ОБЪЕДИНЕННЫЙ
ИНСТИТУТ
ЯДЕРНЫХ
ИССЛЕДОВАНИЙ

Дубна

E4 - 6716

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ANOMALOUS HEATING OF DENSE PLASMA BY LASER RADIATION

ЛАБОРАТОРИЯ ВЫЧИСЛИТЕЛЬНОЙ ТЕХНИКИ
И АВТОМАТИЗАЦИИ

1972

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Invited paper presented at the Symposium on Plasma Heating and Injection organized by International School of Plasma Physics. Varenna, Italy, 1972.

E4 - 6716

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Аномальный нагрев плотной плазмы лазерным излучением

В докладе обсуждаются вопросы аномального нагрева плазмы лазерным излучением. Проведена систематизация порогов возникающих при этом неустойчивостей и сравнивается их вклад в аномальную диссипацию. Особое внимание уделяется плотной плазме, в которой весьма существенны кулоновские соударения частиц.

Препринт Объединенного института ядерных исследований.
Дубна, 1972

E4 - 6716

Makhankov V.G., Tsytovich V.N.

Anomalous Heating of Dense Plasma by Laser Radiation

The enhanced heating of plasma by laser radiation is dicussed in the report. The systematization of thresholds of instabilities arised is given and contributions of various instabilities into the anomalous dissipation are compared. Dense plasma, in which Coulomb collisions being essential, is paid especial attention.

Preprint. Joint Institute for Nuclear Research.
Dubna, 1972

1. There are at least two ways of the heating of plasma by intense electro-magnetic waves and in particular by a laser radiation:

the usual Joule heating, which presents to be a heating due to the inverse bremsstrahlung induced by a laser radiation ;

the anomalous heating occuring because of the generation of collective plasma oscillations and their absorption. If the first one has been investigated quite sufficiently, the last wanted to be paid more attention for the anomalous heating namely is more natural and perspective from the point of plasma physics view[1,2]. It should be mentioned first that we mean the anomalous heating that occuring due to collective processes and exceeding the Joule heating essentially.Here on the basis of the results had derived in[3-8] the analysis of possibilities of the anomalous heating of plasma and dense plasma especially by intense radiation is given and there are determined the ranges of temperatures and densities of plasma in which such heating is sufficiently effective.The purpose of our work is to give also the review and the systematization of thresholds of the radiation intensity above which there arise instabilities and the excitation of collective plasma oscillations and consequently a possibility of the anomalous heating. We have considered the generation of these oscillation with their frequencies both less and greater than the binary collision frequency. The last case is especially essential for dense plasma. Finally the systematization of the effective frequencies corresponding to anomalous processes of the heating is given as functions of incident energy, temperature and density of plasma.

In the present review the dimension of volume of plasma heating is supposed to be much greater than the wavelength of laser radiation only. The opposite limit of an interaction of H.F. fields with plasma has been investigated in detail in literature(for example, so-called parametric instabilities/9,10/). This analysis and the simple qualitative argumants show that the results of works/9,10/ are not longer valid for the case of small wavelengths because they yield wrong conclusions (for example an independence of the excitation of waves on the angular divergence of a radiation).

It is worth to proceed further to give here the definition of the effective frequency describing the anomalous heating process.

Let Q be a capacity of the generation of collective oscillations by the laser radiation of intensity $I = Wc$, where W is a density of radiation energy. Later instead of W we use a dimensionless quantity $u_0 = W/n_0 T_e$. Then the rate of the heating is determined by the relation

$$\gamma_{eff} = \frac{Q}{W} , \tag{1}$$

the rate of the Joule dissipation respectively to be

$$\gamma_J = \nu_e \left(\frac{\omega_{pe}}{\omega_0} \right)^2 \tag{2}$$

ω_0 is the frequency of the laser radiation, $\omega_{pe}^2 = \frac{4\pi n_0 e^2}{m_e}$

Comparing (1) and (2) we can define the effective frequency of the anomalous heating

$$\nu_{eff} = \frac{\omega_0^2}{\omega_{pe}^2} \frac{Q}{W} \tag{3}$$

The quantity Q can be calculated if the growth rate of the excitation of collective oscillations γ_k as a function of I and wavelength is known for the amplitudes of oscillations are determined by the nonlinear saturation. If a distribution of the oscillation energy over wavelengths is charactezized by the dimensionless quantity w_k and $\int w_k \, dk$ corresponds to the ratio of the total oscillation energy to the density of thermal energy of plasma particles, then

$$Q = n_0 T_e \int \gamma_k \, w_k \, dk \tag{4}$$

therefore

$$\nu_{eff} = \frac{\omega_0^2}{\omega_{pe}^2} \frac{1}{u_0} \int \gamma_k \, w_k \, dk \quad . \tag{5}$$

The growth rates γ_k have been derived in a series of works done by Soviet, Italian and American physicists. Their results will be utilized for the further analysis.

