THE STABILITY OF A SPATIALLY INHOMOGENEOUS PLASMA IN A MAGNETIC FIELD

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The oscillations and stability of an inhomogeneous plasma in a magnetic field are studied without assumptions on the smallness of the Larmor ion radius in comparison with the wavelength. The analysis is carried out for the particular case of a low pressure plasma layer in a homogeneous magnetic field. The kinetic equation is used and collisions are neglected. It is shown that in a plasma with an inhomogeneous density and temperature, perturbations exist which are unstable for infinitesimal inhomogeneities and arbitrary relative values of the density and temperature gradients (in this sense, the instability is universal). Perturbations with a wavelength along the magnetic field greater than the characteristic length of density or temperature changes are unstable. The transverse wave length of the most unstable perturbations is of the order or smaller than the ion Larmor radius. Ion-acoustic and Alfven oscillations that move in a direction almost perpendicular to the magnetic field correspond to such disturbances in a homogeneous plasma. The maximum increment of the resultant instability is equal to $(T_i/Ma^2)^{1/2}$.

1. In the present paper, small oscillations and the stability of a plane infinite layer of a low pressure inhomogeneous, collision-free plasma ($\beta = 8\pi p/H^2 \ll 1$) are considered in a magnetic field whose lines of force are straight lines. This problem (which is important for the question of magnetic isolation of a plasma) has been studied in a number of researches.[1-5] The stability of the plasma relative to longitudinal (curl $\mathbf{E} = 0$) long-wave disturbances ($k^2 r_H^2 \ll 1$; k is the wave number, r_H the ion Larmor radius) was considered in [3,4]. These perturbations correspond to ion sound in a homogeneous plasma. These oscillations were considered in [5], but for the case of an arbitrary relation between the wavelength of the oscillations and the mean ion Larmor radius.

In the present work we consider in detail the stability of the oscillations in an inhomogeneous plasma without any assumptions on their longitudinal character, and for arbitrary kr_H. It will be shown that the plasma is unstable for any values of the quantity $\partial \ln T/\partial \ln n$. Both the longitudinal oscillations and disturbances which bend the magnetic lines of force are unstable. (In a homogeneous plasma, these latter disturbances correspond to Alfven waves.)

2. In this section, we shall give a short derivation of the dispersion equation for long wave, $kr_H \ll 1$, oscillations of a low pressure inhomogeneous plasma. The dispersion equation for an arbitrary relation of the wavelength of the excitation and the ion Larmor radius will be given without derivation in Sec. 4.

Oscillations are considered which weakly perturb the density of the magnetic lines of force of the field $\mathbf{H}_0$, $H_z \sim 8\pi p_0 n/H_0 n_0$. In a homogeneous plasma, these are the ion-acoustic and Alfven waves.

We choose a rectangular set of coordinates, the z axis of which is directed along H_0, and the x axis of which is in the direction of the plasma density gradient. We seek perturbations of all these quantities in the form

$$A(x, y, z, t) = A(x) \exp(-i\omega t + ik_y y + ik_z z).$$

If the phase velocity of the oscillations ω/k_z is large in comparison with the thermal velocity of the ions $u_i = \sqrt{2T_i/M}$, then the motion of the ionic component of the plasma can be considered in the hydrodynamic approximation; for excitations with $k_z/k_y \ll 1$, one can neglect the motion of the ions along the magnetic field.

Since we shall be interested in the frequency $\omega \sim k_y c\nabla nT/eHn$, we shall then keep the terms with "magnetic viscosity" in the equation of motion of the ionic component of the plasma:

$$Mn\left(\frac{\partial \mathbf{v}_i}{\partial t} + (\mathbf{v}_i \nabla)\ \mathbf{v}_i\right) = en\,(\mathbf{E} + c^{-1}\,[\mathbf{v}_i\,\mathbf{H}]) - \nabla p_i - \operatorname{div} \pi_{ik},$$

$$\frac{\partial p_i}{\partial t} + \mathbf{v}_i \nabla p_i + \gamma p_i \operatorname{div} \mathbf{v}_i + \operatorname{div} \mathbf{q}_i = 0,$$

$$\mathbf{q}_i = \gamma \frac{cp_i}{eH^2}\,[\mathbf{H}\nabla T_i],$$

$$\pi_{yy} = -\pi_{xx} = \frac{nT_i}{2\omega_H}\left(\frac{\partial v_y}{\partial x} + \frac{\partial v_x}{\partial y}\right),$$

$$\pi_{xy} = -\pi_{yx} = \frac{nT_i}{2\omega_H}\left(\frac{\partial v_x}{\partial x} - \frac{\partial v_y}{\partial y}\right),$$

$$\mathbf{v}_i = \{v_x, v_y, 0\}, \quad \mathbf{H} = \{H_x, 0, H_0\}, \quad \omega_H = eH/Mc. \qquad (1)^*$$

