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ABSTRACT

The cavern inertial deepening occuring after the absorption of collapsed Langmuir waves is not taken into account in the existing theoretical models of a strong Langmuir turbulence of isothermic plasma. Meanwhile, such a cavern deepening is accompanied by the suction of new Langmuir waves and can drastically change the energetic balance of a system. In the work presented here the qualitative theory of Langmuir turbulence is constructed, which takes into account the postcollaptical effects. The spectra obtained for Langmuir waves and accelerated electrons differ substantially from those predicted earlier. An interesting feature of new spectra is their dependence on the collapse symmetry.

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INTRODUCTION

The supersonic collapse of Langmuir waves, soon after its prediction in Ref. 1, was already considered as the most important elementary act of a strong Langmuir turbulence. Usually, it was assumed that, at not too small damping of sonic fluctuations, an energy of turbulence is mainly absorbed in caverns being at the final stage of a collapse (see, for example, Ref. 2). Only recently, it was noticed that even more substantial energy absorption can occur in the process of postcollaptical evolution of caverns via sucking into them new Langmuir waves [3].

Most distinctively this effect is evidenced under the conditions, when the intercavern distance exceeds noticeably the characteristic length of Langmuir waves k_0^{-1}. In order to avoid the onset of caviations elsewhere, the average density of a strong Langmuir turbulence W_0 should estimately be the same as the modulational instability threshold W_{th} for the waves of the energy-containing spatial scale k_0^{-1}, being quantitavely somewhat lower than W_{th} (the extent to what it is lower is determined by the concentration of caverns N_{cav} necessary for the absorption of an energy flux coming into the main scale of a strong turbulence k_0^{-1} from the source external with respect to the strong turbulence). Namely such a regime of a strong Langmuir turbulence is considered below. For the modulational instability threshold the following estimate is valid

$$W_{th} \sim n_0 T k_0^2 r_D^2 , \qquad (1.1)$$

where r_D, n_0, T are respectively the Debye radius, average concentration and temperature (common for electrons and ions). The collapse is further considered to be supersonic, which is justified, if the characteristic group velocity of Langmuir waves v_g is much larger than the ion thermal velocity c_i. Using the well-known estimates for the dispersive adding to the Langmuir wave frequency $\omega_0 \sim \omega_{pe} k_0^2 r_D^2$, one can rewrite the condition $v_g \gg c_i$ in the form

$$1 \gg \left(\frac{m_e}{m_i}\right)^{1/2} \frac{1}{k_0 r_D} \equiv g \tag{1.2}$$

(ω_{pe} is an electron plasma frequency, m_e and m_i are masses of an electron and ion respectively). In the process of a supersonic collapse an energy density $W_{cav}(a)$ of Langmuir waves in the central part of the cavern increases inversely to the cube of its size a, the perturbation of ions concentration $\tilde{n}(a)$ is given by an estimate $|\tilde{n}(a)| \sim n_0 r_D^2/a^2$ and an energy density $W_{cav}^i(a)$ of a directed motion of ions, removing out of the cavern by the trapped Langmuir waves, is of the same order as the value $\frac{|\tilde{n}(a)|}{n_0} W_{cav}(a)$ and grows as a^{-5}. By the moment of absorption of Langmuir waves initially trapped by the cavern, its size decreases down to the value $a_f \ll a_0 = k_0^{-1}$. In this moment the value $W_{cav}^i(a)$ is a_0/a_f times larger than that for the energy density of a sound with a wavelength a_f and an amplitude $|\tilde{n}(a_f)| \sim n_0 r_D^2/a_f^2$. After the absorption of primary Langmuir waves the cavern is continuing to deepen inertially during the time $\tau_f \sim \frac{a_f}{c_i}$ and achieves the depth

$$\tilde{n}_f \sim \left(\frac{a_0}{a_f}\right)^{1/2} |\tilde{n}(a_f)| \, . \tag{1.3}$$

The estimate obtaned is justified under the condition $W_{cav}(a_f) \lesssim n_0 T$, i. e. at

$$k_0 a_f \gtrsim (k_0 r_D)^{2/3}. \tag{1.4}$$

Otherwise, an electron nonlinearity is developed, which changes qualitatively the collapse dynamics [4]. The condition of smallness for the ion nonlinearity $|\tilde{n}_f| \ll n_0$ is somewhat softer than (1.4):

$$k_0 a_f \gtrsim (k_0 r_D)^{4/5}. \tag{1.5}$$

For the time of the cavern postcollaptical deepening the free Langmuir waves transmit to it the following energy

$$\mathscr{E}_g \sim v_g \tau_f a_0^2 W_0 \sim g^{-1} k_0 a_f W_0 a_0^3, \tag{1.6}$$

a noticeable portion of which can be absorbed. An energy trapped when the cavern is formed and absorbed at the final stage of a collapse is equal to

$$\mathscr{E}_0 \sim W_0 a_0^3. \tag{1.7}$$

The postcollapse absorption of an energy can dominate at

$$k_0 a_f \gg g. \tag{1.8}$$

At $k_0 a_f \lesssim g$ an energy $\mathscr{E}_g$ does not exceed $\mathscr{E}_0$, but it is absorbed by the slower, than during the collapse, electrons and can significantly change their distribution over the velocities. From further considerations it will be clear that the diffusion of electrons on Langmuir waves sucked into caverns after the collapse should be taken into account at

$$k_0 a_f \gtrsim g^{8/7}. \tag{1.9}$$

If

$$k_0 r_D \ll \left(\frac{m_e}{m_i}\right)^{6/19}, \tag{1.10}$$

there exists a domain of parameter $k_0 a_f$ values

$$(k_0 r_D)^{2/3} \ll k_0 a_f \ll g^{8/7}, \tag{1.11}$$

where both the electron nonlinearity and electron diffusion on the postcollaptically absorbed Langmuir waves are negligible.

In the region (1.11) the hypotheses are justified, which are the basis of previous models of a strong Langmuir turbulence. The spectral density of wave energies in the so-called inertial range of scales

$$k_0 \ll k \ll k_f = a_f^{-1}$$

is determined in all these models in the following way. The inverse deepening time for the cavern of a size $a \sim k^{-1}$ is estimated as a growth rate for modulational instability of trapped waves:

$$\gamma_{mod}(k) \sim \omega_{pi}\left[\frac{W_{cav}(k)}{n_0 T}\right]^{1/2} \sim k_0 c_i \left(\frac{k}{k_0}\right)^{3/2} \qquad (1.12)$$

(ω_{pi} is an ion plasma frequency).

