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MAGNETIC FIELD GENERATION AND PONDEROMOTIVE FORCE EFFECTS

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The present paper considers the problems of spontaneous magnetic field generation, heat and particle flux and formation in collisionless plasmas under the influence of the ponderomotive force. It was shown [23] that in the hydrodynamic approximation the spontaneously generated magnetic field makes no direct effect on the development of Langmuir plasma turbulence: one has to take into account both magneto-modulational and for purely relativistic effects, i.e. the velocity-dependence of the electron mass.

The effect of spatial and time inhomogeneity of the pumping field on the particle distribution in both isotropic and magnetoactive plasmas is studied. Inhomogeneity of the external field is shown [24] to produce particle and heat fluxes in the plasma.

I. At present, the problem of low-frequency (LF) spotaneous magnetic field generation attracts certain attention due to its importance for understanding the dynamics of Langmuir plasma turbulence (see, e.g. [1], [9], [11]).

Papers [1-3], which we will be henceforth referred to as BKK, study the spontaneous LF magnetic field effect on the nonlinear dynamics of Langmuir waves in plasmas and, in particular,

on the modulational instability and the collapse.

In our opinion, the papers of BKK contain certain inconsistencies. With the latter being removed one comes to modified conclusions on the nature of nonlinear processes under consideration.

1. In the hydrodynamic approximation the spontaneous LF magnetic field makes no direct effect on the nonlinear dynamics of turbulent plasmas. In the expression for the fast current, BKK disregarded the term proportional to the velocity of the slow electron motion. This term compensates the current density due to the spontaneous LF magnetic field.

2. According to the results of BKK and neglecting the skin-effect, the LF vortex current contribution is proportional to the relativistic factor $\frac{V_E^2}{c^2}$, where V_E is the high-frequency (HF) electron velocity amplitude and c is the light velocity. The contribution of the purely relativistic effect, i.e. the HF-electric-field-amplitude-dependence of the electron mass, is shown in the present paper to be of the same order of magnitude. When compared to the magneto-modulational effect, the relativistic effects turn and to be the main reason for the HF wave modulational instability to occur.

We start from the kinetic equations for the electron and ion plasma components and Maxwell equations:

$$\frac{\partial f_e}{\partial t} + (\vec{v}, \nabla) f_e - \frac{e}{m_e}\left(\vec{E} + \frac{1}{c}[\vec{v}, \vec{B}]\right)\frac{\partial f_e}{\partial \vec{v}} = 0 \tag{1.1}$$

$$\frac{\partial f_i}{\partial t} + (\vec{v}, \nabla) f_i + \frac{e}{m_i}\left(\vec{E} + \frac{1}{c}[\vec{v}, \vec{B}]\right)\frac{\partial f_i}{\partial \vec{v}} = 0 \tag{1.2}$$

$$\operatorname{div} \vec{E} = 4\pi e \int (f_i - f_e)\, d\vec{v} \tag{1.3}$$

$$\mathrm{rot}\,\vec{B} = \frac{1}{c}\frac{\partial\vec{\mathcal{E}}}{\partial t} - \frac{4\pi e}{c}\int\vec{v}(f_e - f_i)\,d\vec{v} \tag{1.4}$$

$$\mathrm{rot}\,\vec{\mathcal{E}} = -\frac{1}{c}\frac{\partial\vec{B}}{\partial t} \tag{1.5}$$

$$\mathrm{div}\,\vec{B} = 0 \tag{1.6}$$

where e is the electron charge, m_e and m_i are the electron and ion masses, respectively.

From the kinetic equation (1.1), we obtain the ones for the electron distribution function moments $n_e = \int f_e\,d\vec{v}$, $n_e\vec{u} = \int\vec{v}f_e\,d\vec{v}$

$$\frac{\partial n_e}{\partial t} + \mathrm{div}\,n_e\vec{u} = 0 \tag{1.7}$$

$$\frac{\partial u_m}{\partial t} + (\vec{u},\nabla)u_m = -\frac{e}{m_e}\mathcal{E}_m - \frac{1}{m_e n_e}\nabla_n\bar{P}_{mn} - \frac{e}{m_e c}[\vec{u},\vec{B}]_m \tag{1.8}$$

the pressure tensor being $\bar{P}_{mn} = \int(\vec{v}-\vec{u})_m(\vec{v}-\vec{u})_n f_e\,d\vec{v}$

The system of equations obtained is not closed since it contains higher moments $\bar{P}_{mn}$.

As it is shown in the paper [4], this system can be closed both for the unperturbed state and for the induced oscillations with double plasma frequency $\omega_e\left(\omega_e = \left(\frac{4\pi e^2 n_o}{m_e}\right)^{1/2}\right)$. As for LF oscillations, the relevant procedure is not adequate.

With the HF field present, the values of interest for contain both slow time dependence and fast dependence with the characteristic time $\tau \sim \omega_e^{-1}$. Therefore, we may assume each of the values $A \equiv (\vec{\mathcal{E}}, \vec{B}, n_e, \vec{u})$ to be of the following form:

$$A = \langle A\rangle + \tilde{A} \tag{1.9}$$

where the brackets denote averaging over the time segment τ.

We restrict the consideration to the case when the fast ion motion may be neglected, and the constant magnetic field is absent, $\vec{B}_o = 0$. Moreover, we introduced the following notation: $\langle\vec{J}_e\rangle = \vec{V}_d$; $\langle\vec{B}\rangle = \delta\vec{B}$ (where $\delta\vec{B}$ is the LF perturbation of the magnetic field); $\langle n_e\rangle = n_o + n_d$ (n_o is the unperturbed electron density, while n_d is the LF perturbation of the electron density due to the averaged HF potential influence on the LF motion $(n_o \gg n_d)$).

According to [4], [13], we present all fastly varying values in the series form:

$$\tilde{A} = A_1\exp(-i\omega_e t) + A_2\exp(-2i\omega_e t) + c.c. \tag{1.10}$$

assuming $|\vec{\mathcal{E}}_1| \gg |\langle\vec{\mathcal{E}}\rangle|, |\vec{\mathcal{E}}_2|$; $n_o \gg n_1 \gg n_d, n_2$; $|\vec{V}_1| \gg |\vec{V}_d|, |\vec{V}_2|$. That allows us to develop the perturbation theory. In the first perturbation order, the electron velocity $\vec{u}_1 = \vec{V}_1\exp(-i\omega_e t) + \vec{V}_1^*\exp(i\omega_e t)$, where $\vec{V}_1 = -\frac{ie}{m_e\omega_e}\vec{\mathcal{E}}_1$. The electron distribution function in the unperturbed plasma

$$f_o^e = f_o^e(\vec{v} - \tilde{\vec{u}}_1) \tag{1.11}$$

satisfies the equation

$$\frac{\partial f_o^e}{\partial t} - \frac{e}{m_e}\left[\vec{\mathcal{E}}_1\exp(-i\omega_e t) + c.c.\right]\frac{\partial f_o^e}{\partial\vec{v}} = 0 \tag{1.12}$$

The relation (1.11) implies the major part of electrons to move as a whole. As follows from (1.1) and (1.11), the additional term f_e' in the distribution function f_o^e satisfies the equation

$$\frac{\partial f_e'}{\partial t} + (\vec v,\nabla)(f_0^e + f_e') - \frac{e}{m_e}\left\{\langle \vec E\rangle + \frac{1}{c}[\vec v,\delta\vec B]\right\}\frac{\partial (f_0^e + f_e')}{\partial \vec v} -$$

$$-\frac{e}{m_e}\left(\tilde{\vec E}_2 + \frac{1}{c}[\vec v,\tilde{\vec B}]\right)\frac{\partial f_0^e}{\partial \vec v} - \frac{e}{m_e}\left(\tilde{\vec E} + \frac{1}{c}[\vec v,\tilde{\vec B}]\right)\frac{\partial f_e'}{\partial \vec v} = 0 \qquad (1.13)$$

