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ORIGINATION OF A SELF-OSCILLATING MODE (MAGNETIC STRIATIONS) IN A NONEQUILIBRIUM MAGNETIZED PLASMA

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In this paper, quantitative computations of the nonlinear solution of the problem of ionization instability development in a bounded domain [1], performed by the Lyapunov–Schmidt method [2], are presented. The amplitude of the self-oscillations is computed, the domains of the hard and soft modes of the loss of stability are isolated, a distribution of the electron density and electrical current over the channel section is constructed for the soft mode of the loss of stability – nonlinear magnetic striations. The topology of the striations in the post-critical domain is discussed. It is shown that the maximum of the steady-state wave amplitudes does not correspond to that wave which first lost stability. The results obtained are used for a qualitative analysis of experimental results with a nonequilibrium magnetized plasma in a magnetic field (the existence of oscillations at small wavelengths in a full ionization mode of the admixture).

§1. Let us examine the behavior of a nonequilibrium magnetized plasma in a domain bounded by two nonconducting walls x = 0 and x = b, which are infinite in the y direction. The magnetic field induction vector is directed along the z axis. Let us assume the parameters of heavy particles (atoms and ions) to be independent of the coordinates and time, while the ionization equilibrium build-up time is considerably less than the characteristic time of the problem. We consider the Reynolds magnetic number small and we neglect the effects of radiation. Taking account of these assumptions, the system of equations describing the state of the medium reduces to a dimensionless system of n partial differential equations in the potential Φ_n and the electron concentration Θ_n [1]. The system is solved by the method of a series expansion in the small supercriticality parameter $\varepsilon = (\Omega - \Omega^-)/\Omega^{-1}$. In a zero approximation (n = 0) the system has the form

$$L_{11}^0\Phi_0 + L_{12}^0\Theta_0 = 0,\ L_{21}^0\Phi_0 + L_{22}^0\Theta_0 = 0 \tag{1.1}$$

with the boundary conditions (see [3])

$$\Phi_0(0,\ Y) = \Phi_0(1,\ Y) = 0,\ \Theta_0(0,\ Y) = \Phi_0(1,\ Y) = 0, \tag{1.2}$$

where

$$L_{11}^0 = \frac{d^2}{dx^2} - k_y^2;\ L_{12}^0 = -a_1\frac{d}{dx} - ik_y\,\Omega^-;$$

$$L_{21}^0 = 2\frac{d}{dx};\ L_{22}^0 = -\Lambda L_{11}^0 + f_1;\ Y = y + W_0 t;$$

Λ is a small parameter; Ω^- is the critical Hall parameter; k is the wave vector; and a_1 and f_1 are constant factors [1, 4].

The solution of (1.1) with the boundary conditions (1.2) can be represented in the form

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$$\Phi_0(x,\ Y) = \alpha\left[\varphi(x)\ e^{ik_yY} + \varphi^*(x)\ e^{-ik_yY}\right], \tag{1.3}$$
$$\Theta_0(x,\ Y) = \alpha\left[\Theta(x)\ e^{ik_yY} + \Theta^*(x)\ e^{-ik_yY}\right],$$

where α is the unknown amplitude of the self-oscillations whose magnitude is found from subsequent approximations, and the superscript asterisk denotes the complex-conjugate quantity.

Maps of the level lines of the functions $\Phi_0(x, Y)$ and $\Theta_0(x, Y)$ are presented in Figs. 1 and 2. It must be noted that the level lines of the function $\Phi_0(x, Y)$ are streamlines of the disturbed state which is superposed on the fundamental background (homogeneous and stationary solution); hence, the current direction is indicated by an arrow in Fig. 1.

A simple formula to compute the neutral curve $\Omega^-(k_y): \mathrm{Im}(W_0) = 0$ and $\mathrm{Re}(W_0) = U$, which corresponds to striation motion at a velocity calculated in [5], where the agreement between this velocity and that obtained in experiments is mentioned, is found successfully from the zero approximation.