2. As was already mentioned we consider the collective oscillations which wavelengths are less than geometrical dimensions of plasma. Therefore the boundary and the skin-layer effects were ignored. Under such conditions the main

mechamism of the instability of a laser beam is so-called decay-instability[11] and the aperiodic instability of a second-sound type with the imaginary phase velocity[6,8,12] to be important at sufficiently great intensity of laser (last process we have named as quasidecay).

We assume that the transverse wave beam of high-frequency $\omega_o > \omega_{pe}$ with angular divergence $\Delta\vartheta_o$ goes through plasma. Then the principal types of decay are as follows: t is for laser, ℓ for Langmuir, S for ion-sound $(\omega > \nu_i)$ and a for acoustic $(\omega < \nu_i)$ waves respectively

$$t \longrightarrow t' + s$$

$$t \longrightarrow t' + \ell$$

$$t \longrightarrow t' + \ell \longrightarrow t' + \ell' + s \text{ is cascade process}$$

$$t \longrightarrow t' + a$$

and the quasidecays $t \to t' + \tilde{s}$, $t \to t' + \tilde{a}$ where $\tilde{s}$ and $\tilde{a}$ correspond symbolically to aperiodic growing potential perturbations at respectively $\omega > \nu_i$ and $\omega < \nu_i$.

Becides that there is another process

$$t \to \ell + s$$

when the frequency of the laser radiation approaches the Langmuir frequency $\omega_o \to \omega_{pe}$. Let us consider the first type of a decay in more detail. The remaining types can easily be studied dealing analogously. At first we shall analyse the question about thresholds of the intensity above which the ion-sound decay instability is possible and threshold of the transformation of it into that of the quasidecay type.

Using the results of work [6] we have for the grouth rate of the ion-sound instability inspired by the transverse wave beam

$$\gamma_k = \sqrt{\frac{\pi}{8}}\left(\frac{\omega_{pe}}{\omega_o}\right)^3 \frac{\Delta\vartheta_*}{\Delta\vartheta_o} \frac{\omega_{pe}}{(1+k^2 d_e^2)^{3/2}} u_o \begin{cases} 1, & 2k_o \geq k \geq k_\Delta \\ \frac{k}{k_\Delta}, & k \leq k_\Delta \end{cases} . \qquad (6)$$

Here and throughout the work $\Delta\vartheta_* = \frac{v_s}{c}$ is a critical angular divergence, $k_\Delta = k_o \Delta\vartheta_o$ is a critical wave length of ion-sound oscillations excited and $d_e = v_e/\omega_{pe}$, $\omega_o \simeq k_o c$, $v_e = \sqrt{T_e/m_e}$, $v_s = \sqrt{T_e/m_i}$.

The formula (6) is valid at $\Delta\vartheta_o > \Delta\vartheta_* \sqrt{m_e/m_i}$ and $\gamma < k v_s$.

Taking into account the electron Landau damping of S-oscillation $\gamma_s = -\left(\frac{\pi m_e}{8 m_i}\right)^{1/2} \frac{k v_s}{(1+k^2 d_e^2)^{3/2}}$ we get that only the waves with their total growth rates being positive are unstable

$$\gamma_{total} = \gamma_k + \gamma_s > 0 .$$

From this we have the threshold

$$u_1 = u_{th} \begin{cases} \frac{k}{k_\Delta}, & k \geq k_\Delta \\ 1, & k \leq k_\Delta \end{cases} . \qquad (7)$$

Where

$$u_{th} = \mu^{3/2}\left(\frac{v_e}{c}\right)^2 \left(\frac{\omega_o}{\omega_{pe}}\right)^4 \left(\frac{\Delta\vartheta_o}{\Delta\vartheta_*}\right)^2 \qquad (8)$$

$$\mu = m_e/m_i .$$

In this notation

$$\gamma_k = \gamma_s \frac{u_o}{u_1} \simeq \omega_s \frac{u_o}{u_1} \mu^{1/2} \qquad (9)$$

The conditions $|\gamma_s| < \gamma_k < k v_s$ lead to $\mu^{-1/2} u_1 > u_0 > u_1$.
The quantity $u_2 := \mu^{-1/2} u_1$ is the threshold of the transformation of S -instability into that of the quasidecay type, at $u_0 > u_2$ new aperiodic solutions of the second-sound type appear (see work/6/). More detail analysis of the growth rates of quasidecays gives

$$\tilde{\gamma}_k = \omega_s \left(\frac{u_0}{u_2}\right)^{1/2} , \quad at \quad \left(\frac{\Delta\vartheta_0}{\Delta\vartheta_*}\right)^2 u_2 > u_0 > u_2$$

and

$$\gamma_k = \omega_s \begin{cases} \left(\dfrac{\Delta\vartheta_0 u_0}{\Delta\vartheta_* u_2}\right)^{1/3} , & k \gg k_\Delta \\ \left(\dfrac{\Delta\vartheta_0^2 u_0}{\Delta\vartheta_*^2 u_2}\right)^{1/4} , & k < k_\Delta \end{cases} \quad at \quad u_0 > u_2 \left(\frac{\Delta\vartheta_0}{\Delta\vartheta_*}\right)^2 .$$