Taking into account an independence of the number of caverns, whose dimensions became smaller, than k^{-1}, per time unit, one can easily evaluate the concentration $N_{cav}(k)$ of caverns with the sizes of order k^{-1} and average energy density $W(k)$ of Langmuir waves with lengths of the same order:

$$W(k) \sim \mathscr{E}_0 N_{cav}(k) \sim \mathscr{E}_0 N_{cav}(k_0) \frac{\gamma_{mod}(k_0)}{\gamma_{mod}(k)} \sim \frac{\Pi}{k_0 c_i}\left(\frac{k_0}{k}\right)^{3/2}. \qquad (1.13)$$

Here, Π is the density of energy flux coming from outside to the main scale of a strong Langmuir turbulence. The assumption mentioned above of a large distance between the neigbouring caverns ($N_{cav}(k_0) \ll k_0^3$) is valid at

$$\Pi \ll \Pi_0 \sim \gamma_{mod}(k_0)\, \mathscr{E}_0 k_0^3 \sim k_0 c_i n_0 T (k_0 r_D)^2. \qquad (1.14)$$

Within the dissipative range ($k \gtrsim k_f$) a number of different wave spectra were predicted. Also varied were the predictions for the distribution of electrons accelerated by the turbulence in the region $v \ll v_f = \omega_{pe} a_f$ (the electrons with velocities $v \gg v_f$ were practically absent in any model). A thorough analysis [3] revealed the cause of these disagreements and proved the conclusions made in Ref. 5 (to be more exact, in its part devoted to the supersonic collapse). According to [5], the function $W(k)$ decreases exponentially deep into the region $k \gg k_f$. Correspondingly, the concentration $n_e(v)$ of electrons with velocities of the order of $v \ll v_f(t)$ does not practically depend on time and remains to be the same as that at $v_f(t) \sim v$. The concentration $n_e(v_f)$ of electrons accelerated in the given moment was found from the coincidence condition at $k \sim k_f$ of the modulational instability growth rate $\gamma_{mod}(k_f)$ with the decrement of Landau damping

$$\gamma_L(k) \sim \omega_{pe} n_e\left(\frac{\omega_{pe}}{k}\right)/n_0. \qquad (1.15)$$

As a result, one has the following distribution of electrons

$$n_e(v) \sim n_0 \frac{\gamma_{mod}\left(\frac{\omega_{pe}}{v}\right)}{\omega_{pe}} \sim n_0 \frac{c_i}{v_0}\left(\frac{v_0}{v}\right)^{3/2}, \quad v_0 = \frac{\omega_{pe}}{k_0}, \quad v \lesssim v_f. \qquad (1.16)$$

The law for the shift of electron acceleration front $v_f(t)$ into the region of large velocities is determined by the diffusion equation:

$$\frac{dv_f^2}{dt} = D(v_f), \quad D(v) \sim \frac{\omega_{pe}}{m_e n_0} W\left(\frac{\omega_{pe}}{v}\right) \qquad (1.17)$$

and for the time-independent energy flux Π turns out to be quadratic:

$$v_f(t) \propto t^2. \qquad (1.18)$$

When the growing velocity $v_f(t)$ achieves the value $v_{f1} \sim g^{8/7} v_0$, corresponding to the limit of applicability for the second inequality in (1.11), for electrons with velocities of the order

$$v_{m1} \sim v_{f1}\left(\frac{v_{f1}}{v_0}\right)^{1/4} \sim g^{10/7} v_0 \qquad (1.19)$$

the time of diffusion on the postcollaptically absorbed Langmuir waves becomes equal to the time of $v_f(t)$ variation. Further evolution of spectra for waves and particles is strongly influenced by the postcollaptical effects.

2. POSTCOLLAPTICAL EVOLUTION OF CAVERN

Let $V(r)$ be a characteristic velocity for a directed motion of ions at a distance $r \gg a_f$ from the cavern center at the moment of absorption for primarily trapped Langmuir waves. This velocity is acquired by ions under the action of pressure gradient of Langmuir waves mainly when the size of their collapsing bunch estimately coinsides with r value (that takes place during the time of order $\gamma_{mod}^{-1}(r^{-1})$). Whence one has the estimate

$$V(r) \sim \frac{W_{cav}(r)}{m_i n_0 r \gamma_{mod}(r^{-1})} \propto r^{-5/2} \quad (a_f \lesssim r \lesssim a_0). \qquad (2.1)$$

The perturbation of ion concentrations $\tilde{n}(r, 0)$ at the moment when the collapse is completed can easily be evaluated from the continuity equation:

$$|\tilde{n}(r,0)| \sim \frac{V(r)}{r \gamma_{mod}(r^{-1})} \sim n_0 \frac{r_D^2}{r^2} \quad (a_f \lesssim r \lesssim a_0). \qquad (2.2)$$

During the time τ ($\gamma_{mod}^{-1}(a_f^{-1}) \lesssim \tau \lesssim \tau_f$) after collapse completion the perturbation of the concentration $\tilde{n}(r, \tau)$ changes noticeably due to inertial motion of ions in the region determined by the condition : $\tau \gtrsim \gamma_{mod}^{-1}(r^{-1})$:

$$r \lesssim a_*(\tau) = a_0(k_0 c_i \tau)^{2/3}. \tag{2.3}$$

Here the following value of $|\tilde{n}(r, \tau)|$ is achieved:

$$|\tilde{n}(r, \tau)| \sim \frac{V(r)\tau}{r} \sim |\tilde{n}(r, 0)| \gamma_{mod}(r^{-1}) \tau \infty \frac{\tau}{r^{7/2}}. \tag{2.4}$$

For further considerations the sign of the value $\tilde{n}(\vec{r}, \tau)$ is important, which, apparently, is opposite to the sign of div $\vec{V}(\vec{r})$. An average over the angles value of div $\vec{V}(\vec{r})$ is negative, since the radial component of ion veloocity is positive and decreases with the growth of r faster than r^{-2} (see (2.1). Correspondingly, the value $\tilde{n}(\vec{r}, \tau)$, averaged over angles, is positive. If the cavern is assumed to be centrally symmetric, then after the collapse the barrier would appear around, which practically unpenetrable for the Langmuir waves. Though, actually, the cavern is flattened in some direction z. The calculation shows that for the self-similar solution [6] the value div $\vec{V}(\vec{r})$ is positive throughout the entire plane $z=0$ and in the embracing the plane solid angle of the order of unity. Consequently, in the barrier, surrounding the cavern after collapse, the slit remained wide enough for the Langmuir waves—plasmons—to penetrate into the cavern from the outer space. The noticeable probability of reaching the cavern exists only for plasmons with not too large «orbital momenta» $l \sim 1$ with respect to cavern center, as at $l \gg 1$ the «centrifugal barrier» in the Schrödinger equation describing plasmons (see Ref. 1) is substantially higher than the «attracting potential» (2.2). In the region $r \lesssim a_*(\tau)$, where an attracting potential has the form (2.4), the motion of plasmons becomes quasi-classical and they acquire large orbital momenta. The typical wave number of plasmons flying at a distance r from the cavern centre one can evaluate as follows:

$$k(r, \tau) \sim \left[\frac{|\tilde{n}(r, \tau)|}{n_0 r_D^2}\right]^{1/2} \sim \frac{1}{r}\left[1 + \frac{a_*(\tau)}{r}\right]^{3/4}, \quad a_f \lesssim r \lesssim a_0, \tag{2.5}$$

and an orbital momentum $l(r, \tau)$—as $k(r, \tau)r$. The number of bound states, where plasmons are localized at distances of order r from the cavern center mainly, is estimately equal to

$$N(r, \tau) \sim [k(r, \tau)\, r]^3 \sim \left[1 + \frac{a_*(\tau)}{r}\right]^{9/4}, \quad a_f \lesssim r \lesssim a_0. \tag{2.6}$$