Substituting the real (or imaginary) part of (1.3) into (1.1), we obtain a homogeneous system of equations in the amplitudes, whose condition for nontrivial compatibility

$$\Lambda k_x^4 + (2\Lambda k_y^2 + f_1 + 2a_1)\,k_x^2 + 2\Omega^-\,k_yk_x + k_y^2\,(\Lambda k_y^2 + f_1) = 0 \tag{1.4}$$

has four roots k_{xj}, since the general solution of (1.1) contains four arbitrary constants. Substituting this solution into the boundary conditions, we obtain a homogeneous system of equations to determine the arbitrary constants, whose nontrivial compatibility condition has the form

$$\begin{vmatrix} 1 & 1 & 1 & 1 \\ e^{ik_{x1}} & e^{ik_{x2}} & e^{ik_{x3}} & e^{ik_{x4}} \\ \varkappa_1 & \varkappa_2 & \varkappa_3 & \varkappa_4 \\ \varkappa_1 e^{ik_{x1}} & \varkappa_2 e^{ik_{x2}} & \varkappa_3 e^{ik_{x3}} & \varkappa_4 e^{ik_{x4}} \end{vmatrix} \tag{1.5}$$

Here

$$\varkappa_j(k_{xj}) = -\frac{2ik_{xj}}{\Lambda\left(k_{xj}^2 + k_{yj}^2\right) + f_1}, \tag{1.6}$$

and (1.5) is for determination of the critical Hall parameter Ω^-.

Let us investigate the nature of the roots of (1.4) by using the Sturm theorem [6] and considering all the coefficients of (1.4) positive. Calculations of the quantities a_1 and f_1 for inert gases with an easily ionized admixture show that for a sufficiently high separation of the temperatures this assumption is always valid. The Sturm system for the polynomial $\tau(k_x)$ in the left side of (1.4) has the form

$$\tau_0 = g_0k_x^4 + g_2k_x^2 + g_3k_x + g_4, \tag{1.7}$$
$$\tau_1 = 4g_0k_x^2 + 2g_2k_x + g_3,$$

$$\tau_2 = -\frac{1}{2}\,g_2k_x^2 - \frac{3}{4}\,g_3k_x - g_4,$$
$$\tau_3 = -h_1k_x - h_2,\ \tau_4 = \left(g_2\frac{h_2}{h_1} - \frac{3}{4}g_3\right)\frac{h_2}{h_1} + g_4,$$

where $g_0 = \Lambda$; $g_2 = 2\Lambda k_y^2 + f_1 + 2a_1$; $g_3 = 2\,\Omega^-\,k_y$; $g_4 = k_y^2\ (\Lambda k_y^2 + f_1)$; $h_1 = 2g_2 - 8g_0g_4/g_2 + 9g_0g_3^2/g_2^2$; and $h_3 = g_3 + 12g_0g_3g_4/g_2^2$.

It can be shown that $h_1 > 0$. An analysis of the coefficients of the higher terms in the Sturm system (1.7) results in the deduction that the presence of real roots of (1.4) is governed by the sign of τ_4. The equation

$$\tau_4 = \left(g_2\frac{h_2}{h_1} - \frac{3}{4}g_3\right)\frac{h_2}{h_1} + g_4 = 0 \tag{1.8}$$

defines a certain curve $\Omega_0^-(k_y)$. For $\Omega^- < \Omega_0^-$, Eq. (1.4) has no real roots, while for $\Omega^- > \Omega_0^-$ it has two real and two complex roots.

Let us examine the case $\Omega^- > \Omega_0^-$. Using the Viette formulas, the roots of (1.4) can be represented in the form:

$$k_{x1} = -c - ir,\ k_{x2} = -c + ir,\quad k_{x3} = c + q,\ k_{x4} = c - q, \tag{1.9}$$

where c, r, and q are real and positive. The equalities

$$\varkappa_{1r} = -\varkappa_{2r},\ \varkappa_{1i} = \varkappa_{2i} \tag{1.10}$$

are satisfied for complex roots. The real and imaginary parts of the appropriate quantities are denoted by the subscripts r and i in (1.10).

