1}{2\pi i} \int_{-\infty+i\mu}^{\infty+i\mu} \frac{e^{ikx} dk}{\varepsilon_\ell(0,k)} \left(\int \frac{\bar{q}(v_x)\, dv_x}{ikv_x + v}\right) = n_p\, e^{-\chi x} \mathrm{Re}\, \eta_{eff} \tag{23}$$

The kinetic theory developed above proves that, at least for small η, the surface effects lead to important effects around of the pellet-to appearence of a double layer where the secondary electrons are concentrated and where strong electric fields can arise.
Expression for n(o) following from (23) coincide with that given by the elementary formula (1) if we intruduce η_{eff} instead of η and assume in (1) $\eta \ll 1$. For small coefficients $\eta_b, \eta_{se}, \eta_{si}$ the expression for the field $E(o) = -\partial\varphi/\partial x$ given by (22) also agrees with that following from the elementary considerations. For not small η_e, η_{se}, η_{si}, one can derive a more relible estimate for the field E(x) using (1) as a boundary value for n(x) at the pellet surface and employing the assumptions that both plasma ions and plasma electrons obey around of the pellet the Boltzmann distribution $n_e = n_p e^{e\varphi/T}$, $n_i = n_p e^{-e\varphi/T}$. For plasma ions repelled from the positively charged pellet this approximation must be good. According to (Mitchener, /5/ p.132) this must be a reasonable approximation also for plasma electrons attracted to a solid loaded body. Thus, we can formulate the expression for the potential φ_o arising at the pellet, against the potential far from it, using the rela-

tion $n_o=\alpha n_p=n_p e^{e\varphi_0/T}$, as follows

$$\varphi_o = \frac{T}{e}\ \ell n\ \alpha, \quad \alpha = \frac{\eta_b+\eta_{se}+\eta_{si}}{1-\eta_e-\eta_{se}-\eta_{si}} \qquad (24)$$

The electric field distribution $E=-\partial\varphi/\partial x$ around of the pellet can be now found from the Poisson equation

$$\mathrm{div}\ \vec{E} = \frac{d^2\varphi}{dx^2} = -4\pi e\ (n_i-n_e) = \frac{T}{e}\ \chi^2(e^{\frac{e\varphi}{T}} - e^{-\frac{e\varphi}{T}}) \qquad (25)$$

which provides an expression for the field E(o)

$$E(o) = \frac{2T}{e}\ \chi\ \sqrt{\cosh\frac{e\varphi_0}{T} - 1} = \frac{2T}{e}\ \chi\ \sqrt{\cosh\frac{e\varphi_0}{T} - 1} \qquad (26)$$

For sufficiently big ratio α the field at the pellet surface E(o) is given by the expression

$$E(o) = \frac{\sqrt{2}T}{e}\ \chi\ \sqrt{\alpha} \sim T^{1/2}\ n_p^{1/2} \qquad (27)$$

Note that the field E(o) rises with the reaction chamber plasma density as $n_p^{1/2}$ and, hence, that at low pressures the effect disappears.

A peculiar property of the expressions (1), (27) derived above is that they do not depend on the ion beam current density. This is the consequence of the assumption that the ion beam current is completely neutralized by the plasma return current and of the assumptions $u_i<v_d$, and $v_d<u_e$. For T=50ev, $n_p=10^{18}$ cm^{-3} (Argon plasma) these assumptions correspond to the restrictions on the ion beamcurrent density j_b: $0,1MA/cm^2<ej_b<50MA/cm^2$ Beyond this interval E(o) depends essentially on j_b. To evaluate the fields arising at the pellet surface according to (26), (27) it is necessary to know the coefficients η_b, η_{se}, η_{si} introduced above. The pellet surface must be in a state of a dense gaseous plasma reached shortly after the start of the ion pulse. Experimental date about these coefficients for such targets are missing. However, existing theories of the secondary electron emission and of the electron back scattering from the metal targets can be also applied for the case of dense plasmas, provided they have a sufficiently sharp boundary. Most important of these processes in our case is, probably, the secondary electron emission induced by the beam ions. According to theory of (Sternglass,/6/) corresponding coefficient η_{si} is given for ions of MeV- energy having in plasma an effective charge Z_{eff} by the formula $\eta_{si}=5\cdot 10^{-2}Z_{eff}^2 / (\frac{m_e}{m_b} W_b(MeV))^{1/2}$ where m_b is the mass of the beam ions and W_b is their energy (in MeV).
According to this formula $\eta_{si}\tilde{=}1$ for protons ($Z_{eff}=1$) of energy W_b=5MeV; $\eta_{si}=1,5\div6$ for He- ions of energy 8MeV (charges $Z_{eff}=1\div2$) and $\eta_{si}=1\div9$ for Li- ions of energy 30MeV ($Z_{eff}=1\div3$). The coefficients η_{se} of secondary electron emission and of electron back- scatering can be evaluated with the help of resp. theory of (Kadyshevich, /7/) and of (Everhart, /8/). For electrons of energies 50-100eV we