The function $N(r, \tau)$ decreases over r, which corresponds to the localization of the majority of the bound states in the region $r \sim a_f$. The total number of bound states

$$N(\tau) \sim N(a_f, \tau) \sim \left[\frac{a_*(\tau)}{a_f}\right]^{9/4} \tag{2.7}$$

increases with time up to the moment τ_f, when $a_*(\tau)$ achieves the value $a_*(\tau_f) \sim (a_0 a_f^2)^{1/3}$, and $N(\tau)$—its maximum value:

$$N(\tau_f) \sim \left(\frac{a_0}{a_f}\right)^{3/4}. \tag{2.8}$$

Then the perturbation of concentration starts to damp and, gradually, the bound states are pushed out from the cavern.

In the axially symmetric cavern the projection m of a plasmon orbital momentum on the specified direction is conserved. Therefore, it is of interest the number $N_m(r, \tau)$ of localized at distances of order r from the cavern center bound states with a fixed value m and, especially, the number of states with $m=0, \pm 1$, into which the plasmons are mainly trapped. For $N_m(r, \tau)$ the following estimate is valid

$$N_m(r, \tau) \sim \frac{N(r, \tau)}{l(r, \tau)} \sim \left[1 + \frac{a_*(\tau)}{r}\right]^{3/2},$$
$$m = 0; \pm 1; \quad a_f \lesssim r \lesssim a_0. \tag{2.9}$$

3. TRAPPING OF PLASMONS

The probability for a plasmon to be trapped by the deepening cavern depends essentially on the parameter $\omega_0 \tau_{trap}$, where ω_0 is a typical dispersive addings to the frequency of a free plasmon, τ_{trap} is the time for its sucking into the cavern to the depth of order ω_0 over the frequency. At $\omega_0 \tau_{trap} \ll 1$, most of plasmons flying against the cavern are trapped. At $\omega_0 \tau_{trap} \gg 1$, the main part of the flux passes freely above the cavern and only plasmons of sufficiently large wavelengths are trapped. Under the assumption that the spectral density of turbulence energy W_k is estimately the same within the entire region $k \lesssim k_0$, for the long-wave plasmons the following estimate takes place:

$$W(k)\sim k^3W_k\sim\left(\frac{k}{k_0}\right)^3 W(k_0)\,,\quad k\lesssim k_0. \tag{3.1}$$

Into every new bound state the plasmons are trapped with wavelengths of order k_{trap}^{-1} (located in the volume of order k_{trap}^{-3} around the cavern). Their energy

$$\mathscr{E}_{trap}\sim k_{trap}^{-3}\,W(k_{trap})\sim k_0^{-3}\,W(k_0)\sim\mathscr{E}_0 \tag{3.2}$$

does not depend on a certain value of $k_{trap}\ll k_0$ and estimately is the same as the energy absorbed via the collapse. Because of a large number of states sucked into the cavern, the postcollaptical absorption of an energy dominated obligatory in the regimes with $\omega_0\tau_{trap}\gg 1$.

The evaluation of «trapping time» τ_{trap} depends on the cavern symmetry. Since the question of stability for the axisymmetric self-similar regime of supersonic collapse of Langmuir waves [6] is not yet solved, one has to consider two possibilities. At axisymmetric cavern the projection m of a plasmon orbital momentum on the specified direction is the integral of motion. The probability for plasmons to be trapped into the state with $|m|>1$ is low, since the wave function of such a state, even at a bound energy close to zero, is localized within the region $r\ll a_0$, into which free plasmons with $|m|>1$ do not penetrate because of the centrifugal barrier. The number of states with $|m|\lesssim 1$ sucked by the cavern during the time τ after collapse is approximately equal to

$$N_0(a_f,\tau)\sim\left[\frac{a_*(\tau)}{a_f}\right]^{3/2}. \tag{3.3}$$

Taking into account that one bound state localized in the region $r\sim a_0$ exists always, the time τ_{trap} for sucking plasmon to the depth of order ω_0 one can evaluate as

$$\tau_{trap}(\tau)\sim\frac{\tau}{N_0(a_f,\tau)}\sim\gamma_{mod}^{-1}(a_f^{-1})\,. \tag{3.4}$$

In this case, the condition of «sharp» trapping $\omega_0\tau_{trap}\ll 1$ has the form

$$k_0a_f\ll g^{2/3}. \tag{3.5}$$

In the essentially nonsymmetric cavern the wave function of practically every state with a bound energy close to zero contains

mainly small orbital momenta $l\lesssim 1$ and, to a noticeable extent, is removed out of cavern. Therefore, in fact, all the states sucked by such a cavern are populated by plasmons and the time τ_{trap} has to be evaluated as follows:

$$\tau_{trap}(\tau)\sim\frac{\tau}{N(a_f,\tau)}\sim\left[\frac{a_f}{a_*(\tau)}\right]^{3/4}\gamma_{mod}^{-1}(a_f^{-1})\,. \tag{3.6}$$

In the beginning of the postcollaptical evolution of the cavern τ_{trap} coincides with $\gamma_{mod}^{-1}(a_f^{-1})$ and the sharp trapping condition (3.5) remains to be valid. Further, $\tau_{trap}(\tau)$ decreases proportionally to $\tau^{-1/2}$ and by the moment τ_f it achieves the value

$$\tau_{trap}(\tau_f)\sim\left(\frac{a_f}{a_0}\right)^{1/4}\gamma_{mod}^{-1}(a_f^{-1})\,. \tag{3.7}$$

The corresponding to (3.7) condition of sharp trapping is slightly softer than (3.5):

$$k_0a_f\ll g^{4/7}. \tag{3.8}$$

Under the conditions, where (3.8) is satisfied and (3.5) — not, the plasmon trapping at the initial stage of postcollaptical evolution of cavern proceeds «slowly» or «adiabatically», while at $\tau\sim\tau_f$ the trapping is sharp.