It has been reported in/6/ that besides the kinetic (decay) instability with the growth rate proportional to there may occur the instabilities of hydrodynamic and dissipative types with the growth rates proportional to a square or cubic root of u_0 .

The S -kinetic instability growth rates γ_k and γ_{total} as the functions of the wave number at various values $u_0 > u_{th}$ are shown in Fig. 1. We can see there that the peak of the growth rate lies at $k = k_\Delta$, therefore S - waves with the wave number $k = k_\Delta$ and the frequency $\omega_\Delta = k_\Delta v_s$ have the greatest amplitudes. The thresholds of various types of S - waves instability as the function of $\Delta\vartheta_0$ are plotted in Fig. 2. At certain u_0 and $\Delta\vartheta_0$ we put into different areas of the instability that characterized by different growth rates and their dependences of u_o .

Figure 2 correspond to a certain fixed k . As the detail trearment shows the boundaries of areas change significantly with k .

Figure 3 depicts the thresholds and the regions of various types of instabilities as the functions of the wave number of oscillations excited at certain $\Delta \vartheta_o > \Delta \vartheta_*$ (see the line $S-S$ in Fig. 2). The above mentioned Fig. 1 shows γ_{total} and γ_k as the function of k along the line I in Fig.2. The dependence of γ_{total} on k along the lines 2,3 and 4 is presented in Fig. 4,5 and 6. The growth rate in Fig. 4,5, has a maximum at $k = k_{trans}$ corresponding to the transformation of the kinetic instability into that of qusidecay type. A slight slope of γ at $k > k_{trans}$ is (as in Fig. 1) associated with the Landau damping. The region of wave numbers in which $\gamma \approx \gamma_{max}$ is significantly greater to the right of k_{trans} than to the left: therefore the kinetic instability is to be the main in cases 2 and 3. In case 4 only the quasi-decay instability remains. Finally the growth rate as the function of u_o at certain k (section $\alpha-\alpha$ in Fig. 3) is presented in Fig. 7.

The most interesting range of angular divergence of a laser beam at $k_o c \gtrsim \omega_{pe}$ is

$$\Delta \vartheta_o > \Delta \vartheta_* \approx \mu^{1/2} \frac{v_e}{c} \approx 10^{-3}$$

Let us analogous to Fig. 2 draw the picture of the distribution of the areas of different types of the instabilities in variables u_o and $\Delta \vartheta_o$. We qualify now as the kinetic area that where except the kinetic instability may occur quasidecay instabilities but with growth rate being less than the maximum of the kinetic growth rate in the same area see Fig.4 and 5 (analogous for the quasidecay instabilities with $\gamma \sim u_o^{1/2}$ and $\gamma \sim u_o^{1/3}$). We get then Fig. 8.

We can now see that the kinetic instability area is significantly broader than that shown in Fig. 2 and above mentioned values $\Delta \vartheta_o > 10^{-3}$ lie in it.

These results allow us to calculate the effective rate of the anomalous heating associated with the generation and the absorption of ion-sound osillations. Thus if $u_o > u_{th}$ the kinetic instability arizes that leads to a rapid (exponential) increase of the s-wave energy. In time scale of order $\tau \approx \frac{1}{\gamma_\Delta} \approx \frac{u_{th}}{\mu^{1/2}\omega_\Delta u_o}$ the s-oscillation amplitude grows very much and the nonlinear interaction of s-waves (such as induced scattering of them by plasma ions with a growth rate γ_{sc} or "correlation" s-s decay/13/ with a growth rate of order $8\frac{T_e}{T_i}\gamma_{sc}$)come into the game. As a result of this interaction the s-wave energy goes away from the region of the generation in angle as well as in wave number. At a certain level of the s-wave energy the rates of these two processes (i.e. the process of the generation s-waves by a laser beam and the transformation of their energy due to induced scattering on plasma ions) rush to be equal in the generation region.Therefore we have a steady flux of the laser radiation energy transforming into the s-oscillation energy with its furthe dissipation by plasma particles. The condition $\gamma_k \simeq \gamma_{sc}$ gives the form of the quasisteady s- wave spectrum in the generation region. Using (5) we get the effective rate of the anomalous heating ν_{eff} The equation $\gamma = \gamma_{sc}$ yields

$$\frac{d}{dk} k w_k^s = -\frac{T_e}{T_i} \frac{\Delta\vartheta_*}{\Delta\vartheta_0} \left(\frac{\omega_{pe}}{\omega_0}\right)^3 \frac{\omega_{pe}}{\omega_s} \frac{u_0}{k} \qquad (10)$$

