TEMPORAL EVOLUTION OF WAVES AND INSTABILITIES IN PLASMA WITH SHEAR FLOW

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1. Introduction

The experimental discovery of the transition from the low confinement to the high confinement state, in which the suppression of turbulence and reduction of anomalous transport was detected, open a new page in the theory of drift wave and drift turbulence in toroidal confinement systems. Experiments have shown [1] that, together with the transition to the improved confinement state (H–mode regime), a tokamak plasma develops large variations in the radial electric field, and strong poloidal plasma shear flows. The large variations are also limited to the same small edge layer in which the velocity shear length is found to be much less than the magnetic shear length. In fact, the shearing rate in this region is of the order of or even larger than the typical drift wave frequency. It stands to reason, then, that the nature and evolution of the low frequency waves and instabilities in the edge layer will be different from that in the plasma core, where the magnetic shear is the primary determinant of the spatial structure and temporal evolution of these waves. Unfortunately, the "ballooning transform" method ceases to be useful for problems that involve significant shear flow, and may be suitable only in finding the spectrum (and growth rates) in the limit of vanishing velocity shear. Even for shearless magnetic field rigorous treatment of the shear flow effects encounters a fatal difficulty in the case of the finite flow shear arising from the non-Hermitian (non-self-adjoint) properties of the problem. The standard mode approach breaks down, and the theory may fail to give correct prediction of evolution even if the perturbation fields remain in the linear regime. Therefore other methods have to be developed for the analysis of plasmas with strong flow shear. In this report we present results of application of new approach to the linear analysis of low frequency turbulence of plasma with shear flow. The drift waves and instabilities (considered on the base of Hasegawa–Mima equation and and Hasegawa–Wakatani system), Eta–i instability (considered on the base of Nordman–Weiland system), Rayleigh–Taylor instability, Alfven waves and instabilities of plasmas with shear flow are studied as an initial value problem without the use of spectral expansion in time. It is shown that the conventional modal structure of these waves and instabilities pertains only in the initial stage of their evolution. For larger times, numerous previously overlooked non-modal effects due to the velocity shear determined their temporal evolution. The solutions for electrostatic and magnetic potentials, ion density and temperature perturbation are obtained for any desired finite times of the temporal evolution of the above mentioned waves and instabilities in plasma with homogeneous shear flow. In general these results cannot be obtained practically with normal mode approach or on the base of the Laplace transform in time. The assessment of the importance of these solutions for real physical process in comparison with nonlinear enhanced decorrelation effect is presented.

1. Drift waves in a collisional plasma with homogeneous shear flow.

On the base of the slab model of Hasegama–Wakatani we consider the temporal evolution of the resistive drift wave perurbations in plasmas with homogeneous shear flow. This model is defined by the following system of equations for the dimensionless perturbations of electron density n, $n = \widetilde{n}/n_0$, and potential $\phi = e\varphi/T_e$ (T_e is the electron temperature, n_0 is the electron background density) [2]

$$\rho_s^2 \left(\frac{\partial}{\partial t} + \mathbf{V}\cdot\nabla\right)\triangle\phi = a\frac{\partial^2}{\partial z^2}\left(n-\phi\right) \tag{1}$$

$$\left(\frac{\partial}{\partial t} + \mathbf{V}\cdot\nabla\right)n + v_{de}\frac{\partial\phi}{\partial y} = a\frac{\partial^2}{\partial z^2}\left(n-\phi\right)\,, \tag{2}$$

where

$$\mathbf{V} = \mathbf{v_0}\left(x\right) + \frac{c}{B^2}\left[\mathbf{B}\times\nabla\varphi\right] = v_0' x\mathbf{e}_y + \frac{c}{B^2}\left[\mathbf{B}\times\nabla\varphi\right]\,, \tag{3}$$

$a = T_e/n_0 e^2 \eta_\parallel$, $\eta_\parallel$ is the resistivity along the magnetic field $\mathbf{B} \parallel \mathbf{z}$, ρ_s^2 is the ion Larmor radius at the electron temperature, v_{de} is the electron diamagnetic velocity.