Then it is convenient to make the substitute the variable $\vec V \to \vec v + \vec u_1$ and thus to rewrite (1.13) as

$$\frac{\partial f_e'}{\partial t} + (\vec v + \vec u_1,\nabla)f_e' - \frac{\partial (f_0^e + f_e')}{\partial \vec v}\left\{(\vec v + \vec u_1,\nabla)\vec u_1 + \frac{e}{m_e}\left(\langle \vec E\rangle - \tilde{\vec E}_2\right) +\right.$$

$$\left. + \frac{e}{m_e c}\left[\vec u_1, \delta\vec B + \tilde{\vec B}\right]\right\} - \frac{e}{m_e c}\left[\vec v, \delta\vec B + \tilde{\vec B}\right]\cdot\frac{\partial f_e'}{\partial \vec v} = 0$$

(1.14)

Now we separate the time scales in this equation. We present the function f_e' as a series:

$$f_e' = \langle f\rangle + f_1 \exp(-i\omega_e t) + f_2 \exp(-2i\omega_e t) + c.c.$$

In the first order of the perturbation theory we obtain:

$$f_1 = \frac{i}{\omega_e}(\vec v,\nabla)\vec V_1\cdot\frac{\partial f_0^e}{\partial \vec v} \qquad (1.15)$$

When deriving (1.15), we have assumed that $\omega_e \gg k V_{Te}$ where $V_{Te} = \left(\frac{T_e}{m_e}\right)^{\frac{1}{2}}$ is the electron thermal velocity, k^{-1} is the characteristic scale of the HF motion. Substituting this expression into the pressure gradient which enters (1.8), we obtain:

$$\nabla_m \bar P_{mn} = -\frac{3 i n_0 V_{Te}^2}{\omega_e}\nabla_n \operatorname{div}\vec V_1 = -3V_{Te}^2\nabla_n n_1$$

(1.16)

Substituting (1.16) into (1.8) by means of the well-known procedure [5], [6], we obtain truncated equation for the slowly varying complex amplitude, $\tilde{E} = 2\vec E_1$, of the HF elect-

ric field at the principal frequency ω_e :

$$i\vec E_t + \frac{3}{2}\omega_e r_e^2\nabla \operatorname{div}\vec E - \frac{c^2}{2\omega_e}\operatorname{rot}\operatorname{rot}\vec E - \frac{\omega_e}{2}\frac{n_d}{n_0}\vec E +$$

$$+\frac{i}{2}\left\{\left[\vec E, \frac{e}{m_e c}\delta\vec B - \operatorname{rot}\vec V_d\right] + \vec V_d \operatorname{div}\vec E + \vec V_2 \operatorname{div}\vec E^* - \nabla(\vec E^*\vec V_2) +\right.$$

$$\left. -(\vec V_d\nabla)\vec E\right\} + \frac{\omega_e}{2}\frac{n_2}{n_0}\vec E^* + \frac{e^2}{16 m_e^2 c^2\omega_e}\left[3|\vec E|^2\vec E + \left[\vec E[\vec E^*,\vec E]\right]\right] = 0$$

(1.17)

where r_e is the electron Debye radius.

In this equation we take into account both electron nonlinearities and a weak relativistic effect (the last term in (1.17)) associated with the dependence of the pump-wave-amplitude electron mass. Moreover, equation (1.17) takes into account the effect of the generated IF magnetic field (see BKK) which produces the term $i\frac{e}{m_e c}[\vec E,\delta\vec B]$ in (1.17). We show in what follows this term to be of the same order as the relativistic effect. Therefore, it is necessary to take both these into account in order to obtain a correct pattern of the phenomenon.

To close the equation (1.17), one has to calculate n_d, $\vec V_d$, n_2, $\vec V_2$ in terms of the distribution function moments $\langle f\rangle$ and f_2 :

$$n_d = \int\langle f\rangle d\vec v, \qquad \vec V_d = \frac{1}{n_0}\int \vec v\langle f\rangle d\vec v$$

$$n_2 = \int f_2\, d\vec v \qquad \vec V_2 = \frac{1}{n_0}\int \vec v f_2\, d\vec v$$

(1.18)

(The terms n_2 and $\vec{V}_2$ may be calculated from the hydrodynamic equations with the pressure gradient disregarded). Having calculated the moments f_2 (see (1.15)) made use of (1.18), one can easily obtain the following relations:

$$\frac{n_2}{n_0} = \frac{e^2}{6m_e^2\omega_e^4}\left[\frac{1}{4}\Delta(\vec{E})^2 + \operatorname{div}(\vec{E}\operatorname{div}\vec{E})\right] \tag{1.19}$$

$$\vec{V}_2 = \frac{i e^2}{12 m_e^2 \omega_e^3}\left[\nabla(\vec{E})^2 + \vec{E}\operatorname{div}\vec{E}\right] \tag{1.20}$$

Now we derive the equations for slowly varying values $\vec{\delta B}$, V_1, n_1. We first write the equation for $\langle f \rangle$ that may be obtained from (1.14) by means of averaging over the time interval τ. Here we assume $f_0^e(\vec{v})$ to be the Maxwell distribution function and employ the relation $\vec{B}_1 = \frac{m_e c}{e}\operatorname{rot}\vec{V}_1$. Then the required equation becomes:

$$\frac{\partial\langle f\rangle}{\partial t} + (\vec{v},\nabla)\langle f\rangle - \frac{\partial f_0^e}{\partial\vec{v}}\left(\frac{e}{m_e}\langle\vec{E}\rangle + \nabla|\vec{V}_1|^2\right) + $$
$$+ \frac{i}{\omega_e}(\vec{v},\nabla)\vec{\mathcal{D}}\,\frac{\partial f_0^e}{\partial\vec{v}} = 0 \tag{1.21}$$

where

$$\vec{\mathcal{D}} = (\vec{V}_1^*,\nabla)\vec{V}_1 - c.c. \tag{1.22}$$

Using the solution (1.21) for $\langle f\rangle$ and relation (1.18), one may obtain expressions for $\vec{V}_1$ and n_1. Here it is convenient to use Fourier transformation

$$\langle f\rangle_k = \frac{1}{(2\pi)^4}\int\langle f\rangle\exp(-i\vec{k}\vec{r} + i\Omega t)\,d\vec{r}\,dt \tag{1.23}$$

Having defined from (1.21) the Fourier amplitude $\langle f\rangle_k$ for

the Fourier component of the drift velocity $\vec{V}_1$, we find:

$$\vec{V}_{1k} = \frac{i}{\omega_e}\vec{\mathcal{D}}_k - \vec{k}|V_1|_k^2\frac{1}{\Omega}\left(1 - J_+(\beta)\right) + \frac{ie}{m_e\Omega}\left[\frac{\vec{k}(\vec{k},\langle\vec{E}\rangle_k)}{k^2}\beta^2\left(1 - J_+(\beta)\right) - \right.$$
$$\left. - \frac{[\vec{k}[\langle\vec{E}\rangle_k,\vec{k}]]}{k^2}J_+(\beta)\right] + \frac{i}{\omega_e}\left[\frac{\vec{k}(\vec{\mathcal{D}}_k,\vec{k})}{k^2}\beta^2\left(1 - J_+(\beta)\right) - \right.$$
$$\left. - \frac{[\vec{k}[\vec{\mathcal{D}}_k,\vec{k}]]}{k^2}J_+(\beta)\right] \tag{1.24}$$

where

$$\beta = \frac{\Omega}{kV_{Te}}, \qquad J_+(x) = x\exp\left(-\frac{x^2}{2}\right)\int_{i\infty}^{x} d\tau\exp\left(\frac{\tau^2}{2}\right)$$