For the evaluation of a size of the adiabatic trapping region $k_{trap}^{-1}\gg k_0^{-1}$ it is enough to note that the small displacement of corresponding to the virtual or weakly bound state pole of the scattering amplitude of plasmons on cavern in the plane of complex wave number k is proportional to a little deepening of cavern $\delta n(a_f,\tau)$ and, consequently, to the time $\delta\tau$ during which this deepening is occured. The trapping length k_{trap}^{-1} is found from the condition of the given pole displacement by $\delta k\sim k_{trap}$ during the time of flight $\delta\tau\sim\omega^{-1}(k_{trap})$ of a plasmon with a wave number $k\sim k_{trap}$ above the cavern:

$$k_{trap}\sim\delta k\sim k_0\frac{\delta\tau}{\tau_{trap}}\sim\frac{k_0}{\tau_{trap}\,\omega(k_{trap})}\sim\frac{k_0}{\omega_0\,\tau_{trap}}\left(\frac{k_0}{k_{trap}}\right)^2.$$

Whence,

plasmon to be detected in the central part of cavern is inversely proportional to the scale of plasmons localization region (since beyond the cavern the probability density drops as r^{-2}).

4. ABSORPTION OF TRAPPED PLASMONS

An energy of each plasmon is distributed between various spatial scales including also those, which are lower than a_f. It is necessary to know this distribution for evaluating the Landau damping of plasmons. The character of the sought for distribution in the region of scales, which are small compared to the main energy-content one, practically does not depend on the value of the latter. Therefore, for obtaining the most general picture, it is reasonable to consider plasmons with the sufficiently large energy-content scale $r_{\mathfrak{g}}$. Within the region $r_{\mathfrak{g}} \gtrsim r \gtrsim a_*(\tau)$, where the attracting potential has the form (2.2), the probability $P(r, r_{\mathfrak{g}}, \tau)$ for a plasmon to be detected at a distance of order r from the cavern center is defined by the following chain of estimates:

$$P(r, r_{\mathfrak{g}}, \tau) \sim r \frac{\partial P(r, r_{\mathfrak{g}}, \tau)}{\partial r} \propto \frac{r}{k(r, \tau)} \propto r^2 . \qquad (4.1)$$

The behaviour of the function $P(r, r_{\mathfrak{g}}, \tau)$ in the region $a_f \lesssim r \lesssim a_*(\tau)$, where the potential has the form (2.4), depends on the symmetry of caverns. At the presence of axial symmetry, the expansion of the plasmon wave function over the spherical harmonics at a distance r from the cavern center contains (with estimately the same weights) of about $l(r, \tau) \sim k(r, \tau) r$ terms, hence,

$$P(r, r_{\mathfrak{g}}, \tau) \propto l(r, \tau) \frac{r}{k(r, \tau)} \propto r^2, \quad a_f \lesssim r \lesssim a_*(\tau) . \qquad (4.2)$$

In an essentially nonsymmetric cavern the projection of an orbital momentum is not fixed and the number of harmonics, existing (with estimately the same weights) in the expansion of the plasmon wave function at a distance r from the cavern center, is approximately equal to $l^2(r, \tau)$, correspondingly,

$$P(r, r_{\mathfrak{g}}, \tau) \propto l^2(r, \tau) \frac{r}{k(r, \tau)} \propto r^{5/4}, \quad a_f \lesssim r \lesssim a_*(\tau) . \qquad (4.3)$$

As already mentioned above, the probability $P(r, r_{\mathfrak{g}}, \tau)$ has a meaning of an energy portion, contained at distances of order r from the cavern center, for a plasmon, whose main scale of location is equal to $r_{\mathfrak{g}}$. Since each scale r has its corresponding characteristic wave number $k(r, \tau)$ (see (2.5)), the distribution $\bar{P}(k, r_{\mathfrak{g}}, \tau)$ of the plasmon energy in the k-space can be obtained from $P(r, r_{\mathfrak{g}}, \tau)$ by a simple replacement of r by the function $r(k, \tau)$ inverse to $k(r, \tau)$. With such a replacement the relations (4.1) — (4.3) obtain the following form:

$$\bar{P}(k, r_{\mathfrak{g}}, \tau) \propto k^{-2}, \quad r_{\mathfrak{g}}^{-1} \lesssim k \lesssim a_*^{-1}(\tau) ; \qquad (4.4)$$

$$\bar{P}(k, r_{\mathfrak{g}}, \tau) \propto k^{-8/7}, \qquad (4.5)$$

$$a_*^{-1}(\tau) \lesssim k \lesssim k_M(\tau) \equiv k(a_f, \tau) ;$$

$$\bar{P}(k, r_{\mathfrak{g}}, \tau) \propto k^{-5/7}. \qquad (4.6)$$

The estimate (4.5) relates to the axisymmetric and (4.6) — to the essentially nonsymmetric caverns, respectively.

The main contribution into the Landau damping of the trapped plasmon is given by those values of k, at which the product $\gamma_L(k) \bar{P}(k, r_{\mathfrak{g}}, \tau)$ achieves its maximum. This maximum is located in the region $k_f \lesssim k \lesssim k_M(\tau)$ beyond which one of the factors (at $k \ll k_f$ it is $\gamma_L(k)$, and at $k \gg k_M(\tau)$ it is $\bar{P}(k, r_{\mathfrak{g}}, \tau)$) is exponentially small. If the function $\gamma_L(k)$ increases in the region $k_f \lesssim k \lesssim k_M(\tau)$ faster than $\frac{1}{\bar{P}(k, r_{\mathfrak{g}}, \tau)}$, the main contribution into the plasmon damping is given by the wave numbers $k \sim k_M(\tau)$, if slower, then — by $k \sim k_f$. An analysis shows that the second proposal does not allow the self-consistent description of electron acceleration by the Langmuir waves. Accepting the first alternative, one can evaluate the plasmon damping decrement as

$$\Gamma(r_{\mathfrak{g}}, \tau) \sim \gamma_L(k_M(\tau))\, P(a_f, r_{\mathfrak{g}}, \tau) . \qquad (4.7)$$