This set of equations can be found in the work of Braginskiĭ.[6] It is also used in [7].

Neglecting terms of order $\nabla n_0/n_0 k_y$, $k_y^{-1}\partial/\partial x$, $Hz/H_0 \sim 8\pi p_0 n/H_0^2 n_0$, we obtain the result that

$$\operatorname{div} n\mathbf{v}_i = k_y\left(\omega - k_y c\frac{\nabla p_{0i}}{eH_0 n_0}\right)\frac{Mc^2 n_0}{eH_0^2}E_y - c\frac{E_y}{H_0}\nabla n_0 \qquad (2)$$

in the frame in which $\mathbf{E}_0 = 0$ in the unperturbed plasma.

The drift approximation is applicable for electrons.[4] In the drift approximation,

$$\operatorname{div} n\mathbf{v}_e = ik_z n_0 v_z^e - cE_y(\nabla n_0)/H_0,$$

$$n_0 v_z^e = i\frac{e}{m}E_z\int_{-\infty}^{\infty}\left(\frac{\partial f_0^e}{\partial v_z} - \frac{k_y v_z}{\omega\omega_{He}}\frac{\partial f_0^e}{\partial x}\right)\frac{v_z\,dv_z}{\omega - k_z v_z}. \qquad (3)$$

From the set of equations (1)—(3) and the Maxwell equations

$$\operatorname{rot}\mathbf{H} = (4\pi/c)\,\mathbf{j} = (4\pi/c)\,en\,(\mathbf{v}_i - \mathbf{v}_e),$$

$$\omega H_x/c = k_y E_z - k_z E_y, \qquad (4)^\dagger$$

we find the dispersion equation

$$\left(\omega^2 - \omega k_y c\frac{\nabla nT_i}{eH_0 n} - k_z^2\frac{H_0^2}{4\pi nM}\right)\int_{-\infty}^{\infty}\left(\frac{\partial f_0^e}{\partial v_z} - \frac{k_y v_z}{\omega\omega_{He}}\frac{\partial f_0^e}{\partial x}\right)\frac{v_z\,dv_z}{\omega - k_z v_z}$$

$$= \left(\omega - k_y c\frac{\nabla nT_i}{eH_0 n}\right)\frac{k_y^2 c^2}{\omega_p^2}, \quad \omega_p^2 = \frac{4\pi ne^2}{m}. \qquad (5)$$

3. We investigate the dispersion equation (5) for the Maxwell distribution function

$$f_0^e(v_z) = \frac{n}{\sqrt{\pi}\,u_e}\exp\left(-\frac{v_z^2}{u_e^2}\right), \qquad (5a)$$

when $h_z u_i \ll \omega \ll k_z u_e$. For simplicity, we assume that the ion and electron temperatures are the same: $T_i = T_e = T$. Then Eq. (5) can be written in the form

$$[\omega - \omega_*/(1+\eta)$$
$$+ i(\sqrt{\pi}\omega/k_z u_e)(\omega - \omega_*(1-\eta/2)/(1+\eta))]$$
$$\times(\omega^2 + \omega\omega_* - k_z^2 V_A^2) = Zk_z^2 V_A^2(\omega + \omega_*), \qquad k_z > 0, \qquad (6)$$

where the following notation is used:

$$\omega_* = -k_y c\nabla nT/eH_0 n, \quad \eta = \partial\ln T/\partial\ln n, \quad V_A^2 = H_0^2/4\pi nM,$$
$$Z = k_y^2 r_H^2. \qquad (6a)$$