Equation (1.5) can be written in the form

$$\begin{aligned}(\varkappa_1\varkappa_4 + \varkappa_2\varkappa_3)\left[e^{i(k_{x3}+k_{x4})} + e^{i(k_{x1}+k_{x2})} - e^{i(k_{x2}+k_{x4})} - e^{i(k_{x1}+k_{x3})}\right] + \\ + (\varkappa_1\varkappa_2 + \varkappa_3\varkappa_4)\left[e^{i(k_{x2}+k_{x4})} + e^{i(k_{x1}+k_{x3})} - e^{i(k_{x2}+k_{x3})} - e^{i(k_{x1}+k_{x4})}\right] + \\ + (\varkappa_1\varkappa_3 + \varkappa_2\varkappa_4)\left[e^{i(k_{x2}+k_{x3})} + e^{i(k_{x1}+k_{x4})} - e^{i(k_{x3}+k_{x4})} - e^{i(k_{x1}+k_{x2})}\right] = 0.\end{aligned} \tag{1.11}$$

Using (1.9) and (1.10), it can be shown that the left side of (1.11) is purely imaginary and (1.11) has the form

$$\begin{aligned}2\varkappa_{1r}(\varkappa_{4i} - \varkappa_{3i})\cos(2c) - \left[\varkappa_{1r}^2 + \varkappa_{1i}^2 + \varkappa_{3i}\varkappa_{4i} - \varkappa_{1i}(\varkappa_{3i} + \varkappa_{4i})\right] \times \\ \times (e^r - e^{-r})\sin q - \varkappa_{1r}(\varkappa_{4i} - \varkappa_{3i}) \cdot (e^r + e^{-r})\cos q = 0.\end{aligned} \tag{1.12}$$

For any value of k_y, Eq. (1.12) has an infinite number of solutions: $\Omega_0^-(k_y) < \Omega_1^-(k_y) < \Omega_2^-(k_y) < \ldots (\Omega_0^-(k_y)$ is a solution of (1.12) because two real roots of (1.4) hence agree; therefore, $q = 0$ and $\varkappa_{4i} = \varkappa_{3i}$. This means that for any k_y there exists an infinite set of solutions with different growth increments, where the increment equals zero for $\Omega_0 = \Omega_m^-$ for the m-th solution and is greater than zero for $\Omega_0 > \Omega_m^-$.

Let us consider the case $\Omega^- < \Omega_0^-$. The roots of (1.4) can hence be represented in the form

$$k_{x1} = -c - ir,\ k_{x2} = -c + ir,\quad k_{x3} = c + i\mu,\ k_{x4} = c - i\mu, \tag{1.13}$$

where c, r, and μ are real and positive. Analogously to (1.10), the following equalities are satisfied:

$$\varkappa_{2r} = -\varkappa_{1r},\ \varkappa_{4r} = -\varkappa_{3r},\quad \varkappa_{2i} = \varkappa_{1i},\ \varkappa_{4i} = \varkappa_{3i}. \tag{1.14}$$

Using (1.12) and (1.14), it can be shown that the left side of (1.11) is real:

$$-8\varkappa_{1r}\varkappa_{3r}\cos(2c) + \left(\varkappa_{1r}^2 + \varkappa_{1i}^2 + \varkappa_{3i}^2 - 2\varkappa_{1i}\varkappa_{3i}\right)(e^r - e^{-r})(e^\mu - e^{-\mu}) + 2\varkappa_{1r}\varkappa_{3r}(e^r + e^{-r})(e^\mu + e^{-\mu}) = 0. \tag{1.15}$$

Computations show that (1.15) has no solution in the range $\Omega^- < \Omega_0^-$.

An investigation of the solution whose growth increment vanishes for $\Omega^- = \Omega_0^-$ results in the deduction that it should be discarded, since two roots of (1.14) hence agree.

Therefore, the neutral curve separating the stability from the instability domain is $\Omega_1^-(k_y)$.