get in this way a crude estimate for Pb-targets $\eta_{se}=0,2\div0,4$, $\eta_b=0,1\div0,2$. Using these values of the coefficients η_b, η_{se}, η_{si} we can find with the help of (26) following estimates for the fields arsing at the pellet surface in case the condition $v_d=j_b/n_p<u_e$ is satisfied and for Argon plasma surrounding the pellet with T=50eV and $n_b=10^{18}cm^{-3}$. For protons of energy W_b=5MeV we have $E(o)=(2,8\div4,2)\cdot 10^7V/cm$; for He- ions of energy 8MeV we have $E(o)=(2,2\div6,4)10^7V/cm$; and for Li-ions of energy 30MeV we have $E(O)=(2,1\div7)10^7V/cm$. Thus, the fields corresponding to conditions of breakeven must be, indeed, strong.

It must be noted that electric fields arising at the surface of an ICF- pellet irradiated by Nd- laser light (power densities up to $10^{15}W/cm^2$) were detected earlier and investigated both theoretically and experimentally (see e.g. Elieser, /9/). These fields have another nature than those given by (25), (26) - they are basically due to electron emission from a hot pellet into surrounding nonhomogenious colder plasma. They are, also, at least an order of magnitude smaller than the fields due to pellet irradiation by a neutralized ion beam investigated above.

We must stress here that the analysis performed above is very approximate what concerns the case of large coefficients η_{si}, η_{se}, η_b.

The assumption about Bolzmannian distribution for electrons outside of the pellet used above represents, certainly, a crude approximation and a more detailled non-linear analysis, taking into account pecularities of surface effects, can reveal som new important facts. One can conjucture, in particular, that energy distribution of the secondary electrons ejected from the pellet can influence essentially the structure of the double layer surrounding the pellet. Especially important seems to be the fact that among secondary electrons ejected by the beam ions an important group of electrons having high energies, due to Auger- transitions in the target atoms, is present (see e.g. Schneider et al., /10/). As a consequence, the thickness of double layer L may exceed essentially the Debye radius r_d.

Another limitation of this model is the assumption of the presence of a sharp target boundary. In fact, however, the target surface must become diffuse shortly after start of the ion pulse. One can suppose that the ions of the target surface can be pushed outside of the target under the influence of the field arising at its surface thus reducing the target charge. (Such an effect was considered for the case of target sutuated in the vacuum and irradiated by a heavy ion beam by A Langdon, Report at HIIV 88, Darmstadt, 28-30.8.1988). It seems, however, that in the situation considered here, i.e. for the target embedded in a thick plasma in the reaction chamber, this affect can not be of importance. The ions of the target plasma can not move independently of the plasma electrons who tend to neutralize them. The process will have, probably, the character of the ambipolar diffusion resulting in a redistribution of the target particles in an immediate vicinity from its surface, but the loading provoked by the secondary electron emission will, certainly, persist. One can even conjecture that the diffusiveness of the target surface favouring escape of energetic secondary electrons will enhance the effects considered above.

