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EFFECTS OF ELECTRON MASS VARIATIONS IN A STRONG ELECTROMAGNETIC WAVE

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At present, one of the actual problems of plasma physics is the investigation of the interaction of strong electromagnetic radiation with a plasma in which the coherent wave energy can be transformed into: a) particle random energy and then to be used for heating of plasma [1,2]; b) the generation of a constant current along the propagation vector [3], the phenomenon of which can be interpreted as a "light wind" and is somewhat similar to the acoustic wind; c) the beat-wave acceleration; d) the wake-field accelerations; e) and causes self-focusing and filamentational instability, etc.

The point is that the progress in laser technique and HF generation achieved by this time allows to obtain the electromagnetic fields in which the electrons can acquire relativistic velocities. In space the electromagnetic radiation of cosmic objects (nuclei, galaxies, radiogalaxies, pulsars, etc.) may serve as a source of such strong fields.

Therefore, it is clear that the constant account of this relativistic effect is necessary in order to describe the physical pattern of nonlinear processes occuring at such an interaction.

At present, we have some review publication about relativistic effects [4-6].

Some problems will be discussed in the present talk:

1. Excitation of plasma waves by the beating of electromagnetic radiation [7].

2. Relativistic wake-field generation by an intense

laser pulse in a plasma [8].

3. Compression of plasma at self-focusing of powerful electromagnetic pulses [9].

All these problems were observed ourselves and that is why I am going to give a talk more detaily.

I. At present, one of the promising methods for acceleration of charged particles is the using of the longitudinal plasma waves. At present the different possibilities for the plasma waves excitation are discussed[10]. The idea of the longitudinal waves generation by the laser radiation is suggested in the paper [11]. Here the problem of excitation of plasma wave by beating of two electromagnetic waves is discussed. The frequencies and wavenumbers satisfy the resonance conditions

$$\Delta\omega = \omega_1 - \omega_2 = \omega_p \;, \qquad \omega_1, \omega_2 \gg \omega_p$$

$$\Delta K = K_1 - K_2 = K_p \;, \qquad K_1, K_2 \gg K_p \tag{I.1}$$

where ω_p and K_p are the frequency and the wavenumber of excited electron plasma wave.

The purpose of this paper is to study the excitation of longitudinal wave by beating of two electromagnetic waves taking into account their interaction. It is shown that the amplitude of the longitudinal wave increases in time not linearly but exponentially. The maximum growth is found. The requirement for shortness of injected particles acceleration time against the plasma wave growth time creates certain limitation for the pump waves amplitudes.

The strength of the transversal electric field is represented as a sum of two waves with frequencies ω_1, ω_2 under slowly changing complex amplitudes

$$E=\frac{1}{2}\left\{E_1(z,t)\,e^{-i\omega_1 t+ik_1 z}+E_2(z,t)\,e^{-i\omega_2 t+ik_2 z}+c.c.\right\} \qquad (I.2)$$

Since ω_1, ω_2 are much more larger ω_p the influence of the external fields on electron fluid can be described by averaging HF pressure force. Under the weak relativity for the small perturbations of the electron concentration we obtain

$$\left(\frac{\partial^2}{\partial t^2}+\omega_p^2\right)\frac{\delta n_e}{n_0}=\frac{e^2}{4m_e^2}\frac{\partial^2}{\partial z^2}\left\{\frac{|E_1|^2}{\omega_1^2}+\frac{|E_2|^2}{\omega_2^2}+\frac{1}{\omega_1\omega_2}\left(E_1E_2^*e^{-i\Delta\omega t+i\Delta k z}+c.c.\right)\right\} \qquad (I.3)$$

where n_0 is the equilibrium concentration of electrons.

Below we consider only the case, when the characteristic time of the complex amplitudes E_1, E_2 variation is much more larger than the period of plasma oscillations. The solution of equation (I.3) can be represent as a sum of rapidly and slowly changed parts

$$\delta n_e=\overline{\delta n}+\frac{1}{2}\left\{N(z,t)\,e^{-i\Delta\omega t+i\Delta k z}+c.c.\right\} \qquad (I.4)$$

From (I.3) we obtain $\overline{\delta n}\ll N$. Thus under description of the field complex amplitudes dynamics we shall consider only HF part of perturbation (I.4). From (I.4) as well as from Maxwell equations and the equation of electrons fast motion for envelop amplitudes E_1, E_2 and for the value N we obtain

$$i\left(\frac{\partial}{\partial t}+v_{g1}\frac{\partial}{\partial z}\right)E_1=\frac{\omega_p^2}{4\omega_2}\frac{N}{n_0}E_2 \qquad (I.5)$$

$$i\left(\frac{\partial}{\partial t}+v_{g2}\frac{\partial}{\partial z}\right)E_2=\frac{\omega_p^2}{4\omega_1}\frac{N^*}{n_0}E_1 \qquad (I.6)$$

$$i\frac{\partial}{\partial t}\frac{N}{n_0}=\frac{e^2\kappa_p^2}{4m_e^2\omega_1\omega_2\omega_p}E_1E_2^* \qquad (I.7)$$

where

$$V_{g1}=\frac{\kappa_1 c^2}{\omega_1}\ , \qquad V_{g2}=\frac{\kappa_2 c^2}{\omega_2} \qquad (I.8)$$

Below we assume $V_{g1}\simeq V_{g2}\simeq V_g$.

Let us consider the stability of fields with constant amplitudes and uniform distribution. The terms with the spatial derivatives in (I.5),(I.6) are assumed as small ones. According the unknown quantities are represented as

$$E_1=\left\{E_{10}+\delta E_1' e^{-i\Omega t+iqz}+\delta E_1'' e^{i\Omega t-iqz}\right\}e^{-i\frac{\Omega_\kappa}{2}t-iAt} \qquad (I.9)$$

$$E_2=\left\{E_{20}+\delta E_2' e^{-i\Omega t+iqz}+\delta E_2'' e^{i\Omega t-iqz}\right\}e^{i\frac{\Omega_\kappa}{2}t-iAt} \qquad (I.10)$$

$$N=\left\{N_0+\delta n_1 e^{-i\Omega t+iqz}+\delta n_2 e^{i\Omega t-iqz}\right\}e^{-i\Omega_\kappa t} \qquad (I.11)$$

$$N_0=\frac{e^2\kappa_p^2 n_0}{4m_e^2\omega_1\omega_2\omega_p}\,\frac{E_{10}E_{20}^*}{\Omega_\kappa} \qquad (I.12)$$

where amplitudes of the supporting waves E_{10} , E_{20} are constant. The last two terms in brackets (I.9),(I.10),(I.11) correspond to the perturbation caused by small spatial nonuniformity. The values Ω_κ , A are equal to

$$\Omega_\kappa=\sqrt{\frac{e^2\kappa_p^2\omega_p}{16m_e^2\omega_1\omega_2}\left(\frac{|E_{20}|^2}{\omega_2}-\frac{|E_{10}|^2}{\omega_1}\right)} \qquad (I.13)$$

$$A=\sqrt{\frac{\Omega_\kappa^2}{4}+a^2}\ , \qquad a^2=\frac{\omega_p^2\,|N_0|^2}{16\omega_1\omega_2 n_0^2} \qquad (I.14)$$

They define the frequency shifts of the supporting waves caused by their nonlinear interaction. After the linearization of equations (I.5),(I.6),(I.7) we obtain the dispersion relation:

$$\Omega^2\left\{(\Omega-qV_g)^2-4a^2\right\}-\Omega_K^2\,(2\Omega-qV_g)^2=0 \tag{I.15}$$

The amplitudes of supporting waves are assumed as close ones $|E_{10}|\simeq|E_{20}|$. From (I.13),(I.14) we obtain

$$\beta=\frac{\Omega_K^2}{a^2}\simeq\frac{\omega_p^2}{\omega_1\omega_2}\ll 1 \tag{I.16}$$

The numerical calculation of the dispersion relation (I.15) shows, that the modulation frequency Ω has the imaginary part, which corresponds to the perturbation growth, only in the case when ratio a/qV_g is larger than the certain value. This critical value depends on the parameter β . When $\beta=0.01$ the growth rate acquires its maximum

$$Im\,\Omega_{max}=1,63\cdot\Omega_K=0,41\cdot\frac{V_{01}}{c}\frac{\omega_p}{\omega_2}\omega_p$$

$$V_{01}=\frac{e|E_{10}|}{m_e\omega_1 c} \tag{I.17}$$

which is reached for perturbation wavelength

$$\lambda\simeq 6,296\cdot\lambda_p \qquad (\lambda_p=2\pi/K_p)$$

The characteristic time of the particles acceleration by plasma wave is equal [11]

$$t_0=\frac{2}{\omega_p}\frac{\omega_1^2}{\omega_p^2} \tag{I.18}$$

It is clear that for the efficiency of acceleration the acceleration time must be less or have an order of the plasma wave growth characteristic time $t_c=1/Im\,\Omega_m$

According to (I.16),(I.18) this condition leads to the inequality

$$\frac{V_{01}}{c} \leqslant \frac{1}{0{,}82} \frac{\omega_p}{\omega_1} \qquad (I.19)$$

In the case of the opposite inequality the regular character of plasma wave is destroyed before the particles acquire the maximum energy under acceleration.