In what follows we use the following asymptotics of the function $J_+(x)$:

$$J_+(x) = 1 + \frac{1}{x^2} + \frac{3}{x^4} + \dots - i\left(\frac{\pi}{2}\right)^{1/2} x\exp\left(-\frac{x^2}{2}\right) \tag{1.25}$$

for $x \gg 1$, $|\operatorname{Re} x| \gg |\operatorname{Im} x|$

and $J_+(x) = -i\left(\frac{\pi}{2}\right)^{1/2} x$, when $|x| \ll 1$

We could have obtained a relation similar to (1.24) for the ion drift velocity as well. However, we assume below the phase velocities of LF wave pulsations to be high as compared to the ion thermal velocity, $\Omega \gg kV_{Ti}$, and then general hydrodynamic equations are applicable for the description of ion motion.

From (1.21) and the Poisson equation, one can easily obtain an equation for the Fourier-component n_1 which may be written in the general form as follows:

$$V_{Te}^2 k^2 n_{1k}\left(1 - J_+(\beta)\right)^{-1} = - \frac{\omega_{ci}^2 V_{Ti}^2 k^2 n_{1k}}{V_{Ti}^2 k^2 + \omega_i^2\left(1 - J_+(\beta_i)\right)} - k^2 |\vec{V}_e|^2 n_0 \tag{1.26}$$

where $\beta_i = \frac{\Omega}{k V_{Ti}}$

Now we write one more equation for the LF perturbation of the magnetic field $\overline{\delta B}$. It is easy to obtain from (1.4) anc (1.5)

$$\left(\frac{1}{c^2}\frac{\partial^2}{\partial t^2} - \Delta\right)\overline{\delta B} = \frac{4\pi e}{c} n_0 \,\mathrm{rot}\left(\vec{V}_i - V_d - \left\langle \frac{\tilde{n}_e \tilde{\vec{u}}_e}{n_0} \right\rangle\right) \tag{1.27}$$

where

$$\left\langle \frac{\tilde{n}_e \tilde{\vec{u}}_e}{n_0} \right\rangle = \frac{i}{16\pi n_0 m_e \omega_e}\left(\vec{E}\,\mathrm{div}\vec{E}^* - c.c.\right) \tag{1.28}$$

Note that the ion drift velocity is related to the magnetic field LF perturbation as

$$\mathrm{rot}\,\vec{v}_i = -\frac{e}{m_i c}\overline{\delta B} \tag{1.29}$$

Thus, we have obtained a closed system of equations (1.17), (1.19), (1.20), (1.24), (1.26), (1.27) describing the dynamics of longitudinal and transverse HF and LF fields in a plasma with non relativistic temperature, $T_e \ll m_e c^2$. These equations are of general nature, since they take into account the following non-linear effects in the quadratic approximation with respect to the HFfield amplitude: striction nonlinearity, excitation of LF magnetic fields, relativistic and electron nonlinearities. These equations are applicable for the description of nonlinear effects in the plasma both in the "hydrodynamic" limit (i.e. when the LF motion phase velocities exceed the electron thermal

velocity, $\Omega \gg k V_{Te}$) and in the opposite "kinetic" limit, when $\Omega \ll k V_{Te}$. Below we consider these cases separately.

b) We consider the case of the phase velocity of induced LF oscillations being much higher than the electron thermal velocity $\omega_e \gg \Omega \gg k V_{Te}$. Then we easily obtain the following expression

$$\frac{\partial \vec{V}_d}{\partial t} = -\frac{e}{m_e}\langle \vec{E} \rangle - \nabla |\vec{V}_e|^2 \tag{1.30}$$

from the relation (1.24).

Taking the curl part of the equation (1.30), we obtain, within the context of (1.5) the following useful relation:

$$\overline{\delta B} = \frac{m_e c}{e}\,\mathrm{rot}\,\vec{V}_d \tag{1.31}$$

This enables us to write the final equation for the magnetic field perturbation. Using (1.31) and neglecting the displacement current and the ion contribution in (1.27), we find

$$\left(\Delta - \frac{\omega_e^2}{c^2}\right)\overline{\delta B} = \frac{ie}{4 m_e \omega_e c}\,\mathrm{rot}\left(\vec{E}\,\mathrm{div}\vec{E}^* - c.c.\right) \tag{1.32}$$

It is to be noted that a similar equation for $\overline{\delta B}$ was derived in the paper [3] from the hydrodynamic equations. The right-hand part of their equation is as follows (see Eq. (3.21) of the present paper):

$$\frac{ie}{4 m_e \omega_e c}\,\mathrm{rot}\,\mathrm{rot}\,[\vec{E}, \vec{E}^*] \tag{1.33}$$

As it is shown below, the latter expression is valid only in the kinetic case ($\Omega \ll k V_{Te}$), while the calculation of the "slow" current in the hydrodynamic approximation carried out by the above-mentioned authors is incorrect.

The equation for the electron density perturbation in the hydrodynamic limit may be written from the relation (1.26) under the assumption $\Omega \gg k V_{Te}$ $(\beta \gg 1)$. It is of the following form:

$$\frac{\partial^2}{\partial t^2} n_d = \left(1 + \frac{1}{\omega_i^2}\frac{\partial^2}{\partial t^2}\right) \Delta \frac{|\vec{E}|^2}{16\pi m_i} \tag{1.34}$$

Moreover making use of the Poisson equation

$$\mathrm{div}\langle \vec{E} \rangle = 4\pi e(\delta n_i - n_d) \tag{1.35}$$

(δn_i is the ion density perturbation) and the hydrodynamic equations for ions

$$\frac{\partial}{\partial t}\delta n_i + n_0\, \mathrm{div}\, \vec{v}_i = 0 \tag{1.36}$$

$$\frac{\partial \vec{v}_i}{\partial t} = \frac{e}{m_i}\langle \vec{E} \rangle \tag{1.37}$$

one can easily obtain

$$n_d = \left(1 + \frac{1}{\omega_i^2}\frac{\partial^2}{\partial t^2}\right)\delta n_i \tag{1.38}$$

It follows from (1.34) and (1.33) that

$$\frac{\partial^2}{\partial t^2}\delta n_i = \Delta \frac{|\vec{E}|^2}{16\pi m_i} \tag{1.39}$$

and

$$n_d = \delta n_i + \Delta \frac{|E|^2}{16\pi m_i \omega_i^2} \tag{1.40}$$

Thus, we have obtained a closed system of equations in the hydrodynamic approximation (1.17), (1.19), (1.20), (1.30), (1.32), (1.34), (1.38) with regard for the striction non-

linearity, LF magnetic field excitation, electron and relativistic nonlinearities, and quasi-neutrality violation in the LF motion. Taking into account relation (1.31), one can easily show equation (1.17) not to contain explicitly the LF excitation of the magnetic field $\delta\vec{B}$. According to (1.17), the LF magnetic field can effect the HF field dynamics only through the terms involving the drift velocity $\vec{V}_d$ (note that in general $\mathrm{rot}\, \vec{V}_d \neq 0$ and the curl part of V_d is related to $\delta\vec{B}$).