The analysis of the case with homogeneous shear flow, $v_0' = \text{const}$, is greatly facilitated by a transformation to the coordinates convected with the sheared flow. This transformation is defined by the relations

$$\tau = t, \xi = x, \eta = y - v_0\left(x\right)t = y - v_0' t x, z = z\,. \tag{4}$$

In the new spatial coordinates the linearized system (1), (2) is reduced to the equation for the potential $\phi\left(\tau, \xi, \eta, k_z\right)$,

$$\rho_s^2\frac{\partial}{\partial\tau}\left(\Delta_{\perp c}\phi\right) - \frac{\partial\phi}{\partial\tau} - v_{de}\frac{\partial\phi}{\partial\eta} + \frac{\rho_s^2}{a k_z^2}\frac{\partial^2}{\partial\tau^2}\left(\Delta_{\perp c}\right) = 0, \tag{5}$$

where

$$\Delta_{\perp c} = \frac{\partial^2}{\partial\eta^2} + \left(\frac{\partial}{\partial\tau} - v_0'\tau\frac{\partial}{\partial\eta}\right)^2 \tag{6}$$

Note that the equation (5) does not contain spatially depended coefficients. Fourier transformed equation (5) with wave numbers $k_\perp$ and l conjugated with spatial variables ξ and η, respectively, is found from (5) to be

$$\begin{aligned}&\frac{1}{C}\frac{\partial^2}{\partial T^2}\left[\left(1+T^2\right)\phi\right]\\&+\frac{\partial}{\partial T}\left\{\left[1+l^2\rho_s^2\left(1+T^2\right)\right]\phi\right\}+iS\phi=0,\end{aligned} \tag{7}$$

where a dimensionless time variable T and parameters S and C are defined by

$$T=v_0'\tau-\frac{k_\perp}{l},\ S=\frac{lv_{de}}{v_0'},\ C=\frac{ak_z^2}{\rho_s^2l^2v_0'} \tag{8}$$

In the collisionless limit ($C=\infty$) the electron response is adiabatic ($n=\phi$) and Eqs. (1), (2) are reduced to the Hasegawa–Mima equation, which in considered case of the homogeneous shear flow has a form

$$\frac{\partial}{\partial T}\left\{\left[1+l^2\rho_s^2\left(1+T^2\right)\right]\phi\right\}+iS\phi=0\,. \tag{9}$$

The solution of the initial value problem for Eq.(9) is

$$\begin{aligned}&\phi(\tau,\xi,\eta,k_z)=\frac{1}{(2\pi)^2}\int\int dk_\perp dl\phi(0,\xi,\eta,z)\\&\times\frac{1+\rho_s^2(l^2+k_\perp^2)}{1+\rho_s^2l^2+\rho_s^2\left(v_0'\tau-k_\perp\right)^2}\\&\exp\left[-i\frac{S}{\sqrt{1+\rho_s^2l^2}}\left(\tan^{-1}\frac{\rho_s\left(v_0'\tau-k_\perp\right)}{\sqrt{1+\rho_s^2l^2}}\right.\right.\\&\left.\left.+\tan^{-1}\frac{k_\perp\rho_s}{\sqrt{1+\rho_s^2l^2}}\right)\right]\end{aligned} \tag{10}$$

For small τ $\left(\tau\ll(v_0')^{-1}\right)$ solution (10) gives the trivial result $\phi(\tau,\xi,\eta,k_z)=\phi(0,\xi,\eta,k_z)\exp -i\omega_{de}\tau$ with $\omega_{de}=lv_{de}\left[1+\rho_s^2(l^2+k_\perp^2)\right]^{-1}$. As τ becomes large we find from equation (10) zero frequency convective cell–type solution with decaying as τ^{-2} amplitude. Therefore normal mode solution of the form $\phi=\phi\left(x\right)\exp\left(iky+ik_zz-i\omega t\right)$ is not the steady–state limit for the initial value problem considered.

The evolution in time of the initial disturbance of the wave packet form shows the existence of the stagnation level in both ξ and η directions. Flow shear leads to the vanishing of the group velocity components with time.