Results of computing the neutral curves of the first mode Ω_1^-, the second Ω_2^-, the third Ω_3^-, and the dependence $\Omega_0^-(k_y)$ are presented in Fig. 3. Also presented there for comparison is the netural curve of the first mode (m = 1) computed by means of the approximate formula [1]

$$\left(\Omega_{k_y}^-\right)^2 - k_y^2\left(2\Lambda k_y^2 + f_1 + 2a_1\right)\left(\Lambda k_y^2 + f_1\right) = (\pi m)^2\left(2\Lambda k_y^2 + f_1 + 2a_1\right)^2, \tag{1.16}$$

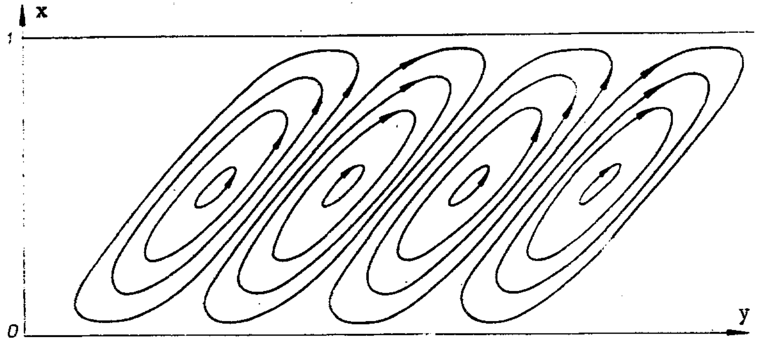

Fig. 1

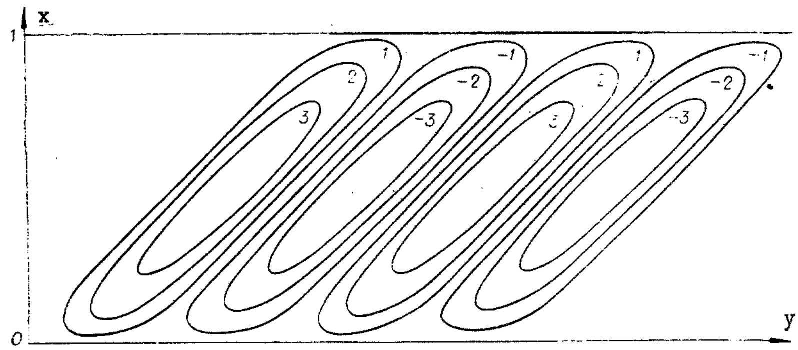

Fig. 2

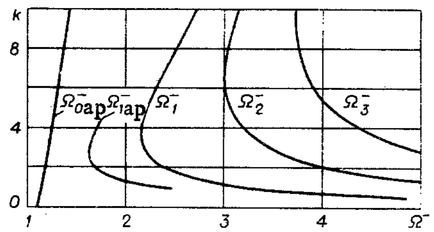

Fig. 3

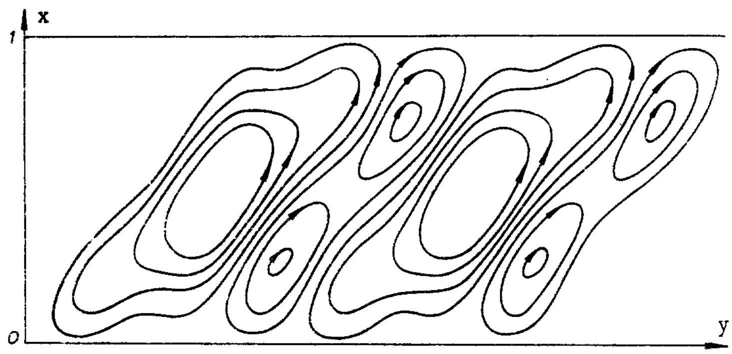

Fig. 4

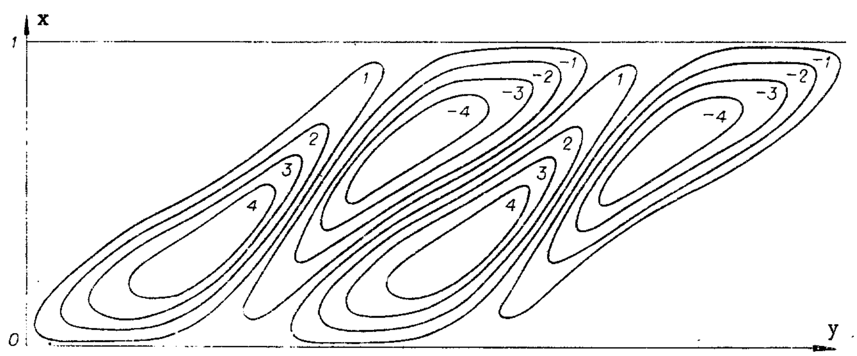

Fig. 5

which can be obtained from (1.12) by using the fact about the presence of a small parameter Λ. Computations are carried out for an Ar $+2.3 \cdot 10^{-5}$ Cs plasma (the pressure $p=1$ kg/cm^2, $T_a=1500$°K, $T_e=3200$°K, and $\Lambda = 0.5 \cdot 10^{-1}$). The coefficients a_1 and f_1 were computed by using the experimental dependences presented in [7]. Comparing the neutral curves computed by using (1.12) and (1.16) shows that the discrepancy can be significant.

§2. To solve the nonlinear problem it is necessary to find the next approximations $n \geq 1$. Values of Φ_1 and Θ_1 are sought from the approximation $n=1$ by the method elucidated in [1]. Qualitative sketches of the level lines of the functions $\Phi = \varepsilon\Phi_0 + \varepsilon^2\Phi_1$ and $\Theta = 1 + \varepsilon\Theta_0 + \varepsilon^2\Theta_1$ (ε is the supercriticality parameter; $\varepsilon^2 = \Omega^- + \Omega$), obtained on the basis of computations of Φ_0, Θ_0, Φ_1, Θ_1, are presented in Figs. 4 and 5. The computed domain of elevated electron concentration (Fig. 5) has the same shape as the luminescing discharge domain in the experiments [8, 9] corresponding to the assumptions made in this paper. Attention should be turned to the presence of closed disturbed streamlines with a characteristic dimension equal to half the channel width b (Fig. 4).