The absorption of a plasmon sucking by cavern occurs at such a scale of its localization $r_{\mathfrak{g}} = r_{abs}(\tau)$, at which the inverse time $\tau_{suck}^{-1}(r_{\mathfrak{g}}, \tau)$ of decreasing $r_{\mathfrak{g}}$ by a factor of 2 turns out to be of order of $\Gamma(r_{\mathfrak{g}}, \tau)$:

$$\Gamma(r_{abs}, \tau) \sim \tau_{suck}^{-1}(r_{abs}, \tau) . \qquad (4.8)$$

This relation establishes the connection between the radius of plasmon absorption in cavern $r_{abs}(\tau)$ and the distribution function of electrons over velocities (in terms of which the Landau damping decrement $\gamma_L(k)$ is easily expressed (see (1.13)). The sucking time $\tau_{suck}(r_{\mathcal{E}}, \tau)$ can be evaluated in the following way. For this time the plasmon frequency is decresed by

$$\delta\omega \sim \omega_{pe} \frac{|\tilde{n}(r_{\mathcal{E}}, \tau)|}{n_0}. \tag{4.9}$$

In terms of the Schrödinger equation $\delta\omega$ can be considered as an increase of the plasmon «bound energy» in cavern during the time $\tau_{suck}(r_{\mathcal{E}}, \tau)$. From this point of view it is clear that $\delta\omega$ is estimately equal to a decrease in the plasmon «potential energy»:

$$\delta\omega \sim \omega_{pe} \frac{|\delta\tilde{n}(a_f, \tau)|}{n_0} P(a_f, r_{\mathcal{E}}, \tau). \tag{4.10}$$

Here, $|\delta\tilde{n}(a_f, \tau)|$ is the cavern deepening in its central part during the time $\tau_{suck}(r_{\mathcal{E}}, \tau)$. According to (2.4),

$$|\delta\tilde{n}(a_f, \tau)| \sim |\tilde{n}(a_f, 0)|\, \gamma_{mod}(a_f^{-1})\, \tau_{suck}(r_{\mathcal{E}}, \tau). \tag{4.11}$$

The substitution of (4.9) and (4.11) into (4.10) results in the estimate

$$\tau_{suck}^{-1}(r_{\mathcal{E}}, \tau) \sim \gamma_{mod}(k_f) \frac{|\tilde{n}(a_f, 0)|}{|\tilde{n}(r_{\mathcal{E}}, \tau)|} P(a_f, r_{\mathcal{E}}, \tau). \tag{4.12}$$

Taking into account the condition of the collapse ceased at a cavern size a_f:

$$\gamma_{mod}(k_f) \sim \gamma_L(k_f) \tag{4.13}$$

and also the estimates (4.12), (4.7), one can rewrite the relation (4.8) in the form

$$\frac{|\tilde{n}(a_f, 0)|}{|\tilde{n}(r_{abs}, \tau)|} \sim \frac{\gamma_L(k_M(\tau))}{\gamma_L(k_f)}. \tag{4.14}$$

It noticeable, that the cavern symmetry dependent probability is not included in (4.14). In this sense, the given by (4.14) relation of the plasmon absorption radius in cavern with the distribution function of electrons is universal:

$$\tau_{abs}(\tau) \sim \begin{cases} a_f \left[\dfrac{\gamma_L(k_M(\tau))}{\gamma_L(k_f)}\right]^{1/2}, & r_{abs}(\tau) \gtrsim a_*(\tau) \\ a_f [\gamma_L(k_M(\tau))\, \tau]^{2/7}, & r_{abs}(\tau) \lesssim a_*(\tau) \end{cases} \tag{4.15}$$

At $\tau \sim \gamma_{mod}^{-1}(k_f)$ the values $a_*(\tau)$ and $r_{abs}(\tau)$ coincide with a_f, and at $\gamma_{mod}^{-1}(k_f) \lesssim \tau \lesssim \tau_f$, according to (2.3) — (2.5),

$$a_*(\tau) \propto \tau^{2/3}, \quad k_M(\tau) \propto \tau^{1/2}. \tag{4.16}$$

Therefore, the first of the alternatives considered in (4.15) is realized in the case, when the function $\gamma_L(k)$ increases in the range $k_f \lesssim k \lesssim k_M \equiv k_M(\tau_f)$ faster than $k^{8/3}$, and the second one — if slower than $k^{8/3}$.

5. STRONG LANGMUIR TURBULENCE SPECTRUM

In an inertial interval of scales $k_0 \ll k \ll k_f$ an average energy density $W(k)$ for the Langmuir waves with lengths of order k^{-1} is determined by the collapse and has the form

$$W(k) \sim \mathcal{E}_0 N_{cav}(k_0) \frac{\gamma_{mod}(k_0)}{\gamma_{mod}(k)}. \tag{5.1}$$

Within the region $k_f \ll k \ll k_M = k_M(\tau_f) \sim k_f(k_f/k_0)^{1/4}$ the value $W(k)$ is determined by the postcollaptical effects. Let $W(k, \tau)$ be the contribution into $W(k)$ of Langmuir wave localized in cavern with the postcollapse lifetime of order τ, $\bar{N}_{cav}(\tau)$ be the concentration of these caverns, and $\mathcal{E}(k, \tau)$ be an energy of localized in each of the caverns Langmuir waves with lengths of order k^{-1}. Then

$$W(k, \tau) \sim \bar{N}_{cav}(\tau)\, \mathcal{E}(k, \tau). \tag{5.2}$$

For $\bar{N}_{cav}(\tau)$ the following estimate is valid

$$\bar{N}_{cav}(\tau) \sim N_{cav}(k_0)\, \gamma_{mod}(k_0)\, \tau \sim N_{cav}(k_f)\, \gamma_{mod}(k_f)\, \tau. \tag{5.3}$$

An energy $\mathcal{E}(k, \tau)$ estimate depends both on the regime of plasmon trapping and the cavern symmetry. In the regime of sharp trapping of plasmons a noticeable portion of energy flux flying upon the cavern is absorbed, therefore,

$$\gamma_L(k_M(\tau))\, \mathcal{E}(k_M(\tau), \tau) \sim \omega_0 \mathcal{E}_0. \tag{5.4}$$

In the regime of slow trapping of plasmons, each bound state gets an energy of order $\mathcal{E}_0$, therefore,

$$\gamma_L(k_M(\tau))\, \mathcal{E}(k_M(\tau), \tau) \sim \tau_{trap}^{-1}(\tau)\, \mathcal{E}_0. \tag{5.5}$$

The relations (5.2) — (5.5), (1.15) enable one to express the energy density $W(k_M(\tau),\tau)$ via the distribution function of electrons accelerated by the turbulence:

$$W(k_M(\tau),\tau)\sim\mathscr{E}_0 N_{cav}(k_f)\,\frac{\gamma_{mod}(k_f)}{\gamma_L(k_M(\tau))}\,\frac{\omega_0\tau}{1+\omega_0\tau_{trap}(\tau)}. \qquad (5.6)$$