The small parameter Z is located on the right hand side of Eq. (6). If we neglect effects of the order of Z, then (6) divides into two equations. One of these describes the degenerate ion sound (drift wave), which is unstable for $\eta < 0$ and which was considered in detail in [4]. It was also shown there that the other branch of ion sound (with the small phase velocity) oscillates if $\eta > 2$. The other equation gives the Alfven oscillations in the inhomogeneous plasma; in this approximation it has a purely real frequency

$$\omega_{1,2} = -\omega_*/2 \pm \sqrt{\omega_*^2/4 + k_z^2 V_A^2}. \qquad (7)$$

Taking into account the terms of order Z, we find that now the frequency of the Alfven oscillations has a small imaginary part, the sign of which can be positive (which corresponds to the oscillations) for waves with

$$k_z < \omega_*/V_A \qquad (8)$$

under the condition that

$$-4 < \partial\ln T/\partial\ln n < 2. \qquad (9)$$

For fixed k_y, the increment depends on k_z and is maximal when the phase velocities of the drift and Alfven waves are equal (for $k_z = \sqrt{2\omega_*}/V_A$):

$$\gamma_{max} = \sqrt{\pi/12}\,\omega_* V_A\sqrt{Z}/u_e. \qquad (10)$$

It is shown in Fig. 1 how the ion sound and Alfven oscillations are deformed in an inhomogeneous plasma. For fixed k_y, the real part of the frequency $\operatorname{Re}\omega$ is plotted as a function of k_z. The dashed curves indicate $\omega(k_z)$ in the inhomogeneous plasma. The curve $\omega = \tilde{\omega} = -c(1 - \eta/2) \times T\nabla n/eHn$ separates the region of the unstable solutions of the dispersion equation (6), $\tilde{\omega}/\omega > 1$ (see Sec. 5).

The assumption made at the beginning of this section that $\omega \ll k_z u_e$ is equivalent for Alfven

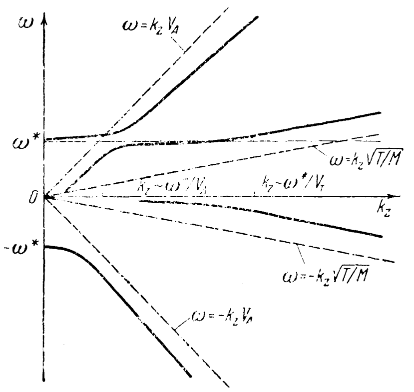

FIG. 1. Dependence of $\omega = \omega(k_z)$ for Alfven and ion-acoustic oscillations in an inhomogeneous plasma.

$^*(\mathbf{v}_i\nabla) = (\mathbf{v}_i\cdot\nabla)$, $[\mathbf{v}_i\mathbf{H}] = \mathbf{v}_i \times \mathbf{H}$.
†rot = curl.

waves, as it is not difficult to see, to the requirement that the plasma pressure be not too low.

In a highly rarefied plasma, $\beta \ll m/M$, the phase velocity of the Alfven waves is larger than the thermal velocity of the electrons u_e, $\omega \gg k_z u_e$, and in this case, Eq. (5) can be written in the form ($T_i = T_e$)

$$(\omega - \omega_*)(\omega^2 + \omega\omega_* - k_z^2 V_A^2) + k_y^2 c^2 \omega^2 (\omega + \omega_*)/\omega_p^2 = 0,$$
$$\omega_* = -k_y c \nabla nT/eH_0 n. \qquad (11)$$

This equation has complex roots, although we have omitted from the derivation the residues of Eq. (5) (they are exponentially small).

The maximum increment is obtained for $k_z = \sqrt{2}\omega_*/V_A$, and is equal to

$$\gamma = \sqrt{2/3}\, k_y c \omega_*/\omega_p. \qquad (12)$$

We note that this "hydrodynamic" increment differs from the "kinetic" increment for the case $\beta \gg m/M$ by a numerical factor only.

4. It is seen from Eq. (10) that the increments increase with decrease in the wavelength of the oscillations (with increase in Z). To estimate the values of the maximum increment, we use the dispersion equation, which is valid for arbitrary Z, and which was obtained by one of the authors (see [8]). It has the form ($k_z u_i \ll \omega \ll k_z u_e$)

$$\left[\omega - \omega_* \frac{1}{1+\eta} + i\frac{\sqrt{\pi}\omega}{k_z u_e}\left(\omega - \omega_* \frac{1-\eta/2}{1+\eta}\right)\right]\left[\omega^2 + \omega\omega_*^i - k_z^2 V_A^2 Z/(1 - I_0 e^{-Z})\right] = Z k_z^2 V_A^2 (\omega + \omega_*^i);$$
$$\omega_*^i = -k_y c T\bar{\varkappa}/eH, \qquad \bar{\varkappa} = d\ln\left[(1 - I_0 e^{-Z})\, n\right]/dx. \qquad (13)$$

Here $I_0 \equiv I_0(Z)$ is the Bessel function of imaginary argument.