It can be shown that if the characteristic space (L) and time (t_s) scales of the slow motion satisfy the condition $\frac{L}{t_s} \gg \left(\frac{m_e}{m_i}\right)^{1/2} c$, the striction nonlinearity may be neglected as compared to the relativistic nonlinearity [7]. Moreover, if $L \gg \frac{c}{\omega_e}$, then the electron nonlinearities may be neglected as compared to the relativistic nonlinearity. As shows, the equation (1.32) in this case the c ontribution of the terms associated with the generated LF magnetic field $\delta\vec{B}$ in (1.17) is also negligible.

Then we obtain the following nonlinear equation for the HF field amplitude:

$$i\vec{E}_t - \frac{c^2}{2\omega_e}\mathrm{rot}\,\mathrm{rot}\,\vec{E} + \frac{\omega_e}{16}\,\frac{e^2}{m_e^2 c^2 \omega_e^2}\left(3|\vec{E}|^2\vec{E} + [\vec{E}[\vec{E}^*,\vec{E}]]\right) = 0 \tag{1.41}$$

It is evident from the equation (1.41) that the characteristic time of nonlinear processes which are governed by the relativistic effect, is of the following order:

$$\frac{1}{\tau_R} \sim \frac{\omega_e}{16}\,\frac{|V_E|^2}{c} \qquad \vec{V}_E = \frac{e\vec{E}}{m_e \omega_e} \tag{1.42}$$

According to (1.32), the generation of LF magnetic fields occurs with the same characteristic times. (Note that the authors of the paper [10] came to a similar conclusion). However, the back influence of this LF magnetic field on the HF field is insignificant.

As one can see from (1.17) and (1.32), in case $L \ll \frac{c}{\omega_e}$, the relativistic effects and the terms associated with $\delta \vec{B}$ are of the same order of magnitude. On the contrary, in case $L \gg \frac{c}{\omega_e}$, the relativistic nonlinearity (as well as the terms associated with the generated LF magnetic field) is much smaller than the electron nonlinearities, and the characteristic time of the processes under consideration is of the order of the electron nonlinearity characteristic time

$$\frac{1}{\tau_e} \sim \omega_e \left(\frac{r_e}{L}\right)^2 \frac{V_E^2}{V_{Te}^2} \qquad (1.43)$$

These are the characteristic times of the LF magnetic field generation.

Thus, the case of so-called "magneto-modulational" instability, i.e. the case of HF wave instability produced by the generation of LF magnetic field, is actually excluded in the hydrodynamic approximation. Nevertheless, the magnetic field generation takes place, and characteristic growth rates of the magnetic field generation are determined by the relativistic and electron nonlinearities in case of $\frac{L}{\tau_s} \gg \left(\frac{m_e}{m_i}\right)^{1/2} c$ or by the striction effects in case of $\frac{L}{\tau_s} \ll \left(\frac{m_e}{m_i}\right)^{1/2} c$

c) We consider the most interesting "kinetic" case, i.e. assume the phase velocity of LF perturbations to be much lower than the electron thermal velocity and much higher than

the ion thermal velocity $(\kappa V_{Ti} \ll \Omega \ll \kappa V_{Te})$. Then it is easy to obtain from (1.24) the following relation for $\beta \ll 1$:

$$\operatorname{rot} \vec{V}_1 = \frac{i}{\omega_e} \operatorname{rot} \vec{\mathcal{D}} + \frac{e}{m_e} \frac{1}{c V_{Te}} \frac{1}{(2\pi)^{3/2}} \frac{\partial}{\partial t} \int d\vec{r}' \frac{\delta \vec{B}(\vec{r}',t)}{|\vec{r}-\vec{r}'|^2} \qquad (1.44)$$

Substituting (1.44) into (1.27), taking into account (1.28) and (1.29), and neglecting the displacement current, one can easily derive an equation for the quasi-static perturbation of the magnetic field

$$\Delta \delta \vec{B} - \frac{\omega_e^2}{c^2} \delta \vec{B} - \frac{1}{(2\pi)^{3/2}} \cdot \frac{\omega_e^2}{c^2 V_{Te}} \frac{\partial}{\partial t} \int d\vec{r}' \frac{\delta \vec{B}(\vec{r}',t)}{|\vec{r}-\vec{r}'|^2} =$$

$$= \frac{ie}{4 m_e \omega_e c} \operatorname{rot} \operatorname{rot} [\vec{E}, \vec{E}^*] \qquad (1.45)$$

Eq. (1.44) was written by BKK without the term $\frac{\omega_e^2}{c^2} \delta \vec{B}$. The integral term in (1.44) describes the nonlinear Landau damping or, in other words, the anomalous skin effect of the generated LF magnetic field. Now we obtain an equation for n_1. It follows from (1.26) for $\kappa V_{Ti} \ll \Omega \ll \kappa V_{Te}$, $T_i \ll T_e$

$$\left(\frac{\partial^2}{\partial t^2} - c_s^2 \Delta\right) n_1 - \frac{1}{(2\pi)^{3/2}} \left(\frac{m_e}{m_i}\right)^{1/2} c_s \left(1 + \frac{1}{\omega_i^2} \frac{\partial^2}{\partial t^2}\right) \Delta \frac{\partial}{\partial t} \int d\vec{r}' \frac{n_1(\vec{r}',t)}{|\vec{r}-\vec{r}'|^2} =$$

$$= \left(1 + \frac{1}{\omega_i^2} \frac{\partial^2}{\partial t^2}\right) \Delta \frac{|E|^2}{16 \pi m_i} \qquad (1.46)$$

In (1.46), we take into account both the terms $\frac{1}{\omega_i^2} \frac{\partial^2}{\partial t^2}$ and the sound wave damping (the integral term in the left-hand part of (1.46)). A simple analysis shows that in contrast to the integral term in (1.45), the one in (1.46) makes a considerable contribution when the phase velocity of the forced waves is of the order of the ion-sound velocity $c_s = \left(\frac{T_e}{m_i}\right)^{\frac{1}{2}}$ [12]. Since we are not interested in this case hereinafter, we shall neglect the ion sound damping in (1.46).

Thus, we have obtained a closed system of general equations (1.17), (1.19), (1.20), (1.24), (1.46) in the kinetic approximation which describes the dynamics of the HF and the related LF fields. Note that this system of equations for the one-dimensional case $\vec{\delta B} = 0$ has been investigated in the paper [6] ; the authors found new regions and branches of modulational instability both for longitudinal and transverse waves and showed the supersonic compression soliton to be able to propagate in the plasma.

Here we are interested in the LF magnetic field generation by HF waves and in the back influence of $\vec{\delta B}$ on the HF wave dynamics.

As it was mentioned above, such problem has been solved in the papers of BKK. However, they took into account only the striction effect and the effect due to the excitation of LF magnetic field ($\frac{e}{c}[\vec{E}, \vec{\delta B}]$ term) in the equation of (1.17) type. Meanwhile, it is clear from Eqs. (1.17) and (1.45) that when the integral term (1.45) may be neglected, the relativistic term and the one associated with the LF magnetic field generation are of the same order of magnitude. Meanwhile, if the nonlinear damping is dominant in (1.45), the relativistic effect in (1.18) is greater than the term given rise to by the LF magnetic field excitation.

To illustrate the above statements, we investigate the stability of monochromatic HF Langmuir wave with respect to poter-

tial perturbations. Here we assume that $\omega_i^2 \gg \frac{\partial^2}{\partial t^2}$ and disregard the electron nonlinearities (as it is shown in [4], the electron nonlinearities make no effect on the modulation instability of Langmuir waves).