In paper [12] the characteristic time for the plasma wave amplitude variation due to the pump depletion is estimated

$$t_p = 2\sqrt[3]{12}\, \frac{\omega_1^2}{\omega_p^3} \left(\frac{c}{V_{01}}\right)^{2/3} \qquad (I.20)$$

Comparing the relation (I.20) with the instability characteristic time t_c , we show, that in the condition

$$\left(\frac{V_{01}}{c}\right)^{1/3} \gg 0{,}54\, \frac{\omega_p}{\omega_1} \qquad (I.21)$$

which is easy to satisfy, the depletion time is larger than the characteristic time t_c . It means that condition for the particles effective acceleration is determined not by the depletion process but by the instability considered above.

II. The excitation of plasma waves by the low-power laser pulses, which generate the linear plasma response in the form of plasma wake-field (PWF) was considered in the papers [11,13,14] . The laser wake-field acceleration (LWFA) scheme based on this mechanism will apparently exhibit a number of principal advantages in comparison with the plasma beat-wave (PBWA) and the plasma wake-field acceleration (PWFA) scheme which are intensively worked out [16,17] . However, it's clear, that the real success of the LWFA scheme can be associated only with the relativistically intense laser pulses [15]. As it is shown in this paper, they are able to generate fully nonlinear longitudinal waves with relativistic electric fields in a plasma $E > mc\omega_p/e$. Only such fields can assure reaching the accelerating gradients a few tens of GeV/m and more.

Starting from the Maxwell's equations and the equations of motion of cold and collisionless electrons we derive the following equation for the plasma electron momentum:

$$\frac{\partial^2 \vec{P}}{\partial t^2} - \Delta\vec{P} + \nabla(\nabla\cdot\vec{P}) + \frac{\partial}{\partial t}\nabla\sqrt{1+P^2} + \frac{\vec{P}}{\sqrt{1+P^2}}\left(1 + \frac{\partial}{\partial t}(\nabla\cdot\vec{P}) + \Delta\sqrt{1+P^2}\right) = 0 \qquad \text{(II.1)}$$

where the dimensionless quantities are introduced:

$\vec{P} \to \vec{P}/mc$; $t \to \omega_p t$; $\vec{r} \to \kappa_p \vec{r}$; $\kappa_p \equiv \omega_p/c$.

We consider the propagation of a circularly polarized laser radiation pulse in the z direction with unchanged shape, thus:

$$\vec{P}_\perp = \frac{P_\perp(z - V_g t)}{\sqrt{2}}(\hat{x} + i\hat{y})\exp(-i\omega_0 t + i\kappa_0 z) + c.c. \qquad \text{(II.2)}$$

where ω_0 and κ_0 are the frequency and the wave number of the electromagnetic field.

For the purposes of accelerating we are interested in longitudinal plasma waves with $V_{ph} \simeq 1$, so we consider the transparant plasma, where $V_g \simeq 1$ or, that is the same in dimensional units, $\omega_0^2 \gg \omega_p^2/(1 + P_\perp^2/m^2c^2)^{1/2}$. From Eq.(II.1) we obtain the equation describing the nonlinear dynamics of longitudinal plasma motion:

$$\frac{d^2 y}{dx^2} = \frac{1}{2}\left(\frac{\gamma_\perp^2}{y^2} - 1\right) \qquad \text{(II.3)}$$

where $\gamma_\perp^2 = 1 + P_\perp^2$, $x = \kappa_p(z - ct)$, $y = \sqrt{\gamma_\perp^2 + P_\parallel^2} - P_\parallel = 1 + \Phi$,

$\bar{\Phi} = e\phi/(mc^2)$ is the dimensionless potential of the longitudinal field.

In the case of a dotting laser pulse with the length of ℓ $p_\perp^2(x) = p_\perp^2[h(x+\ell) - h(x)]$, where $h(x)$ is the Heaviside unit function, we succed in obtaining of an analytic solution of Eq.(II.3). For the natural boundary conditions: $y(0)=1$, $y'(0)=0$, inside the pulse $(-\ell \leq x < 0)$ where $1 < y \leq \gamma_\perp^2$, we have

$$x = -2\gamma_\perp E(\varphi_i, \kappa_i) + 2\left[\frac{(\gamma_\perp^2 - y)(y-1)}{y}\right]^{1/2}$$

$$\varphi_i = \arcsin\left[\frac{\gamma_\perp^2(y-1)}{(\gamma_\perp^2 - 1)\, y}\right]^{1/2}, \quad \kappa_i = \left[\frac{\gamma_\perp^2 - 1}{\gamma_\perp^2}\right]^{1/2} \tag{II.4}$$

here $E(\varphi_i, \kappa_i)$ is the incomplete elliptic integral of the second kind. When the pulse length is $\ell = 2\gamma_\perp E(\kappa_i)$ $(E(\kappa_i) \equiv E(\frac{\pi}{2}, \kappa_i))$ the longitudinal excitation reaches its maximum value and in the area behind the pulse $(x < -\ell)$ where $1/\gamma_\perp^2 \leq y < \gamma_\perp^2$, we have:

$$x = -\ell - 2\gamma_\perp E(\varphi_e, \kappa_e)$$

$$\varphi_e = \arcsin\left[\frac{\gamma_\perp^2(\gamma_\perp^2 - y)}{\gamma_\perp^4 - 1}\right]^{1/2}, \quad \kappa_e = \left[\frac{\gamma_\perp^4 - 1}{\gamma_\perp^4}\right]^{1/2} \tag{II.5}$$

The maximal values of the electric field, of the longitudinal momentum of electrons and the potential in the excited PWF are:

$$\bar{E}_{\parallel max} = \frac{\gamma_\perp^2 - 1}{\gamma_\perp}, \quad P_{\parallel max} = \frac{\gamma_\perp^4 - 1}{2\gamma_\perp^2}, \quad \bar{\Phi}_{max} = \gamma_\perp^2 - 1 \tag{II.6}$$

The space period of the driven waves (in the dimensional form) is $\lambda = 4(c/\omega_p)\gamma_\perp E(\kappa_e)$. In the nonrelativistic case $\gamma_\perp \simeq 1$ we have $\lambda = 2\pi c/\omega_p$, and in the ultrarelativistic case $(\gamma_\perp \gg 1)$ $\lambda \simeq 4(c/\omega_p)\gamma_\perp$

Thus, the intense dotting laser pulse generates a longitudinal waves with the relativistic amplitude and the phase velocity close to C.

The numerical analysis of Eq.(II.3) makes possible to consider the pulses with realistic shapes. In Fig. (a-c) is presented the numerical solution in the case of the pulses with Gaussian rise and Gaussian fall: $P_\perp^2(x) = \{8\exp(-x^2/(2\sigma_r^2))$, when $x \geq 0$, and $8\exp(-x^2/(2\sigma_f^2))$, when $x < 0$. These figures show that the front or the back cutoff of the driving pulse has a substantial effect on the amplitude of PWF and, although the front cutoff is more effective than the back cutoff, in this case the excited longitudinal field inside the pulse is considerably weaker than the laser field, and this favours the pulse propagation without any distortion.

Note, that the cutoff of an electron-beam bunch driving the PWF in the PWFA scheme, is also effective, but its difficulty remains as a large technological barrier in realization of PWFA [16,17], while in the case of laser pulses this is an evidently solvable problem.