For $Z \gg 1$, $\nabla T = 0$, the unstable solution of Eq. (13) is

$$\omega = \frac{k_z u_e}{\pi\sqrt{2Z}}\frac{a^2 + i}{a^4 + 1}; \quad a^2 = \frac{2}{\sqrt{\pi}}\frac{k_z u_e}{\omega_*}\left(1 + \frac{\omega_*^2}{2k_z^2 V_A^2 Z}\right). \qquad (14)$$

The increment of ($\mathrm{Im}\,\omega$) is a function of two parameters: k_z and Z. For fixed Z, the increment is maximum for $k_z^* \sim \kappa u_i/V_A$ (if $Z < \beta M/m$) and for $k_z^* \sim \kappa\sqrt{mZ/M}$ (if $Z > \beta M/m$). For $k_z \sim k_z^*$, the increment increases with increase in Z up to $Z \sim \beta M/m$, and then remains constant:

$$\gamma_{max} = \varkappa u_i/4\sqrt{2\pi}, \quad \varkappa = \nabla n/n. \qquad (15)$$

The real part of ω for $k_z \sim k_z^*$ and $Z > \beta M/m$ does not depend on Z and is of the order of the increment of γ_{max}. The dependences $k_z^* = k_z^*(Z)$ and $\gamma_{max} = \gamma_{max}(Z)$ for $T = 0$ are shown in Figs. 2 and 3.

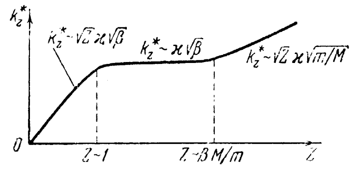

FIG. 2. Dependence of the wave number k_z^*, corresponding to the maximum increment, on Z.

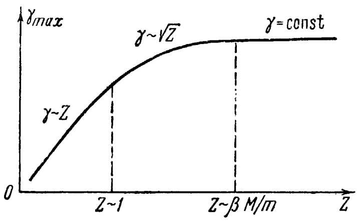

FIG. 3. Dependence of the increment of γ_{max} on Z.

The perturbations are unstable when $\nabla T = 0$ (for fixed Z); they can be shown to be stable in an inhomogeneously heated plasma ($\eta = \partial \ln T/\partial \ln n \neq 0$). If we restrict ourselves to the frequencies $\omega \ll k_z V_A$, then we see from Eq. (13) that

$$\omega - \frac{\omega_*}{1+\eta} - i\sqrt{\pi}\frac{\omega}{k_z u_e}\left[\omega - \omega_* \frac{1-\eta/2}{1+\eta}\right] = -(\omega + \omega_*^i)(1 - I_0 e^{-Z}). \qquad (16)$$

It is then not difficult to get the stability criteria (see also [9]):

$$(1 - \eta/2)(2 - I_0 e^{-Z})/(I_0 e^{-Z} + \eta Z \partial I_0 e^{-Z}/\partial Z) < 1. \qquad (17)$$

The stability regions in the (η, Z) plane are shown in Fig. 4.

In Eq. (13), we have omitted terms which take into account the motion of the ions along the magnetic field, assuming $|\omega| \gg k_z u_i$. Therefore, our consideration is valid for $k_z^* < \kappa$, i.e., for $Z < M/m$.

5. The results of the previous sections can easily be obtained from simple qualitative estimates. In fact, the instability of the vibrations under consideration is associated with the transfer of the energy of the thermal motion of the "resonant" particles ($v_z \approx \omega/k_z$) to the wave. The energy transferred to the wave per unit time and per unit volume $\dot{W} \sim 2\gamma W$ is equal to $\overline{j_z E_z}$, where the bar denotes the average value over the vibrations.