Integrating over k from k to k_{max} we find

$$w_k^s \sim \frac{1}{k}\left(\frac{1}{k} - \frac{1}{k_{max}}\right).$$

The width of this spectrum $\Delta k \approx k_\Delta$. Now we get substituting w_k^s into (5)

$$\nu_{eff} = \left(\frac{\omega_0}{\omega_{pe}}\right)^2 \frac{\gamma^2}{\omega_s} \frac{T_e}{T_i} \frac{1}{u_0} \left(\frac{\Delta k}{k}\right)^2 \approx \gamma_k \frac{1}{(\Delta\vartheta_0)^2} \left(\frac{\omega_{pe}}{\omega_0}\right)^2 \quad \text{(for s-waves)}$$

or

$$\nu_{eff} = \mu^{1/2} \omega_{pe} \frac{\Delta\vartheta_*}{\Delta\vartheta_0} \frac{\omega_{pe}}{\omega_0} \frac{T_e}{T_i} \frac{u_0}{u_{th}} \begin{cases} \left(\frac{u_0}{u_{th}} - 1\right)^2, & \text{at } u_0 \to u_{th} \\ 1 & \text{at } u_0 > u_{th} \end{cases} \qquad (11)$$

where $\mu^{-1/2} u_{th} > u_0 > u_{th}$.

If $u_0 > \mu^{-1/2} u_{th}$ the growth rate goes to be proportional to $u_0^{1/2}$ and $\int w_k^s \, dk = w_0^s \to 1$ so (11) is not more valid . A rude estimation $w_0^s \simeq 1$ at $u_0 > \mu^{-1/2} u_0$ leads us to the result that $\nu_{eff} \sim u_0^{-1/2}$ and decreases with increasing of u_0 , therefore ν_{eff} has a peak at $u_0 = u_{th} \mu^{-1/2}$

$$\nu_{eff}\,(max) \simeq \omega_{pe} \frac{\Delta\vartheta_*}{\Delta\vartheta_0} \frac{\omega_{pe}}{\omega_0} \frac{T_e}{T_i} \qquad (12)$$

when

$$\frac{u_0}{u_{th}} > \mu^{-1/2} (\Delta\vartheta_0)^2 \left(\frac{v_e}{c}\right)^2 \left(\frac{\omega_0}{\omega_{pe}}\right)^2$$

the so-called "correlation decay" is now the main interaction of s-waves and ν_{eff} decreases with T_e/T_i factor.

These estimations for ν_{eff} are based on the equation $\gamma_k = \gamma_{sc}$, but it can be possible if the broadening of the angular divergence of a radiation takes time longer than the nonlinear interaction of s-waves time. In the opposite case the increase of the angular divergence concerning the scattering of t-waves leads to a decrease of the linear growth rate. This decrease of γ_k at small $\Delta\vartheta_0$ (initial) can be significant up to the threshold γ_s. The instability comes then to an end because of the angular relaxation. There is possible the situation when $\gamma_{total} > 0$ even at $\Delta\vartheta_0 \approx 1$ or more correctly at $\Delta\vartheta_0 \gtrsim \Delta\vartheta_{relax}$. Estimations give

$$\Delta\vartheta_{relax} \approx \left(\frac{\omega_{pe}}{\omega_0}\right)^2 . \qquad (13)$$

If $\Delta\vartheta_0 < \Delta\vartheta_{relax}$ then at the initial stage ν_{eff} is determined by the linear growth rate γ_k i.e.

$$\nu_{eff} \simeq \gamma_k \left(\frac{\omega_0}{\omega_{pe}}\right)^2 . \qquad (14)$$

In time of the order $1/\gamma_k$ the divergence $\Delta\vartheta_0$ reaches $\Delta\vartheta_{relax}$ and ν_{eff} is described again by (11) with $\Delta\vartheta_0 \approx \Delta\vartheta_{relax}$

$$\nu_{eff} \simeq \gamma_k (\Delta\vartheta_{relax}) \left(\frac{\omega_0}{\omega_{pe}}\right)^2 = \mathcal{U}_0 \frac{v_s}{c}\,\omega_0 \qquad (15)$$

and

$$u_{th} = \mu^{1/2} . \qquad (16)$$

We can derive from (12) that $\nu_{eff}\,(max) = \omega_o v_s / c$ at $u_o \approx 1$. When $\Delta\vartheta_o > \Delta\vartheta_{relax}$, ν_{eff} and u_{th} are determined by (11) and (8) and the change of $\Delta\vartheta_o$ during the interaction turns out to be small.