This way excited PWF can accelerate the injected in a plasma electrons. Maximally possible increase of electron energy in the wake-field of the dotting laser pulse is $\Delta\mathcal{E} = 2\gamma_g^2 mc^2 \Delta\bar{\Phi}_m$, $\gamma_g = (1 - V_g^2)^{-1/2}$. In the ultrarelativistic case $\Delta\bar{\Phi}_m \simeq P_\perp^2$, and the acceleration length is $\ell_A = 2\gamma_g^2 (c/\omega_p) P_\perp$. Thus, in comparison with the nonrelativistic case the acceleration length is increased $P_\perp$-times, in turn the energy increment over this length increases $P_\perp^2$-times, and as a result the accelerating gradient is increased $P_\perp$-times.

In the considered model we have neglected the influence of the excited longitudinal waves on the pulse. The most fast instability distorting the laser pulse develops during the time $T \geq (\omega_0/\omega_p)^2 P_\perp^{3/2}/\omega_p$. However during this period of time the laser pulse covers much more greater distance than the PWF period is $\lambda \sim (c/\omega_p) P_\perp$.

The preliminary estimations of the nonlinear effects associated with the multidimensional geometry of real pulses give us an assurance in the success of the LWFA scheme.

III. Recently the study of nonlinear propagation of powerful electromagnetic waves in a plasma became very urgent. First of all it is connected with new methods of particle acceleration [18], and also with possibility of heating and compression of plasma in the laser-driven fusion devices. With this purpose the powerful laser short pulses will be used (the density of the energy flow exceeds 10^{17} V/cm^2, the duration is of the order of several femtoseconds [19]).

At the interaction of the intense electromagnetic waves with medium, the latter becomes inhomogeneous optically. The polarization vector, dielectric constant, the index of refraction and other values, characterizing physical properties of the medium, become dependent of the amplitude of the incedent wave. Powerful radiation changes the medium physical properties, that in its turn affects the wave propagation, i.e. there is self-effect of the wave. Unlike other nonlinear effects such as harmonics generation, the stimulated Raman and Mandelstamm-Brillouin scattering and parametric processes, where interactions take place at slightly differing frequency does not change in the process of self-effect, and the effect is observed in the change of its amplitude, polarization, shape of the angular and frequency spectrum.

Out of these phenomena, however, the effect of nonlinear refraction of the light beams, can be especially separated by its importance and influence upon other processes. If nonlinear mismach of the dielectric function grows with the increase of the intensity, in the range, where the intensity is maximum, the medium becomes at the same time the most optically dense. The dispersive relation and phase velocity of the wave vary. The phase velocity is maximum in the points where the wave intensity is minimum (and vice versa). The form of the wave-front is distorted. The beams diverge towards the area where the field is maximum which leads to the further increase of intensity in this area. The effect is snow-slip in character.

In 1962 Askarian [20] supposed that nonlinear refraction can impede beam difraction of the wave beams. Due to the strong inhomogeneity of intensity distribution (precisely the powerful pulse lasers and SHF generators have this kind of distribution) in the transverse cross section of the powerful electromagnetic waves the nonlinear refraction can be so significant that it lead to the beam collapse - its transverse dimansions decrease and almost all energy concentrates in a certain area. This phenomenon is called wave self-focusing. Soon the possibility of self-focusing was collaborated in several works [21-24] showing that the competition of diffraction and nonlinear refraction cause existence of three different regimes of beam propagation in the nonlinear medium: i) at low intensities diffraction predominates over nonlinear refraction and the beam diverges, ii) when the beam power exceeds a certain critical power [25], the beam focuses. In this case medium plays the role of the focusing lens - the transverse dimensions rapidly decrease and the beam is compressed to a point, while intensity increases to infinity. This point is called the focus, and distance on which the beam focuses - the self-focusing length. At the intensity increase other nonlinear effects become significant (generation of electrostatic pulses on the formation of which the part of the pumping wave energy is spend.

The so-called saturation of medium nonlinearity takes place which sets a lower bound on focal dimensions [25].
iii) Under certain conditions the self-trapping of the wave takes place. In this case the effective transverse width of the beam is either almost unchanged or oscillates in some interval. The self-supporting channel is formed in which the central rays are trapped with the intensity slightly exceeding the critical intensity. The outlying rays from the diffraction rings with radii increasing with the wave propagation. They carry away the untrapped energy of the pumping wave. The beam abberation can strongly influence this process and finally lead to the break of the channel [26].

The effect of the nonlinear refraction is the basis of the phenomenon of filamentation. The disturbances on the homogeneous form of transverse dictribution of intensity can grow and lead to the break of the beam into narrow and long structure of the so-called "filament". Depending on the width and form of the beam and on disturbances the formation of filaments can happen either before self-focusing of the beam as a whole or after it. In the case of non-stationary pumping the length of the self-focusing varies with the time. The moving focuses appear, the velocity of their motion depends on the duration and form of the pulse. At self-focusing the effective duration of the pulse also varies.

First publications [20,24] had been devoted to possibilities of self-focusing of electromagnetic waves in plasma, as in the medium, where the nonlinerity is achieved at comparatively low power of radiation and when the problem of LTR arose, the detailed investigation of self-focusing and filamentation in plasma became expedient. In the irradiation experiments of special targets by powerful lasers, transverse inhomogeneities, filaments, focal spots [29] are observed in the plasma corona, which are the causes of the implosion inhomogeneity, cause the enlightening of plasma in certain places, affect the damping and scattering of pumping wave, generation of strong magnetic fields. Self-focusing of the intense laser beam can be the cause of the plasma particle acceleration, due to which one can obtain the electron or ion beam with the energy of some MeV and higher [30], that is rather important for the development of the nuclear and elementary particle physics.

There are three basic mechanisms of the nonlinear refraction of the electromagnetic radiation in plasmas: thermal, strictional and relativistic. The thermal mechanism is caused by plasma particle thermal pressure and is connected with Ohm's heating and ionization.[31] The strictional mechanism is caused by redistribution of the plasma density under the effect of a ponderomotive force [32]. First, electrons are expelled from the localization region, the plasma quasi-neutrality is disturbed in the scale of the order of Debye radius r_D, the electrostatic field appears which causes displacement of the ions as well. The channels with considerably reduced plasma concentration are formed. The excess pressure of radiation inside the channel is compensated by the thermal plasma pressure outside.

With the decrease of the plasma concentration the refraction coefficient of the medium increases locally, that results in the self-focusing and filamentation of the electromagnetic waves. The characteristic time of the development of this process is $t_S \sim a/V_S$, where a is the width of the beam, V_S is ion sound velocity ($t_s \gg \omega_{pi}^{-1}$ where ω_{pi} - plasma frequency of ions). With the increase of the pumping wave intensity the relativistic effect connected with the increase of the electron mass in HF field becomes the basic mechanism of the nonlinear refraction. This nonlinearity is practically inertialees.

The characteristic time for which the electrons achieve relativistic velocity is $t_r \sim \omega_o^{-1}$, where ω_o - the frequency of the pumping wave. First the existence of such mechanism of self-focusing had been mentioned in 1969 [33]. This paper considered the weakly relativistic case without account for the strictional nonlinearity. Other papers [34-36] have not also considered the strictional nonlinearity, discussing the arbitrary relativistic case and assuming that the striction will not have time to develop since $t_r \ll t_s$. One should mention, however, that the characteristic time of self-focusing is not t_r , but $t_f = z_f/v_g \gg t_r$ (z_f - the self-focusing length, v_g - the group velocity). For the time t_f under certain conditions the electron striction can become essential [37-41]. For sufficiently intensive waves the scale of the distortion of quasi-neutrality of plasma can be much higher than Debye length and it could turn out that $t_f \ll t_s$). It means that for the waves with $I = e^2|E|^2/(m_e\omega_o c)^2 \gtrsim 10^{-2}$ (where e and m are the charge and the mass of the electron, respectively), the model given in [38] is incorrect. Even when $t_f \sim t_s$, in such high fields the ion velocity is several times higher than the ion-sound one and instead of the radiation pressure compensation by plasma pressure, the flows and shock waves can appear [42] . If $t_f < t_0 \ll \omega_{pi}^{-1}$ (where t_o - duration of the pulse) the relativistic self-focusing accounting for electron striction is reduced to the "stationary case" [37-41]. For sufficiently narrow beams ($k_o a \sim 10$, k_o - the wave vector of pumping wave) the ijection of plasma from the ranges of the focuses becomes considerable and may lead to the saturation of the self-focusing [38,39] . For broad beams and for very short pulses the condition $t_f, t_o \ll \omega_{pi}^{-1}$ is difficult to observe and one should solve nonstationary problem with account of the ion motion.