Since all the bound states with a scale of localization $r\lesssim r_{abs}(\tau)$ by the time moment τ are already exhausted by the Landau damping, in the region $k_M(\tau)\gtrsim k\gtrsim k(r_{abs}(\tau),\tau)$:

$$W(k,\tau)\propto\mathscr{E}(k,\tau)\propto\bar{P}(k,r_{abs}(\tau),\tau) \qquad (5.7)$$

and, hence,

$$\frac{W(k,\tau)}{W(k_M(\tau),\tau)}\sim\frac{P(r(k,\tau),r_{abs}(\tau),\tau)}{P(a_f,r_{abs}(\tau),\tau)}\sim P(a_f,r(k,\tau),\tau). \qquad (5.8)$$

At $k(r_{abs}(\tau),\tau)\gg k\gg k_0$ the main contribution into an energy density $W(k,\tau)$ is given by the states with the scale of localization $r(k,\tau)\gg r_{abs}(\tau)$, where plasmons are not damped yet. In this region

$$\tau_{suck}^{-1}(r(k,\tau),\tau)\,\mathscr{E}(k,\tau)\sim\frac{\omega_0\mathscr{E}_0}{1+\omega_0\tau_{trap}(\tau)} \qquad (5.9)$$

and, according to (5.2), (5.3),

$$W(k,\tau)\sim\mathscr{E}_0 N_{cav}(k_f)\,\gamma_{mod}(k_f)\,\tau\,\frac{\omega_0\tau_{suck}(r(k,\tau),\tau)}{1+\omega_0\tau_{trap}(\tau)}. \qquad (5.10)$$

The relations obtained enable one to represent the energy density $W(k)=\max_\tau W(k,\tau)$ $(k\gg k_f)$ and, hence, the coefficient of electron diffusion on Langmuir waves $D(v)\sim\frac{\omega_{pe}}{m_e n_0}W\left(\frac{\omega_{pe}}{v}\right)$ in terms of the electron distribution function and, thereby, to close its evolution equation. An analysis shows that self-consistent solution of this equation for the function $W(k,\tau)$ increasing over τ, i. e. at $W(k)\sim W(k,\tau_f)$, is impossible.

Let the function $W(k,\tau)$ decrease over τ. Then, $W(k)\sim W(k,\tau_m(k))$, where $\tau_m(k)$ is the smallest of those τ-values, at which in the cavern there are plasmons with the wave number given $k\,(k_f\ll k\lesssim k_M)$. The function $\tau_m(k)$ is evidently an inverse with respect to $k=k_M(\tau)$. Therefore, for the evaluation of $W(k)$ it is sufficient to substitute $\tau=\tau_m(k_M(\tau))$ into (5.6) and replace $k_M(\tau)$ by k:

$$W(k)\sim\mathscr{E}_0 N_{cav}(k_f)\,\frac{\gamma_{mod}(k_f)}{\gamma_L(k)}\,\frac{\omega_0\tau_m(k)}{1+\omega_0\tau_{trap}(\tau_m(k))}\qquad (k_f\ll k\lesssim k_M). \qquad (5.11)$$

Whence, in the regime of sharp trapping of plasmons $(\omega_0\tau_{trap}\ll 1)$

$$W(k)\propto\frac{\tau_m(k)}{\gamma_L(k)}\propto\frac{k^2}{\gamma_L(k)}\qquad (k_f\ll k\lesssim k_m). \qquad (5.12)$$

Under the assumption that most of caverns are axisymmetric, the function $\tau_{trap}(\tau)$ does not depend on τ (see (3.4)) and the relation (5.12) remains to be valid at the adiabatic trapping of plasmons. If the typical cavern is esentially nonsymmetric, then $\tau_{trap}(\tau)\propto\tau^{-1/2}$ (see (3.6)) and at adiabatic trapping of plasmons $(\omega_0\tau_{trap}\gg 1)$:

$$W(k)\propto\frac{k^3}{\gamma_L(k)}\qquad (k_f\ll k\lesssim k_M). \qquad (5.13)$$

6. THE SCENARIO OF ELECTRON ACCELERATION

The acceleration of electrons by Langmuir turbulence is described by the quasilinear diffusion equation and is going on much slower than the processes of collapse and further evolution of caverns described above. For the characteristic time $t_d(v)$ of the diffusion of electrons with velocities of order v the following estimate is valid

$$t_d^{-1}(v)\sim\frac{D(v)}{v^2}\sim\frac{\omega_{pe}}{m_e n_0 v^2}\,W\left(\frac{\omega_{pe}}{v}\right). \qquad (6.1)$$

Within the applicability region (1.11) for previous models of Langmuir turbulence

$$\gamma_L(k)\propto k^{3/2},\qquad k\gg k_f. \qquad (6.2)$$

(see (1.15), (1.16)). The plasmon trapping by the caverns is sharp for the given range of parameters and, in virtue of (5.12),

$$W(k)\propto k^{1/2},\qquad k_f\ll k\lesssim k_M. \qquad (6.3)$$

The typical time for electron diffusion $t_d(v)$ increases within the interval $v_m\lesssim v\ll v_f$ $\left(v_f=\frac{\omega_{pe}}{k_f},\ v_m=\frac{\omega_{pe}}{k_M}\sim v_f(v_f/v_0)^{1/4}\right)$ as $v^{5/2}$:

$$\frac{t_d(v_m)}{t_d(v)}\sim\left(\frac{v_m}{v}\right)^{5/2}\sim\left(\frac{v_f}{v_0}\right)^{5/8}\left(\frac{v_f}{v}\right)^{5/2},\qquad v_m\lesssim v\ll v_f. \qquad (6.4)$$

In the vicinity of a wave number $k \sim k_f$ the turbulence energy density $W(k)$ decreases sharply (exponentially) over k from the value of $W(k_f) \sim N_{cav}(k_f)\mathscr{E}_0$, determined by the collapse, down to the value of energy density of postcollaptically absorbed waves, calculated by the formula (5.11) with $k = \Lambda k_f (\Lambda = 2 \div 3)$:

$$W(\Lambda k_f) \sim N_{cav}(k_f)\, \mathscr{E}_0 \frac{\omega_0}{\gamma_{mod}(k_f)}.$$

An inverse time of electron diffusion $t_d^{-1}(v)$ is decreasing by the same factor at $v \sim v_f$

$$\frac{t_d(v_f)}{t_d(\Lambda v_f)} \sim \frac{\omega_0}{\gamma_{mod}(k_f)} \sim g^{-1}\left(\frac{v_f}{v_0}\right)^{3/2}. \qquad (6.5)$$