The anomalous heating takes place only when $\nu_{eff} > \nu_e$ therefore we have for $N_D \approx \omega_{pe}/\nu_e$ (the number of particles in Debye radius sphere) at the angular relaxation possible to be neglected

$$N_D > N_D^s \equiv \frac{\Delta\vartheta_o}{\Delta\vartheta_*} \frac{\omega_o}{\omega_{pe}} \mu^{-1/2} \frac{u_{th}}{u_o} \qquad (17)$$

and

$$N_D^s\,(min) = \mu^{-1/2} \Delta\vartheta_o \frac{c}{v_e} \frac{\omega_o}{\omega_{pe}} . \qquad (17a)$$

In the opposite case

$$N_D^s = \mu^{-1/2} \left(\frac{\omega_{pe}}{\omega_o}\right)^3 \frac{1}{\Delta\vartheta_o \Delta\vartheta_*} \frac{u_{th}}{u_o}$$

As it was done above we can find the effective frequencies corresponding to the remaining channels of the decay. Those $t \to t' + \ell$ and $t \to t' + \ell \to t' + \ell' + s$ give approximately the same heating frequencies (the cascade process if it is possible leads to a smaller ν_{eff}) as for the first one associated with the stabilization of the ℓ -wave instability due to their induced scattering on ions and the second due to $\ell \to \ell' + s$ decay. Utilizing the results of[4] and (5) we get

$$\nu_{eff}^{\ell} \approx \omega_{pe} \left(\frac{\omega_{pe}}{\omega_o}\right)^4 \left(\frac{v_e}{c}\right)^6 \frac{m_i}{m_e} \frac{u_o}{(\Delta\vartheta_o)^2} \tag{18}$$

or taking into account the threshold u_{th}^{ℓ} from $\gamma = \nu_e$,
$u_{th}^{\ell} = \frac{1}{N_D}\left(\frac{\omega_o}{\omega_{pe}}\right)^3 \Delta\vartheta_o^2 \frac{c^2}{v_e^2}$ we have

$$\nu_{eff}^{\ell} \approx \nu_e \frac{\omega_{pe}}{\omega_o} \left(\frac{v_e}{c}\right)^4 \frac{m_i}{m_e} \frac{u_o}{u_{th}} \tag{19}$$

and

$$\nu_{eff}^{\ell}(max) \approx \frac{\omega_{pe}}{\omega_o} \left(\frac{v_e}{c}\right)^4 \frac{m_i}{m_e} . \tag{20}$$

The condition necessary for the dissipation due to process to be enohanced is

$$N_D >> \left(\frac{c}{v_e}\right)^6 \frac{m_e}{m_i} \left(\frac{\omega_o}{\omega_{pe}}\right)^4 \frac{(\Delta\vartheta_o)^2}{u_o} \equiv N_D^{\ell} \tag{21}$$

and when ν_{eff}^{ℓ} being maximum

$$N_D^{\ell}(min) = \frac{\omega_o}{\omega_{pe}} \left(\frac{c}{v_e}\right)^4 \frac{m_e}{m_i} . \tag{22}$$

Comparing (17a)and (22) one can find

$$\frac{N_D^s\,(min)}{N_D^{\ell}\,(min)} \approx \Delta\vartheta_o \left(\frac{v_e}{c}\right)^3 \mu^{-3/2} .$$

This means that the process $t \to t' + s$ is more effective at $v_e/c < \mu^{1/2}(\omega_o/\omega_{pe})^{2/3}(if\ \Delta\vartheta_o \approx \omega_{pe}^2/\omega_o^2)$. The processes of the anomalous dissipation studied above took place when $\omega_o > \omega_{pe}$ and $\omega_o > 2\omega_{pe}$ for $t \to t' + \ell$ channel. At

$\omega_o \to \omega_{pe}$ the process $t \to t' + \ell$ is impossible, the $t \to \ell + s$ decay goes to be the main and at the same time the efficiency of the $t \to t' + s$ process enlarges.