Let us compute this energy. We shall assume the resonant particles to be particles which, after

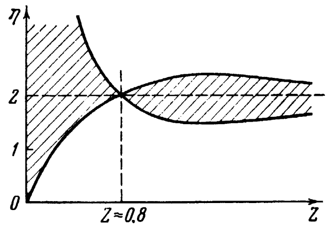

FIG. 4. Stability regions for longitudinal waves.

a time $\tau \gg 1/\omega$, move in phase with the wave along the magnetic field. They can have a velocity scatter Δv_z about the phase velocity ω/k_z of the order of $1/k_z\tau$. If $u_e > \omega/k_z \gg u_i$, then the amount of resonant ions is exponentially small, and of resonant electrons is

$$\delta n_0 \sim f_0^e \, \Delta v_z \sim n_0/u_e k_z \tau. \tag{17a}$$

After a time τ, the resonant electrons found in the field of the wave H_x, E_y are displaced in x by the distance

$$\Delta x = v_x \tau = (H_x v_z/H_0 + cE_y/H_0)\tau \approx cE_z k_y \tau / H_0 kz, \tag{17b}$$

which leads to a change in the density of resonant particles at the point x_0:

$$\delta n = -c\frac{E_z}{H_0}\frac{k_y}{k_z}\tau\frac{\partial \delta n_0}{\partial x} = -c\frac{E_z}{H_0}\frac{k_y \nabla n}{k_z^2 u_e}\left(1-\frac{1}{2}\frac{\partial \ln T_e}{\partial \ln n}\right), \tag{17c}$$

and to a change in electron current $\delta j_z = -e\omega\delta n/k_z$. Consequently,

$$\gamma W \sim e\frac{c}{H_0}\frac{k_y \nabla n \omega}{k_z^3 u_e}\left(1-\frac{1}{2}\frac{\partial \ln T}{\partial \ln n}\right)E_z^2 - \frac{\omega^2 e^2}{k_z^3 T_e u_e}E_z^2. \tag{18}$$

Here we have added to the right hand side of the relation a term, not associated with the inhomogeneity of the plasma, which takes into account the resonant absorption of the wave. This term could also be introduced from much the same descriptive situations. We rewrite Eq. (18) in the form

$$\gamma W \sim -\frac{\omega^2 e^2}{k_z^3 u_e T_e}\left[1 + c\frac{T_e k_y \nabla n}{\omega e H n}\left(1-\frac{1}{2}\frac{\partial \ln T_e}{\partial \ln n}\right)\right]E_z^2, \tag{19}$$

from which follows the criterion for instability:

$$1 - \tilde{\omega}/\omega < 0, \quad \tilde{\omega} = -k_y c\frac{T_e^* \nabla n}{eHn}\left(1-\frac{1}{2}\frac{\partial \ln T_e}{\partial \ln n}\right). \tag{20}$$

Thus the study of instability is reduced to finding the phase velocities of various types of waves.

APPENDIX

In the study of problems related to the stability of the plasma, various approaches are possible: first, one can investigate the temporal development of initial excitations; second, one can find the frequencies (generally complex) of the characteristic vibrations.

Above, we have studied the stability and the plasma oscillations under the assumption that the excited quantities change slowly in the direction of inhomogeneity of the plasma. As a consequence, we have omitted small terms that contain derivatives of the perturbed quantities in the coordinate x, and an algebraic equation is obtained in place of the differential one [for example, Eqs. (6) and (13)]. The frequency ω thus determined is a function of the coordinates, which requires interpretation.

We therefore clarify the meaning of the solutions found in this way from the viewpoint of the two approaches to the problem of instability that have been considered. We take up the evolution of the smooth ($k_x r_H \ll 1$, $k_y = E^{-1}dE/dx$) initial perturbation that is generated in a small region Δx about the point x_0. Let $r_H \ll \Delta x \ll a$ (r_H is the Larmor ion radius, a is the characteristic length of change of the unperturbed density and temperature in the plasma). If the excitation succeeds in growing appreciably before its spatial dependence on the coordinate x is changed, and the assumption that $k_x r_H \ll 1$ is violated, then the increment $\nu(x)$ that we have found is the local increment of growth of such perturbations. One can show that for this increment, the condition $\nu/\omega \ll r_H/a$ must be satisfied (this is the first method).