In the case of the semiinfinite homogeneous shear flow, which occupies the region $x\geq0$ with boundary conditions $\phi\left(\tau,x=0,y\right)=\widetilde{\phi}\left(\tau,y\right);\ \left.\left(\partial\phi/\partial x\right)\right|_{x=0}=\widetilde{\phi}'\left(\tau,y\right)$ the forced oscillations defined by boundary condition have temporal evolution different from determined by equation(10) [9]. In that case we receive the following equation for the perturbed potential $\phi\left(\tau,k,l\right)$:

$$\begin{aligned}&\left[1+\rho_s^2\left(l^2+\left(k-v_0'l\tau\right)^2\right)\right]\frac{\partial\phi}{\partial\tau}\\&-2\rho_s^2\left(k-v_0'l\tau\right)v_0'l\phi\left(\tau,k,l\right)+\\&ilv_{de}\phi\left(\tau,k,l\right)=\Phi\left(\tau,l\right)\,,\end{aligned} \tag{11}$$

where the term $\Phi\left(\tau,l\right)$ is defined by the boundary conditions at $x=0$ and is equal to

$$\begin{aligned}\Phi\left(\tau,l\right)=&-\rho_s^2\frac{\partial}{\partial\tau}\left\{\frac{\partial}{\partial x}\widetilde{\phi}\left(\tau,x=0,l\right)\right.\\&\left.-i\left(k+2v_0'l\tau\right)\widetilde{\phi}\left(\tau,x=0,l\right)\right\}.\end{aligned} \tag{12}$$

In the case of sufficiently strong flow shear , which is observed on the plasma edge, where $v_0'\gtrsim lv_{de}$, even for sufficiently small times τ that solution may be approximately presented in the form

$$\begin{aligned}&\phi\left(\tau,k,l\right)\approx\phi_0\left(0,k,l\right)\frac{1+\rho_s^2\left(k^2+l^2\right)}{1+\rho_s^2l^2+\rho_s^2\left(k-v_0'l\tau\right)^2}\\&\times\exp\left[-i\frac{lv_{de}}{v_0'l\rho_s\sqrt{1+l^2\rho_s^2}}\left(\arctan\frac{\rho_s\left(v_0'l\tau-k\right)}{\sqrt{1+l^2\rho_s^2}}\right)\right]\\&-\frac{\rho_s^2}{1+\rho_s^2l^2+\rho_s^2\left(k-v_0'l\tau\right)^2}\left[\frac{\partial}{\partial x}\phi\left(\tau,x=0,l\right)\right.\\&\left.-i\left(k+2v_0'l\tau\right)\phi\left(\tau,x=0,l\right)\right].\end{aligned} \tag{13}$$

It is followed from this solution that in plasma with strong velocity shear the initial perturbations in shear flow decays as t^{-2}. At the same time the oscillations of plasma , which are excited out of the shear flow and determined by the boundary conditions, decays as t^{-1} in the region $x\geq0$. Therefore the oscillations observed in shear flow may have as their source the inner parts of plasma.

Much more complicated solutions has the equation (7). The asymptotical forms of these solutions were obtained in [2] in two extremes: $C\gg1$ and $C\ll1$. In both cases solutions display strongly non–modal forms with complicate time dependencies. It was obtained that in the case $C\gg S^2l\rho_s\gg1$ the non-modal effects of the potential evolution becomes essential at the time moments considerably earlier than the inverse growth rate time and normal–mode approach in this case is not applicable. Blocking the wave packets at their stagnation levels and the ultimate transformation of drift waves into convective cells are the intrinsic properties of this evolution. In the case of weak shear, when $S^2l\rho_s\gg C\gg1$ non-modal effects of the amplitude evolution will appear after the long period of the exponential growth of the drift wave amplitude, when nonlinear effects may become more essential. In the limit of $C\ll1$ the solution of the equation (7) in times $T\gg(l\rho_s)^{-1}$ is of the power–like form, $\phi(T)=A_1T^{\alpha_1}+A_2T^{\alpha_2}$, and in the times $T\gg Sl\rho_s$, $S\gg1$ it decays as T^{-2}. In the laboratory frame the frequency and wave numbers are time dependent, $k_x(\tau) = k_\perp - lv_0'\tau, k_y = l, k_z = k_z$. These non-modal time dependencies of the frequency and wave numbers might lead to the broad shape of the spectral line of the drift waves–type perturbations that is not associated with strong turbulence effects.