The anomalous dissipation frequency due to the $t \to \ell + s$ decay can be easily derived via the growth rate with allowance for the energy densities of t and ℓ-waves rush to be equivalent. Therefore

$$\nu_{eff}^{\ell_s} = \frac{1}{9} \frac{m_e}{m_i} U_o \, \omega_{pe} \; . \tag{23}$$

The threshold of this instability is followed from the condition $\nu_{eff}^{\ell_s} > \nu_e$ therefore $U_{th}^{\ell_s} = 9 \frac{m_i}{m_e} \frac{1}{N_D}$ and

$$\nu_{eff}^{\ell_s} = \nu_e \frac{U_o}{U_{th}} \; . \tag{24}$$

Formula (11) for $t \to t' + s$ channel in the resonance region $\omega_o \to \omega_{pe}$ is formally the same if

$$\Delta \vartheta_{*r} = \Delta \vartheta_* \frac{\omega_{pe}}{k_o c} < 1$$

being substituted in lieu of $\Delta \vartheta_*$. Here one must keep in mind that $\omega_{pe}/k_o c < c/v_s$ therefore

$$\nu_{eff}^{s} \, (res) = U_o \frac{\Delta \vartheta_*}{\Delta \vartheta_o} \frac{\omega_{pe}}{k_o c} \omega_{pe} < \omega_{pe} \frac{U_o}{\Delta \vartheta_o} \; . \tag{25}$$

We can now find from comparing (25) and (17) that at the same U_o $\nu_{eff}^{\ell_s} = \nu_{eff}^{s} \, (res)$ when

$$\Delta \vartheta_o \simeq \mu^{-1/2} \frac{\omega_{pe}}{k_o c} \frac{v_e}{c} = \Delta \vartheta_1 \; . \tag{26}$$

In the resonance region the $t \to t' + s$ process is also important at $\Delta \vartheta_o < \Delta \vartheta_1$. The generation of s-waves followed by the simultaneous increase of the angular divergence continues untill $\Delta \vartheta_0 \approx \Delta \vartheta_1$, if $\Delta \vartheta_1 \ll 1$ and then the $t \to \ell + s$ process becomes the main.

3. The above results are valid for quite rarefied and sufficiently hot plasma. Let us see how the interaction of a laser beam with plasma changes in the case of the influence of the pair particle collisions to be essential (if not determining), i.e. in dense plasma. The dispersion curve $\omega(k)$ and the damping rate $\gamma(k)$ for low frequency (L.F.) oscillations are plotted in Fig. 8 (see /7/). One can see that in a plasma with temperatures of electrons to be greater than that of ions , $T_e \gtrsim 3T_i$, the ion-sound branch $\omega > \nu_i$ transforms into acoustic that $\omega < \nu_i$ at $k = k_\nu \equiv \nu_e / v_e$. The damping rate $\gamma(k) \approx \mu^{1/2} \omega(k)$ decreases with decreasing k for $k > k_\nu$ and becomes constant, and equal to $\mu^{1/2} \nu_i$ if $k_\nu > k > \mu^{1/2} k_\nu$.

In the wave number region $k \approx \mu^{1/2} k_\nu$ the damping rate is of order of frequency (the region of a significant absorption of a -ossillations), then it decreases with decreasing k as follows

$$\gamma(k) \approx \frac{k^2 v_e^2}{\nu_e} = \frac{k}{\mu^{1/2} k_\nu} \omega(k)$$

As was shown in /7/ the decay probability increases approximately with $(k_\nu / k)^2$ factor when k becomes less than k_ν . Curve III in Fig.8 depicts the growth rate of L.F. oscillations due to the presence of a laser radiation in plasma.

The influence of pair collisions on the generation of L.F. oscillation and consequently on the anomalous dissipation connected with this is generally speaking miserable if k_ν much less than k_Δ where the growth rate of the instability has a maximum i.e.

$$k_\Delta \equiv k_o \Delta Q_o >> k_\nu \equiv \nu_e / v_e$$

or

$$N_D >> N_D(\Delta) \equiv \frac{\omega_{pe}}{\omega_o} \frac{c}{v_e} \frac{1}{\Delta \theta_o} .$$

On the other hand one can see that the growth rate increases with decreasing k being less than ν_i if $k < k_\nu$. This is illustrated in Fig. 9,where thresholds u_1 and u_2 with allowance for the generation of acoustio oscillations are pictured.It is easily seem that ion-sound oscillations are not being excited and acoustic waves become unstable in the region $k_{min} \leq k \leq k_{max}$ when $u_o < u_{th}$. Although the growth rates of a-waves less than ν_i their effect can be quite essential if an interaction of a laser with plasma continues long enough (longer than time of accumulation of a-oscillations in the interaction region, in other words longer than the inverse growth rate of a-wave interaction). Therefore the thresholds of excitation of a-waves are significantly less than those for s-waves.

We shall give the frequencies of the anomalous heating connected with the generation of acoustic waves and range of plasma densities in which these processes are essential.

The influence of pair collisions of plasma particles can in fact be important at yet $k_\Delta \gtrsim k_\nu$. If the impulse of laser radiation is sufficiently short for $\tau_{rad} < 1/\gamma_k$ the accumulation of L.F. oscillations does not occur and the dissipation of laser energy to be of Joule type. This takes place at $\gamma_J > \gamma_k$ and $1/\gamma_k > \tau_{rad} > 1/\gamma_J$.