Electrons have not enough time for the diffusion on the postcollaptically absorbed waves under the condition $t_d(v_m) \gg t_d(v_f)$, which (via (6.4) and (6.5)) is reduced to the right-hand side inequality in (1.11). When the velocity $v_f(t)$, increasing according to (1.18), attains the value $v_{f1} \sim g^{8/7} v_0$, in the vicinity of the velocity $v_{m1} \sim g^{10/7} v_0$ the second front of electron acceleration is formed. Its position $\tilde{v}_f(t)$ at $v_f(t) \gtrsim v_{f1}$ is found from the condition $t_d(\tilde{v}_f) \sim t_d(v_f)$ and is given by the estimate

$$\tilde{v}_f \sim g^{-2/5}\left(\frac{v_f}{v_0}\right)^{3/5} v_f. \qquad (6.6)$$

Within the region $v_f \gg v \gg \tilde{v}_f$ the diffusion is proceeding slowly and electron concentration $n_e(v)$ remains to be the same as it became after the first acceleration (see (1.16)). In the region $\tilde{v}_f \gtrsim v \gtrsim v_m$ the time of electron diffusion should approximately be the same as $t_d(v_f)$, otherwise, it is impossible to construct the self-consistent solution. The condition $t_d(v) \sim t_d(v_f)$ of the function (6.1) independence on the velocity means that

$$W(k) \propto k^{-2}, \qquad k_M \gtrsim k \gtrsim \tilde{k}_f = \frac{\omega_{pe}}{\tilde{v}_f}. \qquad (6.7)$$

According to (5.12), in order to obtain the given wave spectrum, the following condition should be satisfied $\gamma_L(k) \propto k^4$, i. e.

$$n_e(v) \propto v^{-4}, \qquad v_m \lesssim v \lesssim \tilde{v}_f. \qquad (6.8)$$

In the region $v \ll v_m$ there is no diffusion, i. e. the concentration of

electrons remains invariable after the rear front of acceleration $v_m(t)$ passed through the velocity v. At $v_m(t) \gtrsim v_{m1}$:

$$n_e(v_m) \sim n_e(\tilde{v}_f)\left(\frac{\tilde{v}_f}{v_m}\right)^4 \propto \tilde{v}_f^{5/2} v_m^{-4} \propto v_m^{-4/5}$$

with the time-independent propotion factor, hence

$$n_e(v) \propto v^{-4/5}, \qquad v_m \gtrsim v \gtrsim v_{m1}. \qquad (6.9)$$

The law for motion of the first front of electrons acceleration $v_f(t)$ is found out from the diffusion equation (1.17). While, $v_f \lesssim g v_0$, each cavern absorbs an energy of order $\mathscr{E}_0$, therefore, at a constant energy flux Π coming into a strong Langmuir turbulence, the cavern concentration $N_{cav}(k_0)$ is time-independent and the relation $W(k_f) \propto k_f^{-3/2} \propto v_f^{3/2}$ and the law (1.18) are still valid. At $v_f \gtrsim g v_0$ the postcollaptical absorption of energy dominates which value is proportional to v_f in the regime of the plasmons sharp capture (see (1.16)); the cavern concentration $N_{cav}(k_0)$ at an energy flux Π given decreases with time as v_f^{-1}, and the diffusion coefficient $D(v_f)$ grows as $v_f^{1/2}$:

$$D(v_f) \propto W(k_f) \propto N_{cav}(k_f) \propto k_f^{-3/2} N_{cav}(k_0) \propto k_f^{-1/2} \propto v_f^{1/2}.$$

In this case, from (1.17) it follows that

$$v_f(t) \propto t^{2/3}, \qquad v_f \gtrsim g v_0. \qquad (6.10)$$

The picture described above remains correct unless the second front of electrons acceleration $\tilde{v}_f(t)$ catches the first one, which occurs at $v_f(t) \sim v_{f2}$:

$$v_{f2} \sim g^{2/3} v_0 \qquad (v_{m2} \sim g^{5/6} v_0). \qquad (6.11)$$

By this moment, an increment of modulational instability at the final stage of a collapse $\gamma_{mod}(k_f)$ is decreased down to the value ω_0 and the plasmons capture, at least, at the beginning of the cavern postcollaptical evolution ceases to be sharp. The further scenario of electron acceleration depends on the dominant symmetry of caverns. At $v_f(t) \gtrsim v_{f2}$ the diffusion of electrons proceeds equally fast over the entire region $v_m \lesssim v \lesssim v_f$, which means that

$$W(k) \propto k^{-2}, \qquad k_f \lesssim k \lesssim k_M. \qquad (6.12)$$

In the case of the axisymmetric caverns the relation (5.12) is still valid for the adiabatic capture of plasmons, hence, in the region of electron diffusion the following distribution is established

$$n_e(v) \propto v^{-4}, \quad v_m \leqslant v \leqslant v_f. \tag{6.13}$$

In the region $v \leqslant v_m$ there is no diffusion, i. e. the concentration of electrons does not vary after the rear front of acceleration $v_m(t)$ passed through the velocity v:

$$n_e(v) \propto v^{-2}, \quad v_m \leqslant v \leqslant v_{m2}. \tag{6.14}$$

At this stage of evolution the law for the electron acceleration front motion $v_f(t)$, (defined by the diffusion equation (1.17)) turns out to be exponential. Indeed, at adiabatic capture of plasmons an energy absorbed by one axially symmetric cavern is of the order

$$\mathscr{E}_0 N_0(a_f, \tau_f) \propto \left[\frac{a_*(\tau_f)}{a_f}\right]^{3/2} \propto a_f^{-1/2} \propto v_f^{-1/2}.$$

Therefore, at a constant flux of energy Π coming into the strong turbulence, the concentration of caverns $N_{cav}(k_0)$ increases with time as $v_f^{1/2}$ and the diffusion factor $D(v_f)$ — as v_f^2. From the equation $\frac{d}{dt} v_f^2 \propto v_f^2$ it follows that

$$\ln v_f(t) \propto t, \quad v_f(t) \gtrsim v_{f2}. \tag{6.15}$$

This picture is valid until the inertial interval is vanished, if v_0 is less than the light velocity c. If $v_0 \gg c$, the inertial interval cannot vanish and v_f achieves (with time) the relativistic values, for which the diffusion equation in the form (1.17) is not applicable. An analysis revealed that in this case the system comes to the regime of relativistic electrons storage. Since the detailed description of this regime is beyond the scope of present article, further it is assumed that $v_0 < c$.