It is noteworthy that relativistic self-focusing

In the nonrelativistic case, when the change of the electron mass is completely neglected ($\delta m_e / m_e = 0$), at strictional mechanism condition $\delta n_e / n_o < 0$ is necessary for self-focusing. As it follows from (III.3), in the nonmagnetized collisionless plasma the self-focusing and filamentation can be accompanied by plasma compression ($\delta n_e > 0$). It is possible only for relativistic powerful electromagnetic waves, when $\delta m_e / m_e > \delta n_e / n_c$, i.e. when heating electrons predominate against density perturbation. At stationary propagation of the radiation there is no such regime. It turned out that at nonstationary self-focusing of considerably wide and short (in time) pulses in the certain conditions of the running focus, plasma can become denser.

The phenomena found can turn to be advantageous in laser wake-field accelerator. In order to develop well the process of the acceleration it is needed to have self-focusing of these transverse waves, since without it the length of the diffraction broadening of the beams is much less than the characteristic scales of the particle acceleration [30].

Let us consider the propagation of the circularly polarized HF electromagnetic waves in a collisionless plasma. From Maxwell equations for the relativistic intensive wave we can write

$$\Delta \vec{E} - \operatorname{grad} \operatorname{div} \vec{E} - \frac{1}{c^2} \frac{\partial^2 \vec{E}}{\partial t^2} = \frac{\omega_{pe}^2}{c^2} \frac{1}{\sqrt{1+I}} \left(1 + \frac{\delta n_e}{n_o}\right) \vec{E} \tag{III.5}$$

At intensities $I \gtrsim 10^{-2}$ the condition $\frac{T_e}{m_e c^2} \frac{\delta n_e}{n_o} \ll \ll I$ is fulfilled easily for the electron temperatures $T_e \lesssim 10^4$ eV, and therefore in such cases one can neglect the thermal pressure of plasma as compared with the pressure of the incident wave. Then we assume that quasineutrality of plasma if distorted and the electric field of the spatial charge is generated due to the effect and filamentation of the intensive electromagnetic beams with considerable decrease of the plasma electron density will be applied in experiments on lucidation of plasma and acceleration of the charged particles [30,35,40,42]. We think that in order to heat plasma with HF radiation it is more convenient to have such a regime of self-focusing when in the energy localization regions the plasma concentration increases or remains unchanged. In these cases the effectiveness of the pumping wave energy transmission in plasma can enhance considerably that is possible only at relativistic intensities. The considered process can be explained as follows. For circularly polarized wave the refraction index is

$$N^2 = \mathcal{E} = 1 - \frac{\omega_{pe}^2}{\omega_0^2} \frac{1 + \frac{\delta n_e}{n_0}(\vec{r},t)}{\sqrt{1 + I(\vec{r},t)}} = N_0^2 + \delta N \qquad \text{(III.1)}$$

where

$$\delta N = \frac{\omega_{pe}^2}{\omega_0^2}\left(1 - \frac{1 + \frac{\delta n}{n_0}(\vec{r},t)}{\sqrt{1 + I(\vec{r},t)}}\right) \qquad \text{(III.2)}$$

$\omega_{pe}^2 = 4\pi e^2 n_0 / m_e$, $n(\vec{r},t) = n_0 + \delta n(\vec{r},t)$ - electron concentration, n_0 - concentration of unperturbed plasma, $N_0 = \sqrt{1 - \omega_{pe}^2/\omega_0^2}$ - linear part of the refraction index, δN - nonlinear addition. For the weakly relativistic case ($I \ll 1$):

$$\delta N = \frac{\omega_{pe}^2}{\omega_0^2}\left(\frac{\delta m_e}{m_e} - \frac{\delta n_e}{n_0}\right) \qquad \text{(III.3)}$$

where $\delta m_e / m_0 = \frac{1}{2} I$. For nonlinear refraction $\delta N > 0$ is necessary, i.e.

$$\frac{1 + \frac{\delta n_e}{n_0}}{\sqrt{1 + I}} < 1 \quad \text{or} \quad \frac{\delta m_e}{m_e} > \frac{\delta n_e}{n_0} \qquad \text{(III.4)}$$

of the ponderomotive force. If the characteristic times (the pulse duration, the time of the nonlinear refraction) are much larger than ω_{pe}^{-1} $\left(\frac{\partial^2}{\partial t^2} \ll \frac{\omega_{pe}^2}{\sqrt{1+I}}\right)$ that is usually easily fulfilled and is equal to the neglection of the electron inertia, we obtain equation for the electron and ion densities:

$$\frac{\partial^2}{\partial t^2}\frac{\delta n_e}{n_0} = \frac{c^2}{\omega_{pe}^2}\left(\frac{\partial^2}{\partial t^2} + \omega_{pi}^2\right)\Delta\sqrt{1+I} \tag{III.6}$$

$$\frac{\partial^2}{\partial t^2}\frac{\delta n_i}{n_0} = \frac{m_e}{m_i}c^2\Delta\sqrt{1+I} \tag{III.7}$$

We study the propagation of the circularly polarized electromagnetic waves along axis :

$$\vec{E} = (\hat{x} + i\hat{y})\,E\,(z, r_\perp, t)\,e^{-i\omega_0 t + i k_0 z} \tag{III.8}$$

We replace the variables

$$\xi = z$$
$$\tau = t - \frac{z}{V_g} \tag{III.9}$$

superposing conditions

$$\frac{\partial^2}{\partial \xi^2} \ll \Delta_\perp\,,\quad \frac{1}{V_g^2}\frac{\partial^2}{\partial \tau^2} \ll k_0^2\,,\quad \frac{\omega_{pe}^2}{k_0^2 c^2}\frac{\partial^2}{\partial \tau^2} \ll k_0\frac{\partial}{\partial \xi} \tag{III.10}$$

For a underdense plasma $\omega_{pe}^2 \ll \omega_0^2$ we obtain the so-called "shortened" wave equation for the HF field amplitude, where τ will be used as the parameter only:

$$2ik_0\frac{\partial E}{\partial \xi} + \Delta_\perp E + \frac{\omega_0^2 - k_0^2c^2}{c^2}E = \frac{\omega_{pe}^2}{c^2}\,\frac{1 + \delta n_e/n_0}{\sqrt{1+I}} \tag{III.11}$$

and the change of the plasma electron concentration is described by equation

$$\frac{\partial^2}{\partial t^2}\frac{\delta n_e}{n_0}=\frac{c^2}{\omega_{pe}^2}\left(\frac{\partial^2}{\partial\tau^2}+\omega_{pi}^2\right)\left(\frac{1}{V_g^2}\frac{\partial^2}{\partial\tau^2}+\Delta_\perp\right)\sqrt{1+I} \qquad \text{(III.12)}$$

Investigating the linear theory of the filamentational instability, we represent the pumping as the sum of the plane wave of high amplitude A and perturbations with the amplitude a :

$$E=A+a\ ,\quad I_0=\frac{e^2AA^*}{m_e^2\omega_0^2c^2}\ ,\quad I_1=\frac{e^2(aA^*+a^*A)}{m_e^2\omega_0^2c^2}\ ,\quad a\ll A,\ I_1\ll I_0 \qquad \text{(III.13)}$$

$$I_1\sim e^{\varkappa\xi}\exp i(k_\perp r_\perp-\Omega t) \qquad \text{(III.14)}$$

Substituting (III.13) and (III.14) into (III.11) and (III.12) we find the increment of the filamentational instability

$$\varkappa^2=\frac{k_\perp^4}{4k_0^2}\left\{\frac{\omega_{pe}^2}{k_\perp^2c^2}\frac{I_0}{(1+I_0)^{3/2}}+\frac{\Omega^2+k_\perp^2V_g^2}{k_\perp^2V_g^2}\frac{\Omega^2-\omega_{pi}^2}{\Omega^2}\frac{I_0}{1+I_0}-1\right\} \qquad \text{(III.15)}$$

where

$$k_0^2=\frac{\omega_0^2}{c^2}-\frac{\omega_{pe}^2}{c^2}\frac{1}{\sqrt{1+I_0}} \qquad \text{(III.16)}$$

The first term in brackets is caused by relativistic nonlinearity, the second one - by striction, and the third one corresponds to the beam diffraction.