Toroidal ion temperature gradient driven instability in plasma with shear flow.

The basic linear dynamics of the toroidal η_i instability is described by the system of equations for the electrostatic potential ϕ and ion temperature δT_i perturbations [3]. In the variables, determined by the relations (4) these equations are reduced to the following single equation for the potential ϕ [4]:

$$\begin{aligned}
&\left(1+\rho_s^2 l^2\left(1+T^2\right)\right)\frac{\partial^2\phi}{\partial T^2}+\Bigg\{4\rho_s^2 l^2 T+iS\\
&\times\left[1-\frac{\epsilon_n}{\tau}\left(\frac{10}{3}+\tau\right)+\frac{2}{3}\frac{\epsilon_n}{\tau}\rho_s^2 l^2\left(1+T^2\right)\right.\\
&-\rho_s^2 l^2\left(1+T^2\right)\left(\frac{\epsilon_n}{\tau}\left(\frac{7}{3}-\frac{2}{3}\rho_s^2 l^2\left(1+T^2\right)\right)\right.\\
&\left.\left.+\frac{1+\eta_i}{\tau}\right)\right]\Bigg\}+\Bigg\{2\rho_s^2 l^2-S^2\frac{\epsilon_n}{\tau}\Bigg[\eta_i-\frac{7}{3}\\
&+\frac{5}{3}\frac{(1+\eta_i)}{\tau}\rho_s^2 l^2\left(1+T^2\right)-\frac{2}{3}\frac{(1+\eta_i)}{\tau}\rho_s^4 l^4\\
&\times\left(1+T^2\right)^2\Bigg]-2iS\rho_s^2 l^2 T\left[\frac{\epsilon_n}{\tau}\left(\frac{7}{3}+\frac{1+\eta_i}{\tau}\right.\right.\\
&\left.\left.-\frac{2}{3}\rho_s^2 l^2\left(1+T^2\right)\right)\right]-\frac{5}{3}\frac{\epsilon_n^2}{\tau^2}(1+\tau)S^2\Bigg\}\phi(T)=0,
\end{aligned}\tag{14}$$

where dimensionless time $T = v_0't - k/l$ and dimensionless parameter $S = lv_d/v_0'$ are introduced. There is a significant amount of experimental data (see for example [1]) showing that v_0' is comparable to the growth rate $\gamma_{\max}$ prior the formation of the transport barrier and significantly exceeds it after formation. Because the growth rate γ is less than the drift frequency lv_d the L to H transition occurs in the conditions of the weak flow shear, when $S \gg 1$. It was obtained in ref.[4] that the linear stabilization of the toroidal η_i mode by shear flow in the case $\gamma \sim v_0'$ may be dominant only for the narrow part of the poloidal wave number spectrum. The main part of the η_i -mode spectrum will be stabilized due to the nonlinear deccorrelation effect, for which it is sufficient to have $v_0' \sim \gamma$. After the formation of the transport barrier shearing rate significantly exceeds across the whole minor radius the maximum linear growth of all unstable modes in the plasma [1]. Shearing rate is extremely strong at the plasma edge where it exceeds even the drift wave frequency [1]. In this case $S \ll 1$ and because of the condition $v_0' \gg \gamma$ non-modal effects will define the temporal evolution of any plasma perturbations in the edge layer and therefore the stability of the edge layer, in times less than the inverse growth rate of linear instabilities without velocity shear. The solutions of the equation (14) obtained for the asymptotical limit of small $S \ll 1$ in ref.[4] show that in edge layer with strong flow shear any disturbances are stable against the toroidal η_i instability. From these solutions it follows, that a wave packet will be stagnated in the strong flow shear region or will has a group velocity growing with time and will direct into the region opposite to the direction of the gradient of flow velocity. Therefore the region of strong flow shear with $dE_0(x)/dx < 0$ is a barrier for the outward radial transport of the perturbations of the potential ϕ and ion temperature in the case when this region joins to inner part of plasma. It is just the case which was observed in edge layer of D III–D tokamak [1]