We rewrite the HF wave amplitude in the form

$$\vec{E} = \exp i(\vec{k}_0\vec{r} - \omega_0 t)\left[\vec{E}_0 + \vec{E}_1 e^{-i\Omega t + i\vec{k}\vec{r}} + \vec{E}_2 e^{i\Omega t - i\vec{k}\vec{r}}\right]$$

and

$$n_j = n_j e^{-i\Omega t + i\vec{k}\vec{r}} \qquad \vec{\delta B} = \vec{\delta B} e^{-i\Omega t + i\vec{k}\vec{r}} \tag{1.47}$$

Substituting (1.47) into (1.17), (1.45) and (1.46) by means of the standard procedure (see, e.g., BKK), we obtain the following dispersion relation:

$$\det |B_{ik}| = 0$$
$$i, k = 1,2,3,4 . \tag{1.48}$$

where

$$B_{11} = \omega_0 + \Omega - \frac{3}{2}\omega_e r_e^2 (k^2 + \vec{k}_0) + \frac{e^2 |E_0|^2}{8 m_e^2 c^2 \omega_e}\left(1 + 2\cos^2\theta_+\right)$$

$$B_{12} = \frac{e^2 E_0^2}{16 m_e^2 c^2 \omega_e}\left[\cos(\theta_- - \theta_+) + 2\cos\theta_- \cdot \cos\theta_+\right]$$

$$B_{13} = \frac{\omega_e}{2} E_0 \cos\theta_+ \quad , \qquad B_{14} = \frac{e E_0}{2 m_e c} \sin\theta_+$$

$$B_{21} = \frac{e^2 E_0^{*2}}{6 m_e^2 c^2 \omega_e}\left[\cos(\theta_- - \theta_+) + 2\cos\theta_+ \cdot \cos\theta_-\right]$$

$$B_{22} = \omega_0 - \Omega - \frac{3}{2}\omega_e r_e^2(\vec{k}-\vec{k}_0)^2 + \frac{e^2|E_0|^2}{8m_e^2c^2\omega_e}\left(1+2\cos^2\theta_-\right),$$

$$B_{23} = \frac{\omega_e}{2}E_0^*\cos\theta_-, \qquad B_{24} = \frac{eE_0^*}{2m_ec}\sin\theta_-$$

$$B_{31} = \frac{k^2E_0}{16\pi m_i n_0}\cos\theta_+, \qquad B_{32} = \frac{k^2E_0}{16\pi m_i n_0}\cos\theta_-$$

$$B_{33} = \Omega^2 - c_s^2k^2; \qquad B_{41} = \frac{eE_0^*}{4m_e\omega_e c}\sin\theta_+$$

$$B_{42} = \frac{eE_0}{4m_ec\omega_e}\sin\theta_-, \qquad B_{44} = i\left(\frac{\pi}{2}\right)^{1/2}\frac{\Omega\,\omega_e^2}{c^2v_{Te}k^3} - 1$$

where

$$B_{43} = B_{34}; \qquad \theta_\pm = \arccos\left[\frac{(\vec{k}_0, \vec{k}_0 \pm \vec{k})}{|\vec{k}_0|\cdot|\vec{k}_0 \pm \vec{k}|}\right]$$

and

$$\omega_0 = \frac{3}{2}\omega_e r_e^2 k_0^2 - \frac{3}{16}\frac{e^2|E_0|^2}{m_e^2c^2\omega_e} \qquad (1.49)$$

Along with the linear shift of the Langmuir wave frequency, we take into account in (1.49) the nonlinear frequency shift associated with the relativistic the pumping-wave-amplitude dependence of the electron mass [8].

For the sake of simplicity, we restrict the consideration to the case of $\vec{k} \perp \vec{k}_0$. Moreover, we neglect the terms associated with the LF electron density perturbation in the dispersion equation (1.48) (as it was shown in [1], the condition $\frac{v_{Te}}{c} \gg \left(\frac{4}{9}\frac{m_e}{m_i}\right)^{1/4}$ is to be then satisfied; in such case the effect of LF magnetic field exceeds the striction influence. The dispersion relation acquires the following form:

$$\left(\Omega^2 - \frac{9}{4}\omega_e^2 r_e^4 k^4\right) - \frac{3}{16}v_{Te}^2k^2\frac{|V_E|^2}{c^2}\left(1 - 4\frac{k_0^2}{k_0^2+k^2}\right) -$$

$$-\frac{3}{8}v_{Te}^2k^2\frac{|V_E|^2}{c^2}\frac{k_0^2}{k_0^2+k^2}\left[i\left(\frac{\pi}{2}\right)^{1/2}\frac{\Omega\,\omega_e^2}{c^2v_{Te}k^3} - 1\right]^{-1} = 0 \qquad (1.50)$$

The second term in (1.50) is due to relativistic effect, while the last term is given rise to by the quasi-static magnetic field generation. If we neglect the relativistic effect, we obtain the dispersion equation investigated in [1].

Now we consider the case of $k_0 \gg k$; then the term responsible for the generation of the LF magnetic field $\delta\vec{B}$ in (1.50) may be neglected in comparison with the relativistic effect. Then in the case of $|V_E|^2/c^2 > 4(r_e k)^2$ an instability occurs. The instability growth-rate attains its maximum value for $k_m = \frac{1}{2^{3/2}}\frac{|V_E|}{c}r_e^{-1}$ and is equal to

$$\operatorname{Im}\Omega_{max} = \frac{3}{16}\omega_e\frac{|V_E|^2}{c^2} \qquad (1.51)$$

The quasi-static magnetic fields will be excited with such growth rate. Now we consider the case of $k_0 \ll k$. In case of $\frac{|V_E|}{v_{Te}} \gg \frac{v_{Te}}{c}$, the nonlinear Landau damping for the quasi-static magnetic field in (1.50) may be disregarded. Then the relativistic term is equal to a half of the term associated with $\delta\vec{B}$ by a factor of 2, the relativistic effect inhibiting the instability development. The instability governed by the magnetic field perturbation occurs for $|V_E|^2/c^2 > 12(r_e k)^2$. Its growth rate attains its maximum value for $k_m = \frac{1}{2\sqrt{6}}\frac{|V_E|}{c}r_e^{-1}$ and is equal to

$$\operatorname{Im}\Omega_{max} = \frac{\omega_e}{16}\frac{|V_E|^2}{c^2} \qquad (1.52)$$

In the case of $|V_E|/V_{Te} \ll V_{Te}/c$, the nonlinear damping of the magneto-static perturbations is dominant, and the relativistic term in (1.50) is much greater than the term associated with $\delta\vec{B}$. In contrast to BKK results, in this case the modulation instability does not occur and the quasi-static magnetic field generation may be neglected in (1.50). Then the relativistic term is smaller then the term connected with $\delta\vec{B}$ by a factor of 2, the relativistic effect hindering the instability development. The instability defined by the magnetic field perturbation takes place for $|V_E|^2/c^2 > 12(r_e K)^2$. Its growth rate reaches its maximum value at $K_m = \frac{1}{2\sqrt{3}} \frac{|V_E|}{c} r_e^{-1}$ and equals

$$\mathrm{Im}\,\Omega_{max} = \frac{\omega_e}{16} \frac{|V_E|^2}{c^2} \qquad (1.52)$$

In case of $|V_E|/V_{Te} \ll V_{Te}/c$, the nonlinear damping of the magnetostatic perturbations is dominant, and the relativistic term in (1.50) is much greater than the term associated with $\delta\vec{B}$. In contrast to BKK results, in this case the modulation instability does not occur and the quasi-static magnetic field is not generated.