The maximum increment

$$\varkappa_{max}^2=\frac{\omega_{pe}^4}{16k_0^2c^4}\left(\frac{1+\sqrt{1+8\Omega^2/\omega_{pi}^2}}{3+\sqrt{1+8\Omega^2/\omega_{pi}^2}}\right)^2\left[1+\frac{\omega_{pi}^2}{2\Omega^2}\left(1+\sqrt{1+8\frac{\Omega^2}{\omega_{pi}^2}}\right)\right]^{-1} \qquad \text{(III.17)}$$

is achieved at

$$\bar{k}_\perp^2=\frac{1}{2}\frac{\omega_{pe}^2}{c^2}\frac{I_0}{1+\frac{\omega_{pi}^2}{\Omega^2}I_0}\frac{1}{\sqrt{1+I_0}} \qquad \text{(III.18)}$$

for

$$\bar{I}_o = \frac{1}{2}\left(1+\sqrt{1+8\Omega^2/\omega_{pi}^2}\right) \qquad \text{(III.19)}$$

when $\omega_{pi}^2 \gg \Omega^2$; $\bar{I}_o = 1$ and $\varkappa_{max}^2 = \omega_{pe}^4/(32 k_o^2 c^4)$ and
if $\omega_{pi}^2 \sim \Omega^2$; $\bar{I}_o = 2$ and $\varkappa_{max}^2 = \omega_{pe}^4/(100 k_o^2 c^4)$

For the plasma particle concentration we obtain

$$\frac{\delta n_e}{n_o} = -\frac{\omega_o^2}{k_o^2 c^2}\,\frac{\Omega^2 + k_\perp^2 V_g^2}{\Omega^2}\,\frac{\Omega^2-\omega_{pi}^2}{\omega_{pe}^2}\,\frac{I_1}{2\sqrt{1+I_o}} \qquad \text{(III.20)}$$

$$\frac{\delta n_i}{n_o} = \frac{m_e}{m_i}\,\frac{\omega_o^2}{k_o^2 c^2}\,\frac{\Omega^2 + k_\perp^2 V_g^2}{\Omega^2}\,\frac{I_1}{2\sqrt{1+I_o}} \qquad \text{(III.21)}$$

As it is seen from (III.20),(III.21), $\delta n_e > 0$ only in cases when $\omega_{pi}^2 > \Omega^2 > 0$ or $\Omega^2 < -k_\perp^2 V_g^2$.

When $\omega_{pi}^2 \ll \Omega^2 \ll k_\perp^2 V_g^2$ the value of the increment coincides with that of the so-called "stationary case" and the plasma is ejected from the filaments [37]. The maximum increment in this case is

Let us consider the self-focusing of short electromagnetic pulses with Gauss distribution of intensity with the time and also in the direction transverse to the distribution of the beam on the plasma boundary.

$$I(z=0) = I_o(t)\, e^{-\frac{r_\perp^2}{a^2}} = I_{oo}\, e^{-\frac{t^2}{t_o^2}}\, e^{-\frac{r_\perp^2}{a^2}} \qquad \text{(III.22)}$$

where a - is the effective transverse width of the beam on plasma boundary, t_o - is pulse duartion. We solve (Eq.(III.11)) for weak-relativistic intensities ( ) in the paraxial aberrationless approximation.

Substituting:

$$I = \frac{I_{oo}}{f^2(\xi,\tau)}\, e^{-\frac{\tau^2}{t_o^2}}\, e^{-\frac{r^2}{a^2 f^2(\xi,\tau)}} \qquad \text{(III,23)}$$

where dimensionless width of the beam

$$f = \left\{ 1 - \beta \left(I_{oo} e^{-\frac{\tau^2}{t_o^2}} - I_{cr} \right) \right\}^{1/2} \qquad \text{(III.24)}$$

$I_{cr} = \frac{2c^2}{\omega_{pe}^2 a^2}$ - critical intensity, $\beta = \omega_{pe}^2 \xi^2 / (2K_o^2 c^2 a^2)$ into Eqs.(III.6) and (III.7), we can find the formed spatial charge on the beam axis:

$$Q = -e\left(\frac{\delta n_e}{n_o} - \frac{\delta n_i}{n_o}\right)\Big|_{r_\perp = 0} = -e\left\{ -\frac{2c^2}{\omega_{pe}^2 a^2 f^2} \frac{I_o(\tau)}{f^2}\left[\frac{a^2 f^2 (1+\beta I_o/f^2)}{2V_g^2 t_o^2}\left(1 - \frac{2\tau^2}{t_o^2}\right) + 1\right] + \frac{4\tau^2}{t_o^2}\frac{c^2}{\omega_{pe}^2 V_g^2 t_o^2} \frac{I_o(\tau)}{f^4}\left(1 + \beta\frac{I_o}{f^2}\right)\right\} \qquad \text{(III.25)}$$

At $2\tau^2/t_o^2 < \frac{1}{2} + \frac{V_g^2 t_o^2}{a^2 f^2 (1+\beta I_o/f^2)}$ the ion density near the axis of the beam excees the electron density. The spatial charge in the focus region can become significant. If the longitudinal size of the packet is much greater than the beam transverse width ($a^2 \ll c^2 t_o^2$):

$$Q\big|_{\tau=0} = e\,\frac{I_{cr}}{f^2}\,\frac{I_{oo}}{f^2} \qquad \text{(III.26)}$$

and for the opposite approximation

$$Q\big|_{\tau=0} = e\,\frac{1}{\omega_{pe}^2 t_o^2}\,\frac{I_{oo}}{f^2} \qquad \text{(III.27)}$$

The resulting charge can move in the laboratory system with the velocity close to the velocity of light and can excite wake field in plasma, which can be used for particle acceleration [30,43].

For electron density on the axis at $\tau = 0$ we have

$$\left.\frac{\delta n_e}{n_0}\right|_{\substack{r_\perp=0\\ \tau=0}} = \begin{cases} -\dfrac{I_{cr}}{f^2}\dfrac{I_{c0}}{f^2}, & \dfrac{a^2}{c^2} \ll t_0^2 \ll \omega_{pi}^{-2} \\ -\dfrac{1}{\omega_{pe}^2 t_0^2}\dfrac{I_{c0}}{f^2}, & \dfrac{a^2}{c^2}, \omega_{pi}^{-2} \gg t_0^2 \\ -\dfrac{m_e}{m_i}\dfrac{c^2 t_0^2}{a^2}\dfrac{I_{c0}}{f^2}, & \dfrac{a^2}{c^2}, \omega_{pi}^{-2} \ll t_0^2 \\ \dfrac{m_e}{2m_i}\dfrac{I_{00}}{f^2}, & \dfrac{a^2}{c^2} \gg t_0^2 \gg \omega_{pi}^{-2} \end{cases} \tag{III.28}$$

In the first three cases electrons are ejected from the axis area. In the first two cases the pulse durations are so small, that ions can be considered stationary. In the third case after the electron ejection from the focus area in the direction transverse to the beam propagation the ions are also extracted from it. For long pulses ($(m_e/m_i)(ct_0/a)^2 \gg 1$) all plasma will be displaced ($\delta n_e/n_0 \to -1$) from the axis area. The nonlinearity will be saturated. After the self-focusing up to cartain dimensions the beam again begins to axpand. The self-trapping channels are formed.

Unlike other cases, in the forth case plasma compression takes place in the areas of the moving focuses. Due to the non-stationarity of the pumping wave the longitudinal component of the ponderomotive force may exceed the transverse one and the rate of particle raking in the longitudinal direction exceeds the rate of particle ejection in the transverse direction. The beam self-focuses and compresses plasma. This particle is particularly interesting for heating and compression of plasma by intensive relativistic beams, unlike all other cases when the ejection of plasma electrons takes place.