The amplitude of the self-oscillations α^2 can be computed in the next step ($n=2$). An approximate formula is presented for α^2 in [1]. An exact formula for α^2, which has a quite awkward form, is obtained on the basis of the method described in [1]. Comparing the results of computing α^2 by the exact and approximate formulas showed that the approximate formula truly describes the behavior of the self-oscillation amplitude qualitatively for $k \geq 1.5$. However, to construct the self-oscillating mode it turns out to be necessary to study the behavior of α^2 due to k for large wavelengths (small k). As computations showed, in the range of k values from zero to ~ 1.5 $\alpha^2 > 0$, the amplitude changes sign with the increase in k. It is important to determine the maximum positive value, since the decrement of the nonlinear oscillation vanishes at this value. Computations of the self-oscillation amplitude permits extraction of domains of soft and hard modes of loss of stability. An investigation of the nature of the loss of stability is important for clarification of the subsequent plasma behavior.

According to the exact formula for α^2 (which is awkward and is consequently not presented here), computations for the self-oscillating mode were performed for an argon plasma with an easily ionized component — cesium — as admixture. Computations were carried out in the 2000-6000°K electron temperature range for fractional values of the admixture between 10^{-2} and 10^{-5}. The results of computations are shown in Fig. 6.

For $\alpha^2 > 0$ a soft mode of loss of stability holds, and for $\alpha^2 < 0$ a hard mode holds. The point k^+ corresponds to $\alpha^2 = 0$, i.e., the passage from the soft to the hard mode of the loss of stability.