If this condition is not satisfied, then, in order to find the complex frequencies of the characteristic oscillations, one must solve the differential equation in which the small parameter $(r_H/a)^2$ or $(r_H k_y)^2$ replaces the highest-order derivative (second method).

We shall now show an example of longwave vibrations ($k_y r_H \ll 1$, $k_x r_H \ll 1$) such that the solution of the problem again leads to the dispersion equation (6). For this purpose, we keep in Eq. (5) the terms thrown away by us which contain the field derivatives in the x coordinate and which are formally small, since $(k_x/k_y)^2 \ll 1$. Then, for example, in place of the "dispersion equation" (6), we would have had the differential equation

$$r_H E_z'' - U(x, \omega) E_z = 0, \tag{A.1}$$

where

$$U = k_y^2 r_H^2 + \left[1 - \frac{\omega_*}{\omega}\frac{1}{1+\eta} + \frac{i\sqrt{\pi}\omega}{k_z u_e}\left(1 - \frac{\omega_*}{\omega}\frac{1-\eta/2}{1+\eta}\right)\right] \times \left(\frac{1}{1+\omega_*/\omega} - \frac{\omega^2}{k_z^2 V_A^2}\right). \tag{A.2}$$

We now consider the simple example in which $\eta = 0$, $\omega^2/k_z^2 V_A^2 \ll 1$. In this case,

$$U = 1 - \frac{\omega_*}{\omega} + k_y^2 r_H^2 + i\left[\gamma\frac{\omega_*}{\omega_r^2} + \left(1 - \frac{\omega_*}{\omega_r}\right)\frac{\sqrt{\pi}\omega_r}{k_z u_e}\right], \quad \omega = \omega_r + i\gamma, \ \gamma \ll \omega_r. \tag{A.3}$$

In view of the smallness of Im U in comparison with Re U (everywhere except for a small region about the point x, where Re U = 0), the spatial path of the solution is determined by the real part of the potential. Let the density change monotonically.

Then ReU is a potential "well," if the phase velocity of the wave is directed along the drift velocity of the electrons, i.e., $\omega_r/\omega_* > 0$, and there is a potential "bulge," if $\omega_r/\omega_* < 0$. Thus, for $\omega_r/\omega_* > 0$, Eq. (A.1) with the potential (A.3) has solutions that are localized in the region of inhomogeneity between the turning points which are determined approximately by the equation $\operatorname{Re} U = 0$.

Since Eq. (A.2) is determined under the assumption that $k_x r_H \ll 1$, then it is valid only in the range of values of x where $U(x) \ll 1$. In the range of Δx of order a, $1/a = n^{-1} dn/dx$, the potential U changes by a quantity of the order of unity. Therefore, the condition $U(x) \ll 1$ can be satisfied only close to the "bottom" of the potential well, when the dimension of localization of the solution $\Delta x \ll a$. In this region, the change of the potential is small and it can be expanded in a series about the point x_0, [1)] where $\partial U/\partial x = 0$, and reduces Eq. (A.1) to the equation for the oscillator. The characteristic values of such a problem are determined from the equation

$$-U(x_0) = \left(n + \frac{1}{2}\right) r_H \left[-\frac{1}{2}\frac{\partial^2 U(x_0)}{\partial x_0^2}\right]^{1/2}, \qquad n = 0, 1 \ldots \tag{A.4}$$

The dimension Δx of localization of the solution is of the order $\sqrt{n a r_H}$. Therefore, the assumption $\Delta x \ll a$ that we have made is satisfied for $n \ll a/r_H$. It is easy to show that for $k_y \gg n/(r_H a)^{1/2}$, Eq. (A.4) coincides with the corresponding special case of Eq. (6) if we consider it at the point x_0, where $\partial U/\partial x = 0$.

We are grateful to V. P. Silin for discussions that stimulated us to write this Appendix.

[1)]In the case of the potential (A.3), the point x_0 lies on the real axis of the complex plane x. In the general case, x_0 is a complex number. However, if $\gamma/\omega_r \ll 1$, then account of the small imaginary part of x_0, $\operatorname{Im} x_0/\operatorname{Re} x_0 \sim \gamma/\omega_r$ would mean account of the insignificant terms of the order of $(\gamma/\omega_r)^2$. Therefore, for $\gamma/\omega_r \ll 1$, the point x_0 can be determined by the condition $\partial \operatorname{Re} U/\partial x = 0$.

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