Let us consider the case when $\frac{a^2}{c^2} \ll t_0^2 \ll \omega_{pi}^{-2}$ in more detail [37–41]. As it is seen from (III.28) in the focus areas approximation $\delta n_e/n_o << I << 1$ is vilated and the density disturbance becomes significant. The solution of the self-consistent task, i.e. the solution of Eq.(III.11) together with equation for electron concentration becomes necessary:

$$\frac{\delta n_e}{n_o} = \frac{c^2}{\omega_{pe}^2} \Delta_\perp \sqrt{1 + I} \qquad (III.29)$$

For cimplicity we solve the system of equations (III.11)9(III.29) with the help of numerical calculations for the one-dimensional beam ($\Delta_\perp = \frac{\partial^2}{\partial x^2}$, i.e. the width of the beam along axis y is much greater than the width of the beam a along axis X). The disturbances on different harmonics I_1 , I_2 ,... are superimposed on the main homogeneous distribution.

For any distribution of intensity on plasma boundary which is maximum on the axis, in this area $\delta n_e/n_o < 0$. It means that in such case both electron striction and relativistic effect favour beam self-focusing and filamentation. In this case the electron ejection from the axis area increases. Finally, such a moment may come when all the electrons from the focal areas will be ejected (Fig.III.1). Here beam diffraction will exceed nonlinear refraction. Central beams will diverge from the axis, while peripherical ones will converge, i.e. the beam after self-focusing will begin to expand and the so-called "side-band focuses" will be formed on the periphery. Intensity on the axis will become less than on the periphery. The electron displacement from the side-band focuses will take place. Part of them will be ejected into the axis area. The concentration on the axis will grow and nonlinear refraction of the beam will predominate over refraction again.And again the beam will begin focusing (Fig.III.1). The self-supporting channel is formed. This situation will recur till aberrations disturb the unity of the beam. The beam will be divided into many filaments and will finally begin to defocuse.

The more boundary intensity exceeds the critical one the quicker self-focusing takes place and the longer are channels in which electron concentration falls of to zero (Fig.III.2).

As it was seen from Eqs.(III.28) in isotropic and collisionless plasma when the width of the beam exceeds the longitudinal dimensions of the packet ($a^2 > c^2 t_c^2$), while pulse duration exceeds the characteristic time of ion oscillation ($t_o > \omega_{pi}^{-1}$), self-focusing of the beam is accompanied by plasma compression. Analogous picture may be expected at self-focusing and filamentation of relativistically intensive electromagnetic beams in the longitudinal magnetic field, which should interfere with the ejection of plasma electrons from the axis area of the beam in the transverse direction.

If the wave frequency ω_o greatly exceeds cyclotron electron frequency and $\Omega_i t_o \gg 1$ (where Ω_i is ion-cyclotron frequency) then we shall obtain the system of equations which consists of wave equation (III.5) and equation for electron concentration:

$$\frac{\delta n_e}{n_o} = \frac{m_e}{m_i} \frac{c^2}{\Omega^2} \left(\Delta_\perp + \frac{\Omega^2}{V_g^2} \right) \left(\sqrt{1+I} - \sqrt{1+I_o} \right) \tag{III.30}$$

where $\Omega^2 = (\Omega_i \omega_{pi})^2 / (\omega_{pi}^2 + \Omega_i^2)$ (as a rule $\Omega^2 \simeq \Omega_i^2$).

Increment of filamentation instability has the form:

$$\gamma = \frac{K_\perp^2}{2K_o} \left\{ \frac{\omega_{pe}^2}{K_\perp^2 c^2} \frac{I_o}{(1+I_o)^{3/2}} + \frac{\omega_{pi}^2}{\Omega^2} \frac{I_o}{1+I_o} - 1 \right\}^{1/2} \tag{III.31}$$

If $(\omega_{pi}/\Omega_i)^2 I_o < 1$, at $\bar{K}_\perp^2 = \frac{\omega_{pe}^2}{2c^2} \frac{I_o}{1+I_o} \left(1 - \frac{\omega_{pi}^2}{\Omega_i^2} I_o\right)^{-1}$ the increment has the maximum value:

$$\gamma_{max} = \frac{\omega_{pe}^2}{4K_0C^2}\frac{I_0}{1+I_0}\left(1-\frac{\omega_{pi}^2}{\Omega_i^2}I_0\right)^{-1/2}$$

(III.32)

At the self-focusing of the beam with Gauss profile of intensity distribution in the point $z=0$, in the weak-relativistic case ($\delta n_e/n_0 \ll I \ll 1$) in paraxial aberrationless approximation the disturbance of electron density on the axis has the form:

$$\left.\frac{\delta n_e}{n_0}\right|_{r_\perp=0} = \frac{m_e}{2m_i}\frac{I_0}{f^2}\left(1-f^2-\frac{4V_g^2}{a^2\Omega^2 f^2}\right)$$

(III.34)

It is seen that of $\delta \equiv (a\Omega/4V_g)^2 > 1$ at self-focusing when the beam dimansionless width varies in the interval $\left[\frac{1}{2}(1-\sqrt{1-1/\delta}), \frac{1}{2}(1+\sqrt{1-1/\delta})\right]$ the compression of plasma takes place in the focus. At $f_0^2 = 8V_g^2/(\Omega a)^2$ the dicturbance of electron density reaches the maximum value:

$$\frac{\delta n_e}{n_0} = \frac{m_e}{2m_i} I_0 \left(\frac{a^2\Omega^2}{V_g^2}-1\right)$$

(III.34)

The similar picture is observed in the case of arbitrary relativistic beams. The filamentation instability of plane wave has been investigated with the help of numerical calculations. The dynamics of formation and evolution of one filament has been investigated. The obtained results can be subdivided into two quantitavity different groups. The first, where self-focusing is accompained by electron ejection from the focus, and at high intensities ($I \gg 1$) these striction effect may become appreciable ($\frac{\delta n_e}{n_0} \to 1$) (Fig.III.3). In the second group in the filament area the plasma electron density first increases and then begins to decrease. Each filament begins to divide and several narrower filaments are formed (Fig.III. 4). In the case of sufficiently broad beams the large longitudinal magnetic fields prevent the electron ejection from the focus or filament region and the longitudinal component of ponderomotive force may exceed the resulting force acting in the transverse direction. The wave "rakes up" electrons in the focus region.

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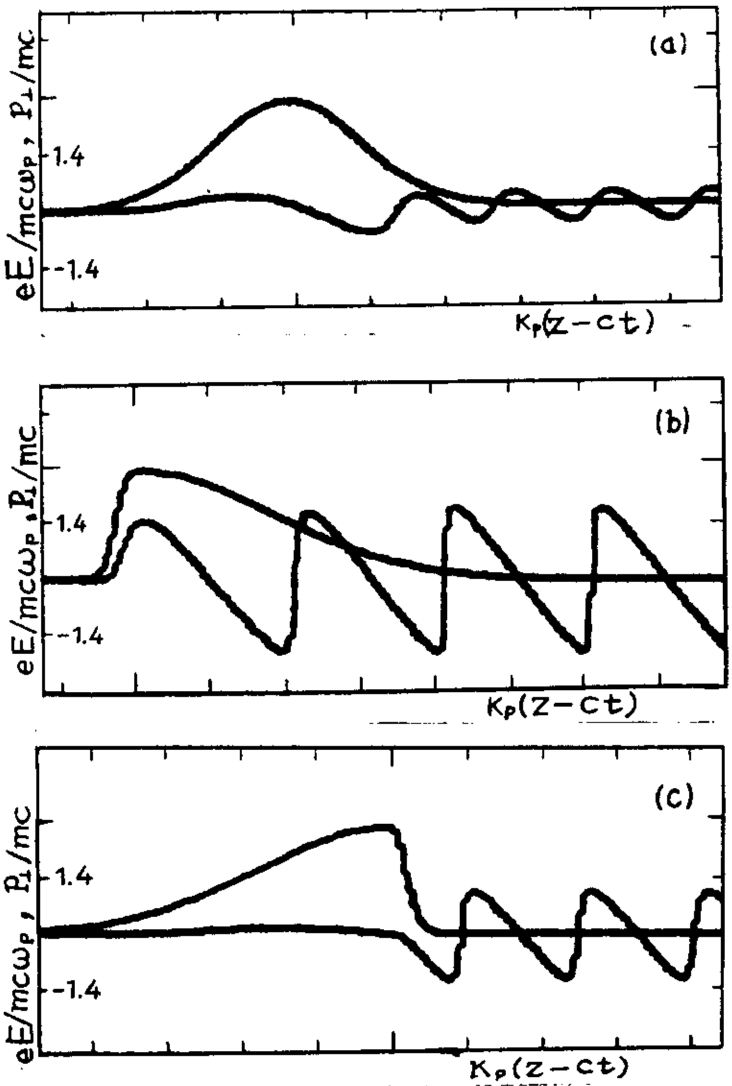

Fig. II.1 Numerical solutions of wake-fields produced by various pulse shapes: (a) Gaussian rise and fall, $\sigma_r = \sigma_f = 5\,c/\omega_p$; (b) Gaussian rise, $\sigma_r = c/\omega_p$, Gaussian fall, $\sigma_f = 9\,c/\omega_p$; (c) $\sigma_r = 9\,c/\omega_p$, $\sigma_f = c/\omega_p$.