d) Let us consider the problem of the excited LF magnetic field influence on the Langmuir collapse [5] . This problem has been also discussed in BKK papers. The authors of this paper have come to the conclusion that the magnetic field excitation favours the Langmuir collapse development. It may be easily shown that the relativistic effect favours the collapse too. In fact, neglecting the electron nonlinearities in (1.17) and repeating the calculations carried out in the above-mentioned

paper, one finds that with regard for the relativistic effect, there appears an additional term

$$-\frac{1}{2} \frac{V_{Te}^2}{c^2} \int d\vec{r} \left(2|\vec{E}|^4 + (\vec{E}^*)^2 \vec{E}^2 \right) \qquad (1.53)$$

in the expression for the plasmon energy (see Eq. (1.13) in the paper [2]). This term is of the same order and sign as the term due to the generated LF magnetic field; according to [2], the latter is equal to

$$-\frac{1}{2} \frac{V_{Te}^2}{c^2} \int d\vec{r} \, |[\vec{E}, \vec{E}^*]|^2 \qquad (1.54)$$

However, as it was mentioned above, the electron nonlinearities may be neglected in comparison with the relativistic effect and the effect associated with $\delta\vec{B}$ only under the condition that $L \gg \frac{c}{\omega_e}$. Thus, these effects influence the collpase only at the initial stage of the caverne decrease. After the characteristic dimension of the caverne becomes smaller than $\frac{c}{\omega_e}$, the electron nonlinearities arise, which determine the character of the further process. The problem of the electron nonlinearity influence on the Langmuir collapse has been discussed in [4] , [14].

II. Recently, much attention has been paid to the problem of HF electromagnetic field effect on the kinetic phenomena in plasma, such as particle velocity redistribution and formation of hot electrons [15],[16] . It was shown in [16] that, in a collisional plasma hot electron "tails" are formed on the Maxwell of distribution function due to the anomalous diffusion. Similar particle redistribution, however, corresponding to the formation of hot electrons, can occur in collisionless plasma, as well as under the effect of the ponderomotive force. For instance, the ponderomotive force effect on the phase-space particle distribution was studied in [17]. Paper [18] is devoted to the kinetic treatment : the powerful electromagnetic field interaction with the plasma. Relativistic motion of electrons in this field is taken into account. However, the results of the above papers were obtained with regard for the spatial dependence of the field amplitude $\vec{E}(z)$ only, the slow-time-dependence being neglected. A new ponderomotive effect was found and analysed in [19]. It is associated with the origin of particle and heat fluxes. This effect is caused by the time-dependence of the pump field amplitude. The plasma was assumed to be charged (ambipolar diffusion, decaying plasma), then there occurs a strong electromagnetic field in the equilibrium state, and the equilibrium distribution function is given by the Maxwell-Boltzmann distribution, i.e.

$$f = \left(\frac{m_e}{2\pi T_e}\right)^{1/2} \exp\left[-\frac{m_e v^2}{2T_e} + \frac{e\varphi(z)}{T_e}\right]$$

which implies that the wave energy is much lower than the electrostatic energy

$$e\varphi/T_e \gg |E|^2/nT_e$$

On the other hand, under the effect of the powerful electromagnetic field with plasma, the ponderomotive force causes the charge separation and formation of the potential field, which in turn modifies the particle distribution in the ordinary space. Just such approach was employed in [18,20].

This paper considers the ponderomotive force effect on the kinetic phenomena in a magnetoactive plasma. The mathematical procedure is similar to that given in [18,20]. In contrast to

[19], the time-dependence of the pump wave amplitude was shown not to be the only reason of heat and particle flux formation. Availability of the magnetic field provides conditions required for the formation of heat and particle fluxes even in the case when the wave amplitude is time-independent. This paper discusses also the problem of the hot electron origin and gives relevant expressions for density and pressure associated with the diffusion term in the kinetic equation.

Consider an electron plasma, exposed to a HF field of a circularly polarized wave

$$\vec{E} = \left\{ E_0(z,t)\cos\left(\omega t - \int k\,dz\right);\ E_0(z,t)\sin\left(\omega t - \int k\,dz\right);\ 0 \right\} \tag{2.1}$$

The external magnetic field $\vec{B}_0$ is directed along the axis Oz and the particle distribution is assumed to vary along the axis Oz.

The induced magnetic field possessess only x - and y - components

$$B_x(z,t) = c\int_{-\infty}^{t} dt'\,\frac{\partial E_y}{\partial z}(z,t')\ ;\ B_y(z,t) = -c\int_{-\infty}^{t} dt'\,\frac{\partial E_x}{\partial z}(z,t') \tag{2.2}$$

Apart, the ponderomotive force results in the shift of the bulk which gives rise to the potential electric field $E_z = -\nabla_z \varphi$. This system is described by the particle distribution function ($f(v_x, v_y, v_z, z, t)$), which satisfies the Vlasov kinetic equation [21]. Let us change the variables and introducing a new function

$$f(v_x, v_y, v_z, z, t) = f\Big\{ v_x - \frac{e}{m}\int\left(E_x + \frac{1}{c}[\vec{v}\,\vec{B}]_x\right)dt',$$
$$v_y - \frac{e}{m}\int\left(E_y + \frac{1}{c}[\vec{v}\,\vec{B}]_y\right)dt',\ v_z,\ z,\ t \Big\} \tag{2.3}$$

which is the distribution function in the coordinate system oscillating with the external field frequency. For the new distribution function (2.3), the kinetic equation becomes

$$\frac{\partial f}{\partial t} + v_z \frac{\partial f}{\partial z} + \frac{\partial}{\partial z}\left(-\frac{e\varphi}{m}\right)\frac{\partial f}{\partial v_z} + \frac{e}{mc}\left\{\frac{\partial}{\partial v_z}\left[\left((v_x + \tilde{v}_x)B_y - (v_y + \tilde{v}_y)B_x\right) f\right]\right\} = 0 \qquad (2.4)$$

where m is electron mass at rest, v_x and v_y are the electron thermal velocity components, $\tilde{v}_x$ and $\tilde{v}_y$ are the electron oscillation velocity components in the external field.

Solving characteristic equations of motion for electrons, it is not difficult to obtain (due to high inertia the ions are assumed to be stationary):

$$\tilde{v}_x = \frac{eE_0}{m\omega}\cdot\frac{\omega - k v_z}{\omega - k v_z - \Omega}\sin\left(\omega t - \int k\,dz\right)$$

$$\tilde{v}_y = \frac{eE_0}{m\omega}\cdot\frac{\omega - k v_z}{\omega - k v_z - \Omega}\cos\left(\omega t - \int k\,dz\right)$$

$$(2.5)$$

where $\Omega = eB_0/mc$ is the electron cyclotron frequency. Using (2.5), equation (2.4) can be rewritten as

$$\frac{\partial f}{\partial t} + v_z \frac{\partial f}{\partial z} + \frac{\partial}{\partial z}\left(-\frac{e\varphi}{m}\right)\frac{\partial f}{\partial v_z} - \frac{\partial}{\partial z}\left(\frac{v_E^2}{2}\right)\frac{\partial}{\partial v_z}(af) +$$

$$+\frac{\partial}{\partial z} v_E\left[v_y \cos\left(\omega t - \int k\,dz\right) - v_x \sin\left(\omega t - \int k\,dz\right)\right]\frac{\partial f}{\partial v_z} = 0 \qquad (2.6)$$

where

$$v_E = \frac{eE_0}{m\omega} \qquad a = (\omega - k v_z)/(\omega - k v_z - \Omega)$$