Temporal evolution of drift–Alfvén instabilities

It is well known that in shearless case inhomogeneous plasma with hot ions is unstable against the hydrodynamic and resistive drift–Alfven instabilities [5] . The governing equations for these instabilities are the equations for the longitudinal motion of electrons, for the quasineutrality of the current density, and for the perturbations of the electron pressure $\tilde{p}_e$ and ion pressure $\tilde{p}_i$. In the drift approximation, the system of equations reduces to the system for the parallel component of the perturbed magnetic potential $\tilde{A}_\parallel$, the electrostatic potential $\tilde{\phi}$, and for $\tilde{p}_e$ and $\tilde{p}_i$[6]. In ref.[7] the effect of flow shear on the temporal evolution of these instabilities in the regime of weak flow shear, which corresponds to the stage of the period of the low-to-high (L–H) transition, and in the regime of strong flow shear, which corresponds to the stage of the developed transport barriers was considered separately.

It was obtained in [7] that in low pressure plasma with $\beta \ll m_e/m_e$ and cold ions the flow shear has fundamentally altered the modal time behaviour from the conventional to a power–like time dependence of the spatial Fourier harmonic of the perturbed potentials:

$$\phi \approx \frac{1}{T^{3/2}}\left(C_1 T^{i\omega_1} + C_2 T^{-i\omega_1}\right), \tag{15}$$

$$A_\parallel(T) \approx \frac{ic}{k_\parallel v_A^2}\frac{1}{T^{5/2}}\left(\frac{C_1}{\frac{1}{2}+i\omega_1}T^{i\omega_1} + \frac{C_2}{\frac{1}{2}-i\omega_1}T^{-i\omega_1}\right) \tag{16}$$

An algebraic decay is imposed with the magnetic potential decaying more rapidly with time than the electrostatic potential. This feature – the different time dependence of different perturbations in a linear system, is a strictly non–modal element introduced by the velocity shear.

It was found in [7] that for a plasma with finite pressure ($m_e/m_i \ll \beta \ll 1$) and cold ions Alfvén waves in their "classical″ modal form (with thermal corrections),

$$\phi_{1,2}(\tau) \approx \exp\left\{\pm ik_{\|} v_A \tau \left(1+\left(k_\perp^2+l^2\right)\rho_s^2\right)^{\frac{1}{2}}\right\}, \quad (17)$$

will exist only within the time interval $0<\tau<(v_0')^{-1}$ ($|T|<1$), i.e., for physical times less than the inverse of the shearing rate. In the next well–defined time interval

$$1 \ll T \ll (l\rho_s)^{-1}, \quad (18)$$

the Alfvén waves begin to acquire the non–modal slow–decay imposed on a modal form,

$$\phi_{1,2} \sim \frac{1}{T}\exp(\pm iST). \quad (19)$$

This mixture of the modal and non–modal behaviour continues in the interval

$$(l\rho_s)^{-1} \ll T \ll (l\rho_s)^{-1}\frac{v_{Te}}{v_A} \quad (20)$$

in which the solution

$$\phi_{1,2} \sim T^{-\frac{3}{2}}\exp\left(\pm i\frac{S}{2}l\rho_s T^2\right) \quad (21)$$

is characterized by a "frequency" increasing linearly with time. Finally, for asymptotic times (for times larger than the various "times" in the problem),