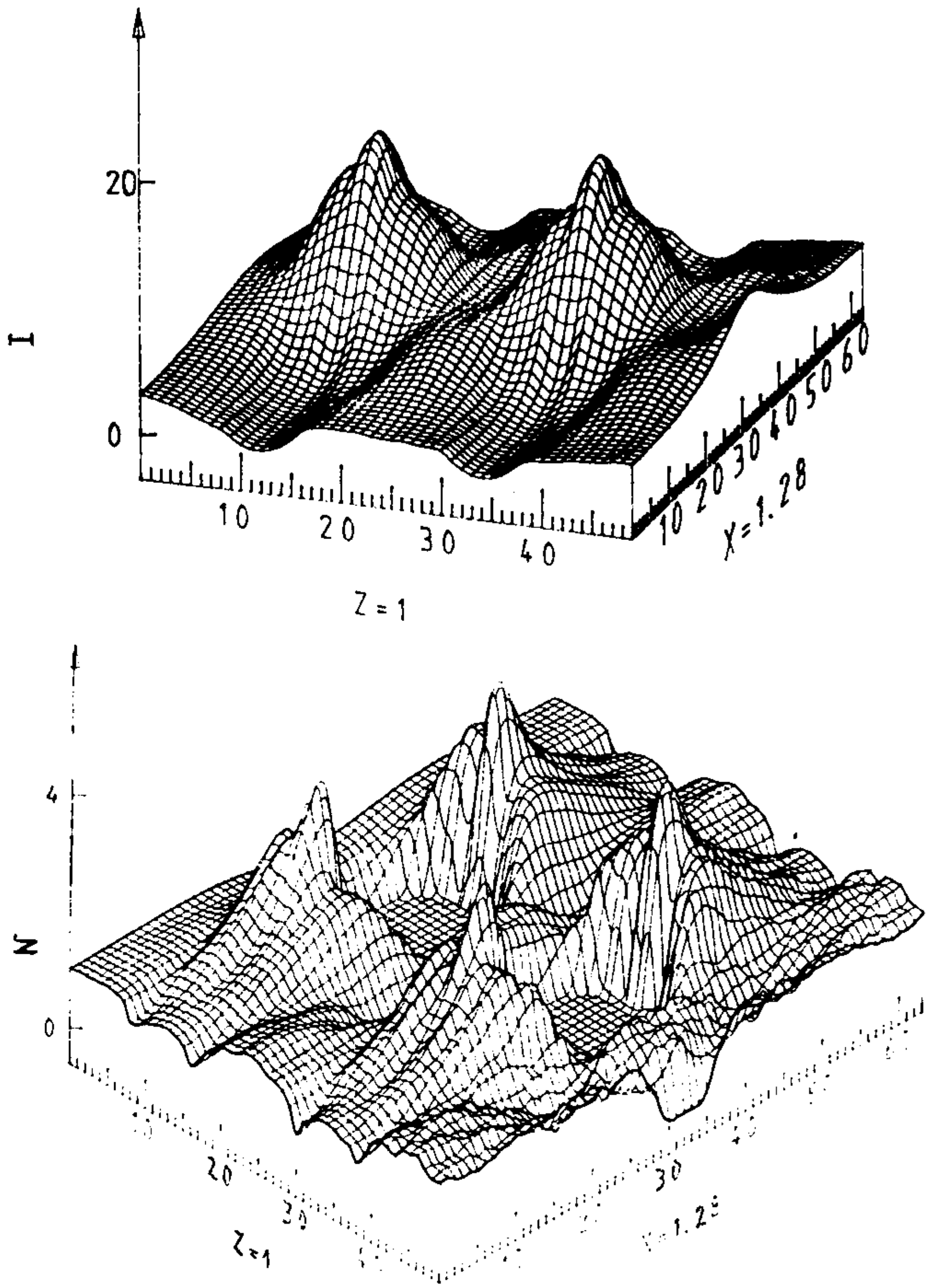

Fig. III.1 Distribution of intensity and concentration , when

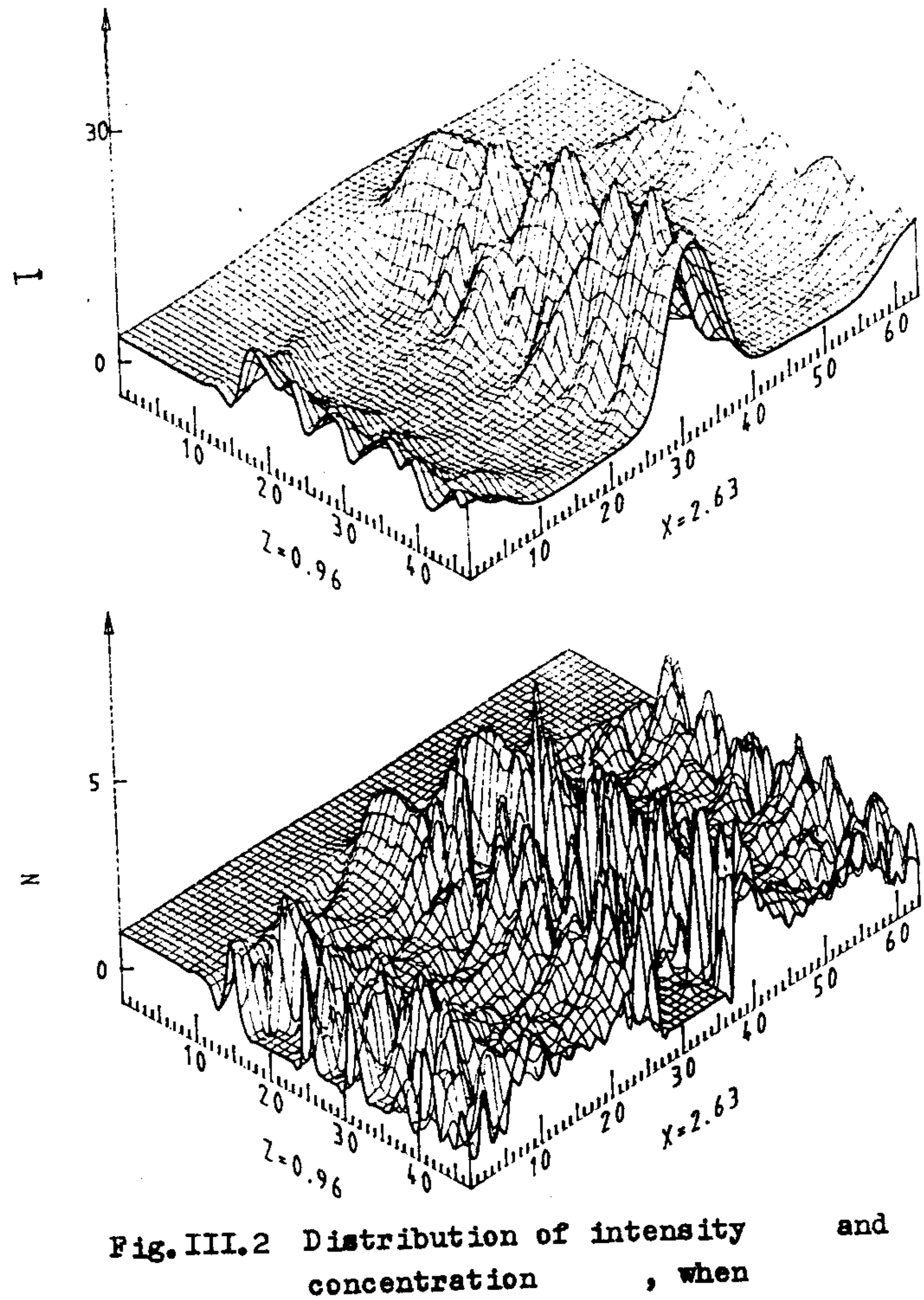

Fig. III.2 Distribution of intensity and concentration , when