As is usual for the quasilinear approximation, we present the function f as a slowly varying part and a fast oscil-

lating addition:

$$f = F + \tilde{F}, \qquad \langle f\rangle = F$$

The relevant equations for $\tilde{F}$ and F are of the form

$$\frac{\partial \tilde{F}}{\partial t} + v_z \frac{\partial \tilde{F}}{\partial z} + \frac{\partial}{\partial z}\left(-\frac{e\varphi}{m}\right)\frac{\partial \tilde{F}}{\partial v_z} - \frac{\partial}{\partial z}\left(\frac{v_E^2}{2}\right)\frac{\partial}{\partial v_z}(a\tilde{F}) =$$

$$= -\frac{\partial}{\partial z} v_E\left[v_y \cos\left(\omega t - \int k\,dz\right) - v_x \sin\left(\omega t - \int k\,dz\right)\right]\frac{\partial F}{\partial v_z}$$

$$(2.7)$$

$$\frac{\partial F}{\partial t} + v_z \frac{\partial F}{\partial z} + \frac{\partial}{\partial z}\left(-\frac{e\varphi}{m}\right)\frac{\partial F}{\partial v_z} - \frac{\partial}{\partial z}\left(\frac{v_E^2}{2}\right)\frac{\partial}{\partial v_z}(aF) =$$

$$= -\frac{\omega}{2\pi}\int^{2\pi/\omega} dt\left\{\frac{\partial}{\partial z} v_E\left[v_y \cos\left(\omega t - \int k\,dz\right) - v_x \sin\left(\omega t - \int k\,dz\right)\right]\frac{\partial \tilde{F}}{\partial v_z}\right\} \qquad (2.8)$$

Here the averageing over the period of the external field oscillations has been carried out.

Solving equation (2.7) by the method of characteristics, we obtain

$$\tilde{F} = -\int_{-\infty}^{t}\frac{\partial}{\partial \tilde{z}} v_E\left[v_y \cos\left(\omega\tau - \int k\,d\tilde{z}\right) - v_x \sin\left(\omega\tau - \int k\,d\tilde{z}\right)\right]\frac{\partial F}{\partial \tilde{v}_z}\,d\tau \qquad (2.9)$$

where

$$\tilde{v}_z = \frac{d\tilde{z}}{d\tau}, \qquad \frac{d\tilde{v}_z}{d\tau} = \frac{\partial}{\partial \tilde{z}}\left(-\frac{e\varphi}{m} - \frac{\omega}{\omega - \Omega}\frac{v_E^2}{2}\right) = \frac{\partial}{\partial \tilde{z}}\psi$$

Substituting solution (2.9) into (2.8) and taking into account $k v_z \ll |\omega - \Omega|$, we obtain an equation for the slowly varying part of the distribution function

$$\frac{\partial F}{\partial t} + v_z \frac{\partial F}{\partial z} + \frac{\partial}{\partial z}\left(-\frac{e\varphi}{m}\right)\frac{\partial F}{\partial v_z} - \frac{\alpha}{2}\frac{\partial v_E^2}{\partial z}\frac{\partial F}{\partial v_z} - \frac{\beta}{2}\frac{\partial v_E^2}{\partial z}\frac{\partial}{\partial v_z}(v_z F) -$$

$$-\frac{\gamma}{2}\frac{\partial v_E^2}{\partial z}\frac{\partial}{\partial v_z}(v_z^2 F) = \frac{k^2 v_E^2}{4\omega^2}\left[\frac{\partial v_E^2}{\partial z}\frac{\partial F}{\partial v_z} + \frac{\partial v_E^2}{\partial t}\frac{\partial^2 F}{\partial v_z^2} + v_z\frac{\partial v_E^2}{\partial z}\frac{\partial^2 F}{\partial v_z^2} - 2v_E^2\frac{\partial^2 F}{\partial z\,\partial v_z}\right],$$

$$(2.10)$$

where the following notation is introduced:

$$\alpha = \frac{\omega}{\omega - \Omega}, \quad \beta = \frac{\Omega k}{(\omega-\Omega)^2}, \quad \delta = \frac{\Omega k^2}{(\omega-\Omega)^2}$$

$V_T^2 = \int\limits_{-\infty}^{\infty} (v_x^2 + v_y^2) F \, dv_x \, dv_y$ is the electron thermal velocity.

Equation (2.10) describes the basic quasistationary state of the plasma. Equation (2.10) may be solved by the perturbation method, under the assumption that $F = F_0 + F_1$, $F_1 \ll F_0$ then in the zero-order approximation we have

$$v_z \frac{\partial F_0}{\partial z} + \frac{\partial}{\partial z}\left(-\frac{e\varphi}{m}\right)\frac{\partial F_0}{\partial v_z} - \frac{\alpha}{2}\frac{\partial v_E^2}{\partial z}\frac{\partial F_0}{\partial v_z} = 0 \qquad (2.11)$$

Equation (2.11) is solved by the function

$$F_0 = C \exp\left\{-\frac{m v_z^2}{2T} + \frac{e\varphi}{T} - \frac{\omega}{\omega - \Omega}\cdot\frac{m v_E^2}{2T}\right\} \qquad (2.12)$$

where T is the plasma temperature in energy units, C is the integration constant.

In the first approximation (for the function F), we obtain

$$\frac{\partial F_1}{\partial t} + v_z \frac{\partial F_1}{\partial z} + \frac{\partial}{\partial z}\left(-\frac{e\varphi}{m}\right)\frac{\partial F_1}{\partial v_z} - \frac{\alpha}{2}\frac{\partial v_E^2}{\partial z}\frac{\partial F_1}{\partial v_z} = \frac{\beta}{2}\frac{\partial v_E^2}{\partial z}\frac{\partial}{\partial v_z}(v_z F_0) -$$
$$- \frac{\delta}{2}\frac{\partial v_E^2}{\partial z}\frac{\partial}{\partial v_z}(v_z^2 F_0) + \gamma \frac{d v_E^2}{dt}\frac{\partial^2 F_0}{\partial v_z^2} - 2\gamma v_E^2 \frac{\partial^2 F_0}{\partial z \partial v_z} +$$
$$+ \gamma \frac{\partial v_E^2}{\partial t}\frac{\partial F_0}{\partial z} - \frac{\partial F_0}{\partial t}, \qquad (2.13)$$

where $\gamma = V_T^2/4\omega^2$. The last term in tho oquation (2.13)

appears due to the fact that when equation (2.11) was solved, the time derivative of the zero distribution function was not considered since, for $t \gg \frac{1}{\omega_i}\frac{L}{\tau_D}$, it is much smaller than the second and the third terms of the equation (2.11). However, one cannot neglect the contribution of this correction in the first approximation (L is the spatial scale of the plasma, τ_D is the Debye radius).