$$T \gg (l\rho_s)^{-1}\frac{v_{Te}}{v_A}, \quad (22)$$

the solution

$$\phi_{1,2} \sim T^{-2}\exp\left(\pm iST\frac{v_{Te}}{v_A}\right). \quad (23)$$

corresponds to an algebraically decaying but oscillating mode. In the ultimate stage, the wave that began as an ordinary Alfvén wave ends up propagating with the electron thermal speed v_{Te}. This metamorphosis is due to the finite electron mass effect which is usually omitted in the study of Alfvén waves in a plasma with finite pressure, $\beta > m_e/m_i$. Through the agency of the shear which transfers energy between different wave numbers (in particular induces upward cascading), this normally neglected term becomes dominant for large enough times when the effective wave numbers become large. Since the waves propagating with the the electron thermal speed will be subject to strong electron Landau damping, the flow shear may become a rather effective mechanism for damping the Alfven or Drift–Alfven waves in the shear–layer of the high-confinement tokamak discharges; the shear connects the wave (which may be growing) to the highly damped part of the spectrum causing it to ultimately decay. Needless to say that for these conditions a kinetic description for the Alfvén wave evolution becomes necessary.

In plasma with hot ions hydrodynamic drift–Alfvén and resistive drift–Alfvén instabilities may be developed. For the regimes of low flow shear, which corresponds to the period of the low–to–high transition, the non–modal effects may actually control the temporal evolution of instabilities on time scales less than the their inverse growth rates. However the long time evolution of these instabilities as well as their saturation are determined by the nonlinear effects such as the nonlinear decorrelation effect. In contrast, the plasma with strong flow shear, which correspond to the regime of the developed transport barriers [1], is stable against the development hydrodynamic drift–Alfvén and resistive drift–Alfvén instabilities. For plasma with hot ions ($T_i \leq T_e$) , the frequency of the electron drift wave ultimately reduces to the frequency of the ion drift wave, lv_{di}. The drift–wave transformation into a convective cell with zero frequency and an amplitude decaying as $1/t^2$ seems to be an inherent property of the temporal evolution of the drift mode in a plasma with cold ions [7]

Rayleigh–Taylor instability in plasma with homogeneous shear flow

In paper [8] the temporal evolution of the Raleigth–Taylor instability of plasma with shear flow was considered.The governing equations of the problem considered are the equation for the quasineutrality of the current density and the continuity equation for electrons. In the drift approximation the system of equations for the flute perturbations in the case of cold ions, $T_i \to 0$, but for the finite electron temperature, T_e, reduces to the following system for the perturbed electrostatic potential ϕ and for the perturbation of the electron density n

$$\left(\frac{\partial}{\partial t}+\mathbf{V}\cdot\nabla\right)\nabla_\perp^2\phi = -\frac{v_A^2}{c^2}\frac{4\pi e}{T_e}v_{Re}\frac{\partial n}{\partial y}, \quad (24)$$

$$\left(\frac{\partial}{\partial t}+\mathbf{V}\cdot\nabla\right)n + v_{Re}\frac{\partial n}{\partial y} + \frac{en_{0e}}{T_e}\left(v_{de}-v_{Re}\right)\frac{\partial\phi}{\partial y} = 0, \quad (25)$$

where velocity $\mathbf{V}$ is determined by equation (3). The solution of the equations (24)–(25) was obtained perturbatively as an expansion in powers of the initial values of the perturbation of the electrostatic potential,

$$\phi(k,l,T) = \phi_{(1)}(k,l,T) + \phi_{(2)}(k,l,T) + \phi_{(3)}(k,l,T) + ... \quad (26)$$

For the times T in the interval $1 \ll T \ll (l\rho_s)^{-1}$, where the shear flow effects become important, the linear solution $\phi_{(1)}(T)$ in this interval is

$$\phi_{(1)}(T) \approx e^{-\frac{i}{2}ST}T^{-2} \times \left(C_1(k,l)\,T^{\nu_1} + C_2(k,l)\,T^{\nu_2}\right) \qquad (27)$$

where

$$\nu_{1,2} = \frac{1}{2} \pm \sqrt{\frac{1}{4} + \left(\frac{\gamma_0}{v_0'}\right)^2} \qquad (28)$$

and

$$\gamma_0 = \frac{(v_{Re}v_{de})^{1/2}}{\rho_s}. \qquad (29)$$

It was shown that flow shear leads to the linear damping of electrostatic potential when the condition $(1/\sqrt{2})\gamma_0 \leq |v_0'|$ holds. This condition however is insufficient for the suppressing of the nonlinearly excited perturbations. It is obtained that in the range of the velocity shear $(1/\sqrt{2})\gamma_0 < v_0' < (2/\sqrt{3})\gamma_0$ nonlinearly excited perturbation of potential $\phi_{(2)}$,