3. Effects due to the electric fields arising outside of the pellet

Strong fields E arising at the pellet surface can provoke different instabilities most important of which, being, probably the capillary wave instability characteristic for a loaded, conducting liquid surface (see e.g. Landau, /11/). (Fig. 2). This instability can arise also at the surface of a loaded dense plasma as a result of action of electrostatic force equal (per cm^2 of the surface) $p_o=\sigma E/2$, $\sigma=E/4\pi$ being the surface charge density, i.e. $p_o=E^2/8\pi$.
This force becomes the stronger the bigger σ, i.e. the bigger the curvature of the surface is. Thus, the electrostatic force acting on the crest of the capillary wave tends to break it away. The grows rate γ of this instability in case the conducting liquid borders the vacuum is equal to (Landau,/11/ , §. 5) $\gamma=Ek/(4\pi\varrho)^{1/2}$, $k=2\pi/\lambda$, where ϱ is the mass density of the liquid and λ is the wavelength of the capillary wave. The ICF- problem discussed above is more complicated in a sense that the pellet surface borders not the vacuum but a layer containing noncompensated negative charge due to back-scattered and secondary eledtrons, and, besides, a neutralized ion beam impinges on it. Similar problem (excitation of the capillary wave instabilities in the small liquid metal droplets surrounded by a plasma) was investigated recently both experimentally and theoretically (Vladimirov, 1982) and it was found that the presence of a noncompensated electron layer at the liquid metal surface can strongly increase the growth rate of the capillary wave. We will repeat in a simplified form the analysis described in (Vladimirov/12/). The approach used below is valid if the thikness of the double layer L is small, $L \ll \lambda$.
Suppose, on the border of the liquid metal a surface perturbation ζ arisies (Fig. 2). It leads to appearance of an additional electrostic pressure $p-p_o=((E+E')^2-E^2)/8\pi \approx EE'/4\pi$ acting on the liquid surface. We can, hence, formulate, using the Bernoulli equation $\varrho \frac{\partial \Psi}{\partial t} + \frac{1}{2}\varrho v^2 = p - p_0$ and neglecting $\frac{1}{2}\varrho v^2$ ($\vec{v}=\nabla\psi$, ψ is the velocity potential obeying the equation $\Delta\psi=0$), a relation

$$\varrho \frac{\partial \psi}{\partial t} = \frac{E}{4\pi}\frac{\partial \varphi'}{\partial z} \qquad \text{for } z = 0, \quad E' = -\partial\varphi'/\partial z \tag{28}$$

Potential φ' due to the surface perturbation is governed by the Poisson equation

$$\Delta\varphi' = \left(\frac{\partial^2}{\partial x^2} + \frac{\partial^2}{\partial z^2}\right)\varphi' = 4\pi e\delta n = \frac{4\pi e^2}{T} n_o\varphi' \tag{29}$$

where δn is an electron density perurbation in double layer connected with n_o (given by (1)) by the Boltzmannian relation $(\delta n/n_o) = e\varphi'/T$. (We are using here the model of the double layer of thickness L formed by the positively charged pellet surface and negatively charged sheath characterised by the constant electron density n_o defined by (1)).
As the boundary condition for (29) we will use the relation

$$\varphi' = E\zeta \qquad \text{for } z = 0 \tag{30}$$

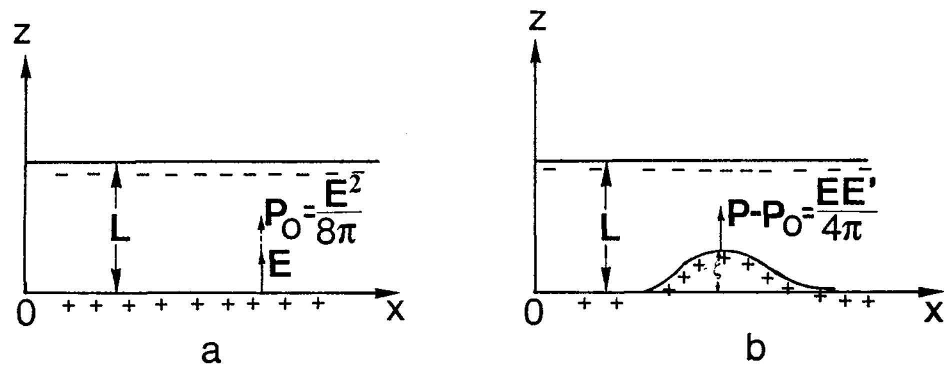

Fig. 2. Schematic representation of the capillary wave instability arising in a loaded liquid bordering the double-layer of thickness L. a. Pressure $p_0 = E^2/8\pi$ due to the presence of noncompensated positive charge in the liquid acts on its surface. b. The appearance of a surface disturbance $\zeta < L$ leads to an additional pressure $E\,E'/4\pi$ acting on its crest. This additional pressure is bigger in case the charged liquid borders not the vacuum but the double-layer containing the free electrons since the electrons tend to gather nearby the crest of the perturbation thus increasing the additional electric field and additional pressure acting there.