In the case of essentially nonsymmetric caverns the scenario of electron acceleration turns out to be somewhat more complex due to the existence of a relatively narrow transitive region

$$v_{f2} \leqslant v_f(t) \leqslant v_{f3} = g^{4/7} v_0 \quad (v_{m2} \leqslant v_m(t) \leqslant v_{m3} = g^{5/7} v_0), \tag{6.16}$$

where the capture of plasmons at the initial stage of postcollaptical evolution of a cavern is adiabatic and at the final stage it is sharp. The regime of plasmons capture changes at $\omega_0 \tau_{trap}(\tau) \sim 1$, which means that (taking into account (3.6), (2.5)) $k_M(\tau) \sim \tilde{k}$:

$$\tilde{k} = k_f \frac{\omega_0}{\gamma_{mod}(k_f)} \sim g^{-1} k_0 \left(\frac{k_0}{k_f}\right)^{1/2}. \tag{6.17}$$

In the region $k_f \leqslant k \leqslant \tilde{k}$ the main contribution into an energy density is given by the adiabatically captured plasmons. By virtue of (6.12), (5.12), (5.13), (1.15), it follows that

$$n_e(v) \propto v^{-5}, \quad v_f \gtrsim v \gtrsim \tilde{v} = \frac{\omega_{pe}}{\tilde{k}};$$

$$n_e(v) \propto v^{-4}, \quad \tilde{v} \gtrsim v \gtrsim v_m. \tag{6.18}$$

With an increase in the velocity $v_f(t)$ from v_{f2} up to v_{f3}, the velocity

$$\tilde{v} = \frac{\omega_{pe}}{\tilde{k}} \sim g v_0 \left(\frac{v_0}{v_f}\right)^{1/2} \tag{6.19}$$

decreases from v_{f2} down to v_{m3}, in other words, the electron distribution reconstruction wave passes along the region of diffusion in the direction from $v_f(t)$ to $v_m(t)$. In the region $v \ll v_m(t)$ the concentration of electrons is time-independent and remains the same as it was at $v_m(t) \sim v$:

$$n_e(v) \propto v^{-4/5}, \quad v_m \gtrsim v \gtrsim v_{m2}. \tag{6.20}$$

A major part of energy is absorbed by caverns at the final stage of their postcollaptical evolution via the sharp capture of plasmons, therefore, the dependence of caverns concentration on $v_f(t)$ remains the same as previous and $v_f(t)$ continues to grow by the law (6.10).

Upon reaching the value v_{f3} by the velocity $v_f(t)$ the capture of plasmons becomes adiabatic even at the final stage of the caverns postcollaptical evolution and for the entire region of diffusic the first distribution from (6.18) is established. At $v_f(t) \gtrsim v_{f3}$ the distribution of electrons has the form:

$$n_e(v) \propto v^{-5}, \quad v_f \gtrsim v \gtrsim v_m,$$

$$n_e(v) \propto v^{-11/5}, \quad v_m \gtrsim v \gtrsim v_{m3}. \tag{6.21}$$

An energy absorbed by a single cavern in this regime is estimately equal to

$$\mathscr{E}_0 N(a_f, \tau_f) \propto \left[\frac{a_*(\tau_f)}{a_f}\right]^{9/4} \propto a_f^{-3/4} \propto v_f^{-3/4}.$$

Therefore, at a constant power supply Π to a strong turbulence the concentration of caverns $N_{cav}(k_0)$ grows with time proportionally to $v_f^{3/4}$ and the diffusion factor $D(v_f)$ — proportionally to $v_f^{9/4}$. At such a dependence $D(v_f)$, from the diffusion equation (1.17) the explosive growth of $v_f(t)$ follows:

$$v_f(t) \propto (t_0 - t)^{-4}, \quad v_f \gtrsim v_{f3}. \tag{6.22}$$

This regime is kept until the inertial interval is vanished.

7. CONCLUSION

In searching the wave and particle spectra described above a number of assumptions were used that were formulated explicitely or assumed. The result obtained justifies all asumptions taken, in particular, the cavern postcollaptical deepening turns out to be really inertial, i. e. practically not affected by the sucked Langmuir waves, and modulational instability of these waves has not enough time to be developed during their presence in the cavern. Quite numerous (and therefore, not considered in detail) alternative variants of assumptions were also analyzed, in fact, but no other self-consistent solutions were found.

As is seen from the content of this paper, the postcollaptical effects can be neglected only at an early stage of the accelerated electrons «tail» formation. Nearly all the energy contribution into a strong Langmuir turbulence occured at the later (and longer) stages, when the postcollaptical effects dominate and change qualitatively the wave and particle spectra. The most remarkable difference from the predictions of previous models of a strong Langmuir turbulence is in the dependence of the form (and even total energy) of accelerated electrons «tail» on the symmetry of majority of caverns. The versions of electron «tail» corresponding to axisymmetric and essentially nonsymmetric caverns, which is formed by the moment of the inertial interval vanishing, are given by the estimates (see §6):

$$n_e(v) \sim n_0 \frac{c_i}{v_0} \left(\frac{v_0}{v}\right)^2, \quad v_0 \gtrsim v \gtrsim g^{5/6} v_0; \tag{7.1}$$

$$n_e(v) \sim n_0 \frac{c_i}{v_0} \left(\frac{v_0}{v}\right)^{11/5}, \quad v_0 \gtrsim v \gtrsim g^{5/7} v_0. \tag{7.2}$$

In the region $g^{5/6}v_0 \gtrsim v \gtrsim g^{10/7}v_0$ (for axially symmetric) or $g^{5/7}v_0 \gtrsim v \gtrsim g^{10/7}v_0$ (for nonsymmetric collapse) the final concentration of electrons with velocities of the order of v has the form

$$n_e(v) \sim n_0 \frac{c_i}{v_0} g^{-1} \left(\frac{v_0}{v}\right)^{4/5}. \tag{7.4}$$

In the region $g^{10/7}v_0 \gtrsim v \gtrsim v_0 (k_0 r_D)^{5/6}$ (existing under condition (1.10)) the prediction of the previous models of a strong turbulence is still valid:

$$n_e(v) \sim n_0 \frac{c_i}{v_0} \left(\frac{v_0}{v}\right)^{3/2}. \tag{7.4}$$

The total energy density of electron «tail» in the case of axisymmetric collapse is mainly determined by the region $v_0 \gtrsim v \gtrsim g^{5/6}v_0$ and estimately is equal to:

$$W^e \sim n_0 m_e c_i v_0 \ln g^{-1}. \tag{7.5}$$

In the case of nonsymmetric collapse the major part of the «tail» energy is concentrated in electrons with velocities of the order of $g^{5/7}v_0$ and is estimately equal to:

$$W^e \sim n_0 m_e c_i v_0 g^{1/7}. \tag{7.6}$$

In fact, the difference in values of (7.5) and (7.6) is minor, but in principle, there is an interesting possibility to discuss the Langmuir collapse symmetry using only on a macroscopic parameters of system. It is also quite difficult to distinguish the distribution (7.1) from (7.2) experimentally; much simpler is to distinguish them from the distribution (7.4) predicted earlier and, especially, from (7.3). The detection of characteristic break of the accelerated electrons distribution function between section (7.1) (or (7.2)) and (7.3) is, probably, the most accessible way of experimental verification of the presented here theory.

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