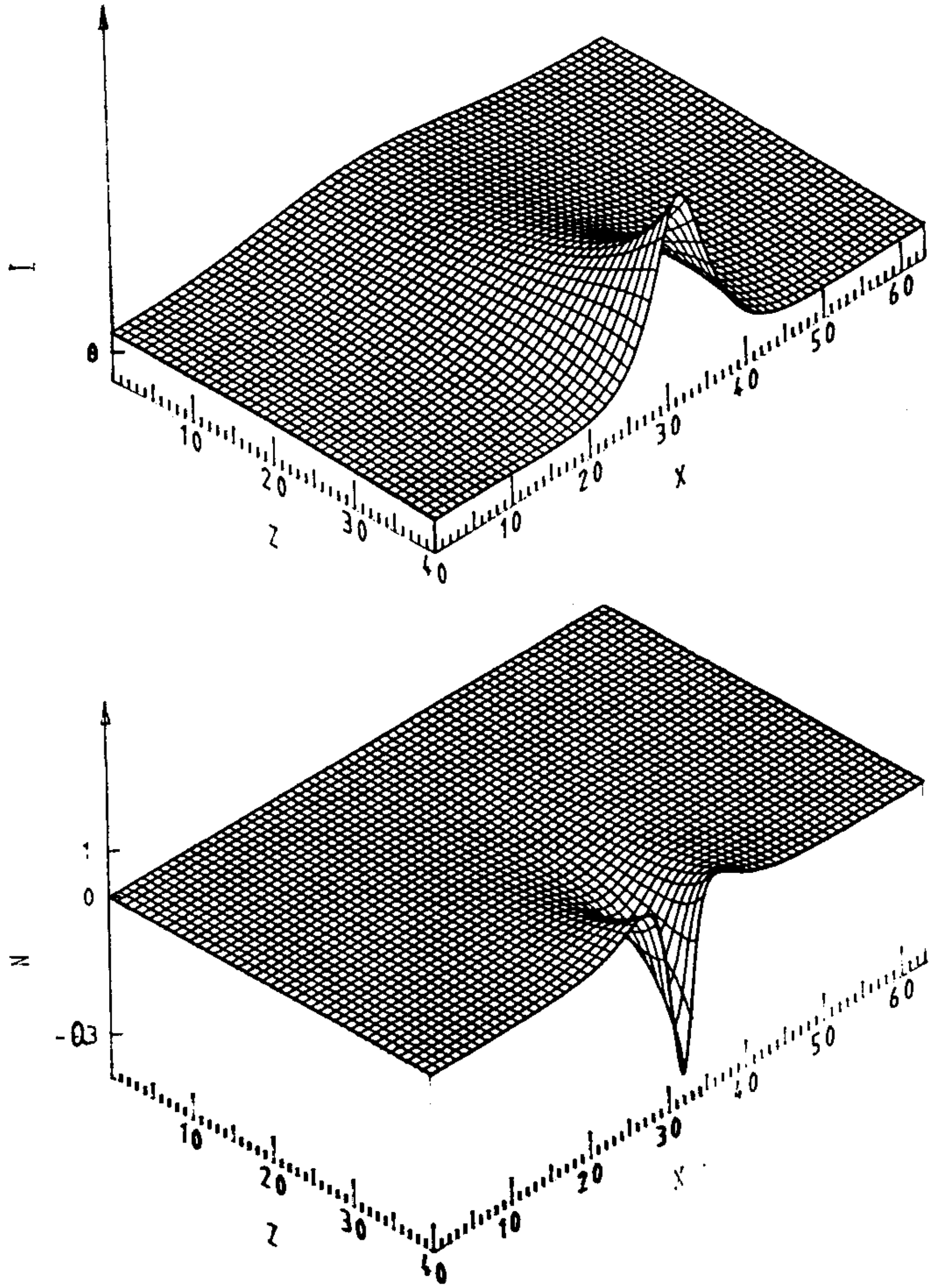

Fig. III.3 Distribution of intensity and concentration , when

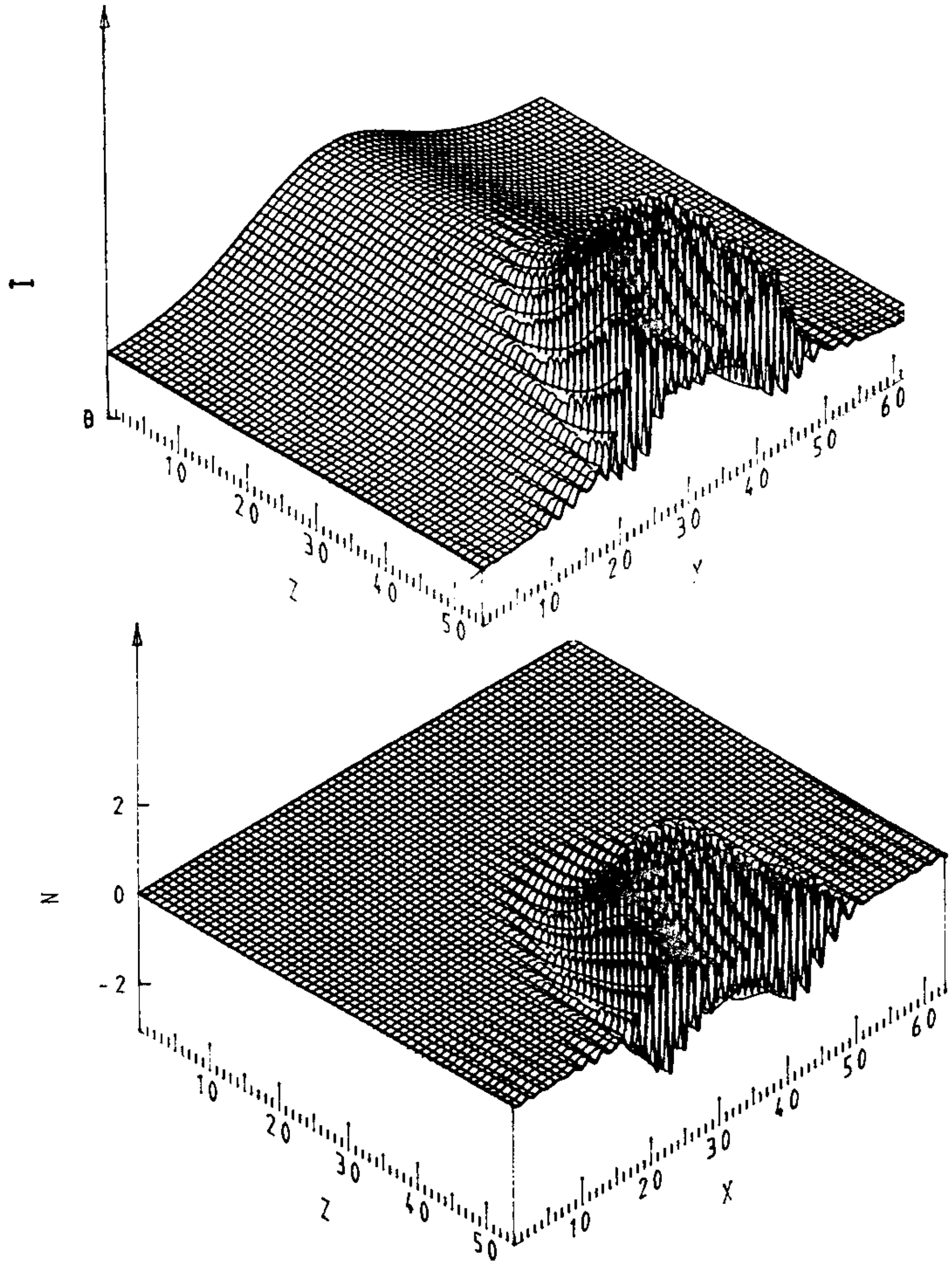

Fig. III.4 Distribution of intensity and concentration , when