which is valid provided $\zeta \ll L$. Solution of (29) valid for $0 \leq z \leq L$ is

$$\varphi' = E\zeta \exp\left(ikx-\sqrt{k^2+\kappa^2}\,z\right), \quad \kappa^2=4\pi n_0 e^2/T \tag{31}$$

i.e. $E = \partial\varphi/\partial z = E\zeta\sqrt{k^2+\kappa^2}\exp\left(ikx-\sqrt{k^2+\kappa^2}\,z\right)$.

Note that the electric field at the border of the pellet $|E|=E\zeta\sqrt{k^2+\kappa^2}$ proves to be bigger by the factor (κ/k) than the field $|E|=E\zeta k$ corresponding to the case $\kappa=0$ when the pellet borders the vacuum.

Using the solution (31), solution $\psi=A\exp(i(kx-\omega t)-kz))$ for the velocity potential ψ (A is a constant) and inserting these definitions into (28) and into the definition of the normal velocity $v_z = \frac{\partial\zeta}{\partial t} = \frac{\partial\psi}{\partial z}$ we get the dispersion relation $\omega^2 = -\frac{E^2k}{4\pi\varrho}\sqrt{k^2+\kappa^2}$ allowing to find the expression

$$\gamma = \frac{Ek}{\sqrt{4\pi\varrho}}\left(1+\frac{\kappa^2}{k^2}\right)^{1/4}, \quad k = 2\pi/\lambda \tag{32}$$

for the growth rate. As follows from (32) the growth rate γ exceeds that for the vacuum case by the factor $(\kappa/k)^{1/2}$ for $\kappa \gg k$. Assuming in (32) $E=10^8$ V/cm, $\rho=11.3$ g/cm^3 (Pb-shell), $\lambda=0,1$mm $n_0=10^{19}$ cm^{-3} (for $n_p=10^{18}$cm^{-3}, $\alpha=10$), we get $\gamma\approx10^9$ sec^{-1}. This means that the tamper shell can be disrupted within a few nanoseconds.

In this treatment we have neglected the presence of plasma electrons drifting to the pellet. Taking into account this fact requires consideration of coupled equations for the plasma electron hydrodynamical velocity, plasma electron density and potential φ' as it was done in (Vladimirov /12/).

This modification appears to be necessary if $u_e \lesssim v_d$. Calculations performed in (Vladimirov, /12/) lead to conclusion that the new Pierce-type instabilities corresponding to the coupled excitation of the capillary waves in the pellet surface and waves in the surrounding plasma can appear. Corresponding growth - rates can exceed that given by (32). The problem of appearance of such Pierce - type instabilities for the ICF case requires further consideration.

4. Conclusions

We have performed investigation of the loading arising in a ICF- pellet irradiated by a neutralized beam of light ions with parameters corresponding to breakeven. It was found for a generally adopted scheme of a freely falling pellet situated in a reaction chamber filled with a dense gas (i.e. with a dense plasma produced by the ion beams) that strong noncompensasted charges can arise herewith in the pellet as a result of emission of secondary electrons and of backscattering of the plasma electrons from the pellet surface. The loading of the pellet can give rise to the fields at its surface of the order of 10^8V/cm in case of a light ion beam moving in a plasma of the reaction chamber with parameters: electron density $n_p \approx 10^{18} cm^{-3}$, temperature $T \approx 50-100$ eV. These fields can provoke capillary wave instabilities of the pellet surface leading to disruption of the tamper shell within a few nanoseconds. Thus, it seems that phenomena considered here can lead to serious complications in case light ions will be used as drivers since high currents of the order of 10MA are required in this case and, accordingly, the plasma surrounding pellet in the reaction chamber must be dense. For neutralized heavy ion beams these effects are not dangerous according to the formula (27) since much lower ion currents are required in this case. It must be noted here , however, that using of non- neutralized heavy ion beams moving to the pellet in a vaccum chamber can result in strong loading of the pellet provoking dangerous capillary wave instability on its surface with the growth rate given by the formula $\gamma = Ek/(4\pi\rho)^{1/2}$

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