Taking into account (2.12), we find from the equation (2.13)

$$F_1 = C\exp\left[-\frac{m v_z^2}{2T} + \frac{e\varphi}{T} - \frac{\omega}{\omega-\Omega}\frac{m v_E^2}{2T}\right]\left[\gamma\frac{m}{T}\left(\frac{m v_z^2}{T} - 2\right)v_E^2 +\right.$$
$$+ \gamma\frac{m}{T}\int\limits_{-\infty}^{t} d\tau \frac{\partial v_E^2}{\partial \tau} - \frac{\beta}{2}\left(\frac{m v_z^2}{T} - 1\right)\int\limits_{-\infty}^{t} d\tau \frac{\partial v_E^2}{\partial z} + 2\gamma\frac{m v_z^2}{T}\int\limits_{-\infty}^{t} d\tau \times$$
$$\times v_E^2 \frac{\partial}{\partial z}\left(-\frac{e\varphi}{T} - \alpha\frac{m v_E^2}{2T}\right) - \int\limits_{-\infty}^{t} d\tau \frac{\partial}{\partial \tau}\left(-\frac{e\varphi}{T} - \alpha\frac{m v_E^2}{2T}\right) -$$
$$\left. - \frac{\delta}{2} v_z \left(\frac{m v_z^2}{T} - 2\right)\int\limits_{-\infty}^{t} d\tau \frac{\partial v_E^2}{\partial z}\right] \qquad (2.14)$$

Substituting the solutions of (2.12) and (2.14), one can write the expession for F as

$$F = C\exp\left[-\frac{m v_z^2}{2T\left(1 + \delta\frac{m v_E^2}{T}\right)} + e\varphi - \alpha\frac{m v_E^2}{2T} - 2\gamma\frac{m v_E^2}{T}\right]\cdot$$
$$\times\left[1 - \frac{\beta}{2}\left(\frac{m v_z^2}{T} - 1\right)\int\limits_{-\infty}^{t} d\tau \frac{\partial v_E^2}{\partial z} - \frac{\delta}{2} v_z\left(\frac{m v_z^2}{T} - 2\right)\int\limits_{-\infty}^{t} d\tau \frac{\partial v_E^2}{\partial z} + \int\limits_{-\infty}^{t} d\tau \frac{\partial}{\partial \tau}\xi\right] \qquad (2.15)$$

where $\xi = \gamma\frac{m v_E^2}{T} + \frac{1}{T}\left(-e\varphi + \alpha\frac{m v_E^2}{2}\right)$

In (2.15) integrals in the square brackets are, obviously, much less than unity, but they are just the reason for the formation of the particle and heat fluxes. The dominant conti:ibution in particle density and pressure is due to the functic f without integral terms, i.e.

$$n_0 = \int_{-\infty}^{\infty} \tilde{f}\, dv_z$$

this condition yields

$$C = N_0 m^{1/2} / \left[2\pi T \left(1 + 2\gamma \frac{m v_E^2}{T} \right) \right]^{1/2}$$

and

$$n_0 = N_0 \exp\left[\frac{e\varphi}{T} - (1 + 4\gamma) \frac{m v_E^2}{2T} \right]$$

and the pressure caused by the function $\tilde{f}$ is of the form

$$P_0 = n_0 T \left(1 + 2\gamma \frac{m v_E^2}{T} \right)$$

Let us find the corrections to the particle density and pressure as well as the expression for the particle and heat fluxes caused by the integral terms.

For the particle density

$$\Delta n = n_0 \int_{-\infty}^{\infty} \varphi \, dv_z \tag{2.17}$$

The mean velocity is determined by

$$\langle v \rangle = \int_{-\infty}^{\infty} v_z \varphi \, dv_z \tag{2.18}$$

For the pressure

$$\Delta P = m n_0 \int_{-\infty}^{\infty} (v_z - \langle v \rangle)^2 \varphi \, dv_z \tag{2.19}$$

To find the heat flux, we use the standard expression

$$q = \frac{m n_0}{2} \int_{-\infty}^{\infty} (v_z - \langle v \rangle)^3 \varphi \, dv_z \tag{2.20}$$

where

$$\varphi = \left[m / 2\pi T \left(1 + \gamma \frac{m v_E^2}{T} \right) \right]^{1/2} \exp\left[- \frac{m v_z^2}{2T \left(1 + 2\gamma \frac{m v_E^2}{T} \right)} \right] M(\tilde{z}, \tilde{v}_z, t) \tag{2.21}$$

$$M(\tilde{z}, \tilde{v}_z, t) = -\left[\frac{\beta}{2} \left(\frac{m v_z^2}{2T} - 1 \right) - \frac{\delta}{2} v_z \left(\frac{m v_z^2}{T} - 2 \right) \right] \int_{-\infty}^{t} d\tau \frac{\partial v_E^2}{\partial z} + \int_{-\infty}^{t} d\tau \frac{\partial \tilde{z}}{\partial \tau} \tag{2.22}$$

If the external magnetic field is absent $\Omega = 0$, then the equations (2.15-2.22) are in accordance with the results obtained in [20] under the condition $v_E^2/c^2 \ll 1$ (nonrelativistic limit).

Unlike the isotropic plasma, equations (2.21-2.22) show that the constant magnetic field appreciably modifies the equation (2.22) and the terms associated with the magnetic field are much greater than the last term in the equation (2.22).

Particle and heat fluxes, the density and pressure modifications are mainly due to the magnetic terms.

Note that the integrand in the equation (2.22) depends on $\tilde{z}(\tau)$ and τ (i.e. $v_E = v_E(\tilde{z}(\tau), \tau)$).

Let us consider for example the effect of a stationary soliton moving with a group velocity v_g on the transport phenomena. In this case v_E in the expression (2.22) is explicitly dependent only on $\tilde{z}(\tau)$; bearing this in mind, one can integrate in the expression (2.22), which becomes

$$M(\tilde{z}, v_z, t) = \frac{1}{v_g} \left[\frac{\beta}{2} \left(\frac{m v_z^2}{T} - 1 \right) - \frac{\delta}{2} v_z \left(\frac{m v_z^2}{T} - 2 \right) \left(1 + \frac{v_z}{v_g} \right) \right] v_E^2 \tag{2.23}$$

where

$$v_E^2 = \frac{v_{E\,max}^2}{\mathrm{ch}^2 \frac{z - v_g t}{\lambda}} \tag{2.24}$$

and $v_z \ll v_g$, λ – is the characteristic soliton width.

In this case, one can make use of (2.17-2.20) and thus find the particle density, mean velocity, modifications of the pressure and the heat flux for the field of the type (2.24).

The calculations yield

$$\Delta n = - \frac{\Omega_0 E_0^2 I N_0}{2 m v_g^2 (\omega - \Omega)^3} v_{E\,max}^2 \mathrm{ch}^{-2} \frac{1}{\lambda} (z - v_g t) \tag{2.25}$$

$$\langle v \rangle = \frac{\Omega K T}{m v_g(\omega-\Omega)^2} v_{E\max}^2 \, ch^{-2} \frac{1}{\lambda}(z - v_g t) \qquad (2.26)$$

$$\Delta p = \frac{\Omega K T N_0}{v_g(\omega-\Omega)^2} v_{E\max}^2 \, ch^{-2} \frac{1}{\lambda}(z - v_g t) \qquad (2.27)$$

$$q = \frac{3 \Omega K T^2 N_0}{m v_g(\omega-\Omega)^2} v_{E\max}^2 \, ch^{-2} \frac{1}{\lambda}(z - v_g t) \qquad (2.28)$$

As follows from (2.25), Δn can change the sign depending upon the sign of the difference $\omega - \Omega$. Physically it is associated with the fact that the striction interaction leads to the plasma migration out of the strong field range for $\omega > \Omega$, and for $\omega < \Omega$ plasma is compressed [22]. It is worth mentioning that under the action of HF electromagnetic field on the plasma, the quasistationary electron distribution is be Maxwellian with the new temperature $T + \frac{k^2 v_E^2}{2\omega^2}\frac{m v_E^2}{2}$ (it follows from (2.15). The addition to the temperature results in the electron redistribution and the increase of the number of fast particles on the tail of the Maxwell distribution.

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