Curves 1 and 3 in Fig. 6 are computed for $T=3000$°K and $\Lambda = 10^{-3}$, but for different fractions of the Cs admixture: 1) $\delta = 10^{-5}$; 3) $\delta = 10^{-2}$. It is seen from a comparison of the dependences $\alpha^2(k)$ for different fractions of the admixture that a diminution in δ shifts the point k^+ into the domain of shorter wavelengths (large k), the maximum amplitude hence increasing substantially.

Curves 4 and 1 in Fig. 6 have been computed for $\delta = 10^{-5}$ and $\Lambda = 10^{-3}$, but for different temperatures: 4) 2000°K; 1) 3000°K. Comparing curves 4 and 1 shows that a rise in temperature influences the behavior of the dependence $\alpha^2(k)$ substantially: the maximum amplitude grows and the transition point (k^+, Ω^+) shifts into the domain of large k.

Let us compare curves 2 and 3 computed for $T=3000$°K and $\delta = 10^{-2}$, but for different Λ : 2) $\Lambda = 10^{-2}$, 3) $\Lambda = 10^{-3}$. It is seen from the computations that a change in the thermal conductivity (a small parameter) in the band mentioned will influence the perturbation amplitude slightly.

A diminution in the fraction of the admixture or a rise in the temperature (or both factors simultaneously) leads to an increase in Ω^- as well as a growth in the maximum of α^2 and shifts the point k^+ of the passage from the soft to the hard modes of the loss of stability into the domain of large values of the number k (small wavelengths).

Self-oscillations corresponding to the magnetic striations observed in experiments [9] occur in the soft mode of the loss of stability. The question of selecting the amplitude of the steady wave is an independent problem, since $\alpha^2 > 0$ for the whole wavelength range in the soft mode of the loss of stability. In some cases the amplitude of the steady wave visibly depends on the kind of initial perturbation. The amplitude of the steady wave is selected corresponding to α_{max}, since the decrement of the nonlinear oscillations vanishes at this value. It should be noted that the wavelength corresponding to the steady amplitude does not correspond to that wave which first loses stability.

This is apparently explained by the fact that waves with minimal linear and nonlinear decrements can be distinct in the subcritical domain. Only waves with a linear decrement are taken into account in a linear analy-

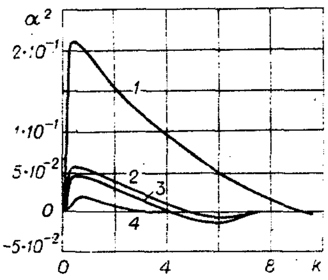

Fig. 6

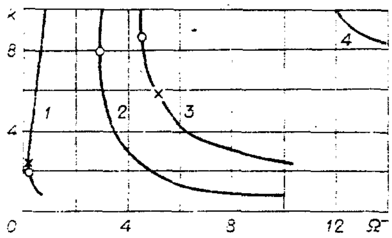

Fig. 7

sis of the stability. In the post-critical domain (the decrement goes over into an increment) a wave corresponding to a nonlinear increment, and not the wave which first lost stability according to the linear increment, plays the main part.

§3. The methods elucidated above for computing the neutral curves and amplitudes of the self-oscillations can be used for a qualitative explanation of experiments [9-11]. An Ar + Cs plasma placed between coaxial electrodes was studied in these experiments. The dimensions of the cathode and anode, and the height of the electrodes, were 13.38 and 40 mm, respectively. The fraction of the admixture varied between $0.5\cdot10^{-5}$ and $0.5\cdot10^{-4}$ and the electron temperature varied between 2000 and 4000°K. For the geometry of the experiment, the small parameter Λ took on values from 10^{-3} to 10^{-2}, while the wave number k varied between 1 and 10.

The computations were performed for the temperatures 2000 and 4000°K and the following values of the admixture fraction for each temperature: 10^{-3}, 10^{-4}, and 10^{-5}.