$$\phi_{(2)}(k,l,T) \sim exp\left(-i\frac{ST}{2}\right)T^{2\nu_1-3}, \qquad (30)$$

grows algebraically. For $v_0' > (2/\sqrt{3})\gamma_0$ nonlinear perturbations $\phi_{(2)}$ also will be damped, but the nonlinearly excited potential $\phi_{(3)}$,

$$\phi_{(3)}(k,l,T) \sim \exp\left(-i\frac{ST}{2}\right)T^{3\nu_1-4}, \qquad (31)$$

will grow. The damping of the potential $\phi_{(3)}$ needs more severe condition for flow shear parameter v_0', which is determined by the condition $|v_0'| > (6/4)\gamma_0$. This condition admits the nonlinear growth of the potential $\phi_{(4)}$,

$$\phi_{(4)}(k,l,T) \sim \exp\left(-i\frac{ST}{2}\right)T^{4\nu_1-5}. \qquad (32)$$

This growth will be suppressed when the condition $|v_0'| > 4/\sqrt{5}\gamma_0$ is fulfilled. Clearly that the suppression of the nonlinear perturbations of potential $\phi_{(i)}$ of more higher order, $i \geq 5$, will demand even more strong flow shear. This damping of the nonlinearly excited perturbations, as well as linear ones is non-modal effect which appeared in times $t \geq (v_0')^{-1}$. This effect will be appreciable for the instability evolution, when this time is comparable with inverse growth rate, i.e. for $v_0' \sim \gamma_0$. So even under above conditions the suppression of the perturbation of the electrostatic potential by flow shear may be real physical process in cases when the terms of the order of $O(C^5_{1,2}(k,l))$ and of higher order may be safely omitted. At the same time perturbation of the electron density n does not suppressed by flow shear on the interval $1 \ll T \ll (l\rho_s)^{-1}$. In times $T > (l\rho_s)^{-1}$ potential oscillates with frequency lv_{Re} and with decaying with time as τ^{-2} amplitude and linear perturbation of the electron density evolves as $e^{ilv_{Re}\tau}$ with the amplitude which stays constant for these times.

Obtained non–modal solutions appear in times $T \geq 1$ or for dimensional time $t \geq v_0'$ and obviously absent in shearless case. The above condition for the stabilization of the linear RT instability and the condition $t \geq v_0'$ give $\gamma_0 t \sim 1$. That means that for shear flow with $(1/\sqrt{2})\gamma_0 \leq |v_0'|$ the non–modal solution of the algebraic form is settled in time of the order of the inverse growth rate. For this case any nonlinear processes with modal solutions can't develop. In the case of more weaker shear the non–modal solutions also may be developed, but in times which are much more longer then the inverse growth rate. In that case prior the development of such non-modal solutions nonlinear modal processes will define the development of the RT instability. So the condition $(1/\sqrt{2})\gamma_0 \leq |v_0'|$, as well as the other limitations on v_0' obtained look like as the bifurcation conditions, which lead to the principally different non–modal linear and nonlinear evolution of the RT instability.

At the condition $v_0' \sim \gamma_0$ the Rayleigh–Taylor instability may be stabilized also due to the enhanced by shear flow nonlinear decorrelation effect. The above analysis shows, that in the case of the RT instability the theory of the nonlinear decorrelation effect have to include the non–modal development of the potential and density perturbations and may be developed on the base of the non–sine type solutions (24), (27), (28), (29) for the electrostatic potential perturbations.

Conclusions.

The above results show, that in times of the order of the inverse velocity shear parameter, $(v_0')^{-1}$, the drift–like waves and instabilities in plasma with shear flow develop as a non–modal formations with time dependent wave numbers and frequencies. The stagnation and reflection by shear flow the wave packets of these perturbations may result in the reducing the anomalous transport of particles and energy in the regions of plasma with sufficiently strong velocity shear.

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