When $\gamma_J > \gamma_k$ and the impulse of a laser is long the angular relaxation of the laser beam occurs and ν_{eff} decreases to be less or of order γ_k in about time-scale $1/\gamma_k$. If $\gamma_J < \gamma_k$ or

$$N_D > N_{D1}(\gamma) \equiv \mu^{-3/2} \left(\frac{c}{v_e}\right)^2 \frac{\Delta\vartheta_*}{\Delta\vartheta_0} \left(\frac{\omega_{pe}}{\omega_0}\right)^3 \tag{28}$$

the laser flash energy can be anomalously dissipated.

When

$$N_{D1}(\gamma) > N_D > \mu^{1/2} N_{D1}(\gamma) \tag{29}$$

the condition $\gamma_k > \gamma_J$ yields

$$u_0 > u_{th}(\nu) \equiv u_{th} \frac{N_{D1}(\gamma)}{N_D} . \tag{30}$$

When (27), (28) and (30) hold the effective frequency of absorption is determined by (11).

Finally the dissipation becomes that of Joule type at

$$u_{th}(\nu) > u_0 > u_{th} . \tag{31}$$

Let us now study the range of plasma parameters necessary for the effect of Coulomb collisions to be decisive. It is clear this occurs at $k_\Delta < k_\nu$ or

$$N_D < \frac{\omega_{pe}}{\omega_0} \frac{c}{v_e} \frac{1}{\Delta\vartheta_0} . \tag{32}$$

Suppose $k_\Delta \approx k_\nu$ and $u_0 > u_{th}$. Then as it follows from Fig. 9

$$\gamma_k^{(s)}(max) = \gamma \equiv \mu^{1/2}\omega_\Delta \frac{u_0}{u_{th}} \quad , \quad \gamma_k^{(a)}(max) = \gamma \frac{k_\nu^2}{k_\Delta k_{min}}$$

and k_{min} is given by equation $\gamma_k^{(a)}(max) = k v_s$ therefore

$$k_{min}^2 = k_\nu^2 \mu^{1/2} \frac{u_0}{u_{th}} \quad and \quad \gamma_k^{(a)}(max) = \gamma \frac{k_\nu}{k_\Delta} \mu^{-1/4} \sqrt{\frac{u_{th}}{u_0}} .$$

By comparing $\gamma_k^{(s)}(max)$ and $\gamma_k^{a}(max)$ we find at

$$\frac{k_\nu}{k_\Delta} > \mu^{1/4} \sqrt{\frac{u_0}{u_{th}}} \tag{33}$$

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When the level of laser beam energy is sufficient to lay $\Delta \vartheta_0 \approx (\omega_{pe}/\omega_0)^2$ in u_{th} the anomalous dissipation due to acoustic instability appears at a higher N_D

$$N_D < \mu^{-3/4} \left(\frac{u_{th}}{u_0}\right)^{3/2} \frac{\omega_{pe}}{\omega_0} \frac{c}{v_e} \Delta \vartheta_0 . \tag{35}$$

Suppose now u_0 to be less than u_{th} there remains then only the a-wave instability. In this case and when (34), (35) hold we have for ν_{eff}

$$\nu_{eff}^{(a)} = \nu_{eff}^{(s)} \mu^{-3/4} \frac{k_\nu}{k_\Delta} \left(\frac{u_{th}}{u_0}\right)^{3/2} < \nu_i . \tag{36}$$

The effective dissipation frequency due to the excitation of L.F. fluctuations is a sharp decreasing function of laser radiation frequency. Moreover, the L.F. instability threshold increases with inoreasing ω_0. These facts are illustrated in Figs. 11,12. Figure 11 shows that the instability breaks down when a plasma density decreases up to $\omega_0/\omega_{pe} \geq \alpha_{cr}^{u} \equiv (u_0/u_{th}^{0})^{1/4}$, $u_{th}^{\bullet} = u_{th}$ $(\alpha = 1)$ at fixed u_0 and ω_0.

The characteristic dependence of the anomalous dissipation due to the generation of both s and a-waves as a function of ω_0 is depicted in Fig. 12.

4. The main conclusions of this review are as follows.

(1) In dense plasma the long anomalous dissipation is possible only at $\omega_0/\omega_{pe} > \mu^{-1/4}$.

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(3) The processes following the interaction of a laser beam with plasma used to be of two types: the anomalous dissipation and scattering of radiation. Last effect can be more easily observed.

(4) The utilization of beam with a small angular divergence for the enchanced heating in many cases is unsuitable as the anomalous scattering following heating enlarges rapidly the angular divergence. For the clear observation of the anomalous dissipation the threshold corresponding to a sufficiently large angular divergence should be exceeded.