Results of the neutral curve computations are presented in Fig. 7: 1) T = 2000°K, $\delta = 10^{-4}$; 2-4) T = 4000°K and $\delta = 10^{-3}$, 10^{-4}, 10^{-5}, respectively. As is seen from Fig. 7, for the conditions of the experiments [9-11], the plasma went from the instability into the stability domain with the diminution in the fraction of admixture and with the rise in temperature; the instability domain occurs at higher values of Ω^-, and the critical value of the Hall parameter shifts into the short-wavelength domain.

Let us compare the neutral curves 1 and 3 in Fig. 7. Both neutral curves were computed for $\delta = 10^{-4}$, but for different temperatures: 1) T = 2000°K; 3) T = 4000°K. Points of minimum $\Omega^-(k^-)$ are denoted by a circle, and the point $k^+(\alpha^2 = 0)$ by a cross. Comparison shows that the minimum value of the Hall parameter on the neutral curve from the soft mode domain goes over into the domain of the hard mode of the loss of stability with the increase in temperature (k^+ is greater than k^- for T = 2000°K, while k^+ is below k^- for T = 4000°K; $\alpha^2 > 0$ from k^+ to $k = 0$ while $\alpha^2 < 0$ from k^+ and above). The computations confirmed that a temperature rise resulting in an increase in the degree of admixture ionization results in broadening the stability domain, especially for high wavelengths.

It was detected in the experiments [10, 11] that saturation of the effective Hall parameter was nevertheless observed in the full ionization mode for the admixture. The assumption that this effect can be explained by the presence of a hypothetical microscale inhomogeneity, not recovered by the instruments, was expressed.

However, another explanation of the effect noted can be given. As the warming current increases (the fraction of the admixture and the magnetic field are fixed), it is possible to go successively from the unstable domain with $\Omega_1 > \Omega^-$ into the stable domain where $\Omega_1 < \Omega^-$ (Ω_1 is some fixed value of the Hall parameter).

Let $\Omega_1 > \Omega^-$; fluctuations are observed in the plasma. As the temperature (current) rises, Ω^- shifts to the right and upward from the origin, as is seen in Fig. 7 (toward high values of both the magnitude of the critical value of the Hall parameter and the magnitude of the wave number corresponding to Ω^-).

If the parameter Ω^- was in the domain of the soft mode of the loss of stability prior to the heating, then as the temperature rises Ω^- can turn out to be in the domain of the hard mode of the loss of stability (in this case a self-oscillating mode exists for $\Omega < \Omega^-$). In this latter case, since the full ionization state of the admixture is achieved by a transition from the unstable domain, the fluctuations do not vanish at once. Although the temperature corresponds to full ionization of the admixture and the corresponding Hall parameter Ω_1 is less than the critical value Ω^-, in this state, oscillations which result in saturation of the effective Hall parameter can exist. It is not excluded that the characteristic wavelengths of the self-oscillating mode which can occur in the subcritical domain in this case were not fixed in the experiments [10, 11]. If the appropriate physical diagram is true, the results of the experiment will depend on how the full ionization mode of the admixture is achieved (hysteresis occurs). If the full-ionization mode of the admixture is achieved from the stable state in which there are no fluctuations, the effective plasma parameters (conductivity and Hall parameter) should agree with the laminar values. If the full-ionization mode of the admixture is achieved from the state in which oscillations exist, the effective parameters can visibly differ from the laminar values.

For a more detailed comparison between computations and the experiments [10, 11], it is certainly necessary to perform additional computations with more exact values of the transfer coefficients.

The results presented here have been obtained for small values of the supercriticality parameter, but, as has been shown in [12] even for higher values of the supercriticality parameter, ionization waves of finite amplitude – magnetic striations – can exist. Only when the second critical value exceeds the Hall parameter does an ionization turbulence mode occur. The model proposed by Landau for the origin of turbulence holds for the soft mode of the loss of stability. As more detailed investigations show, magnetic striations can exist even for moderate values of the supercriticality parameter. As the supercriticality parameter increases further, a modulation instability of the magnetic striations starts to develop. The mutual interaction between the waves should be taken into account in the hard mode of the loss of stability.

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