(5) The thresholds for the enhanced absorption are not large, but an explicit experimental study of this requires that a series of above conditions and restrictions have to be fulfilled. For example having used formula (16) at $n_0 \simeq 10^{19}$ and $T_e \approx 3_{KeV}$ we get that the electric field tensity of laser radiation has to exceed $E_{th} \approx 3 \cdot 10^{8}$ volt/cm. The rate of the anomalous absorption is then $\nu_{eff} \approx 10^{11} sec^{-1}$, that exceeds the rate of Joule dissipation by two orders. This value of ν_{eff} according to (27) concernes with excitation of ion-sound oscillations.

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References

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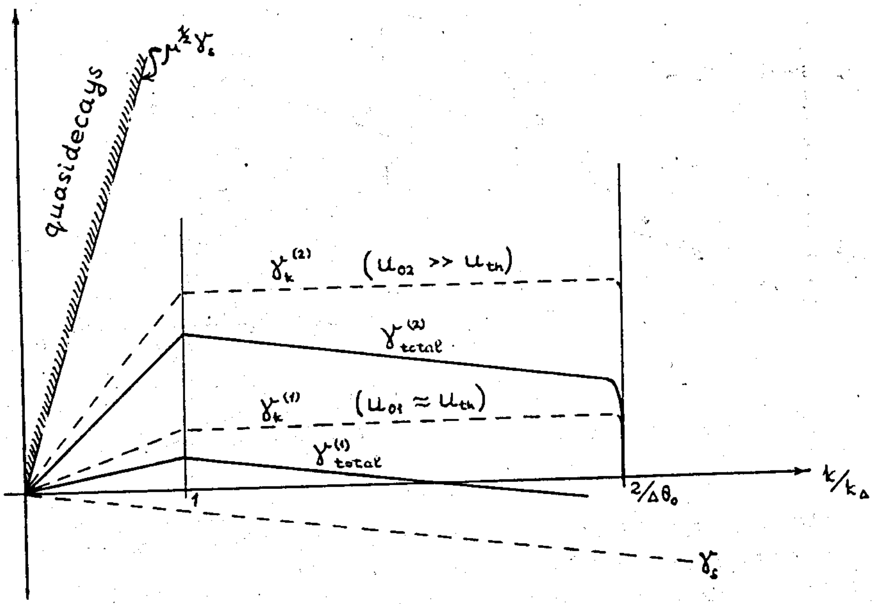

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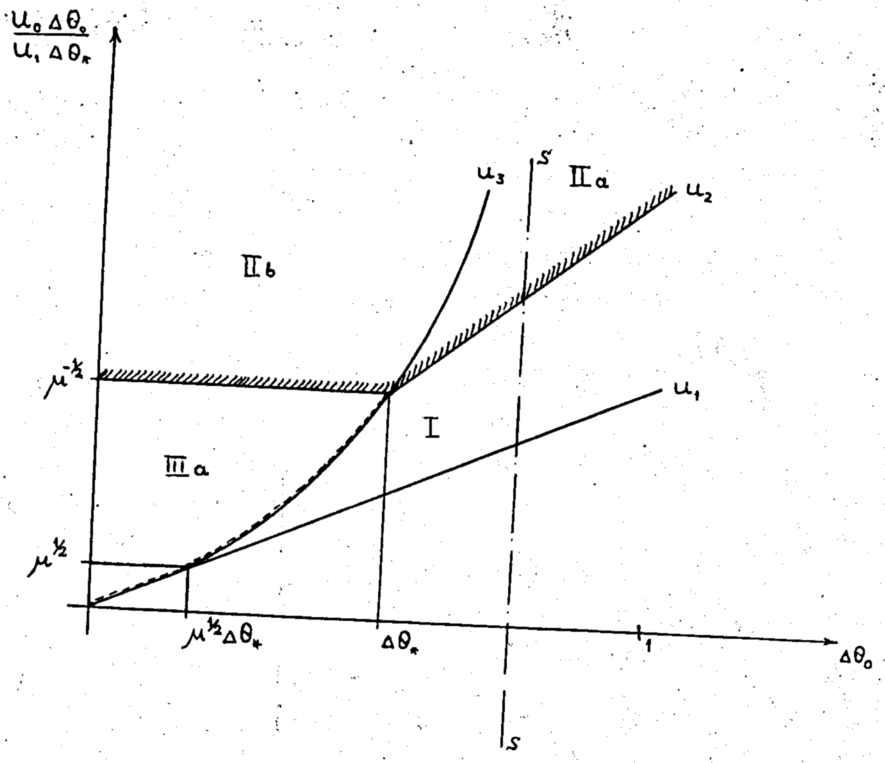

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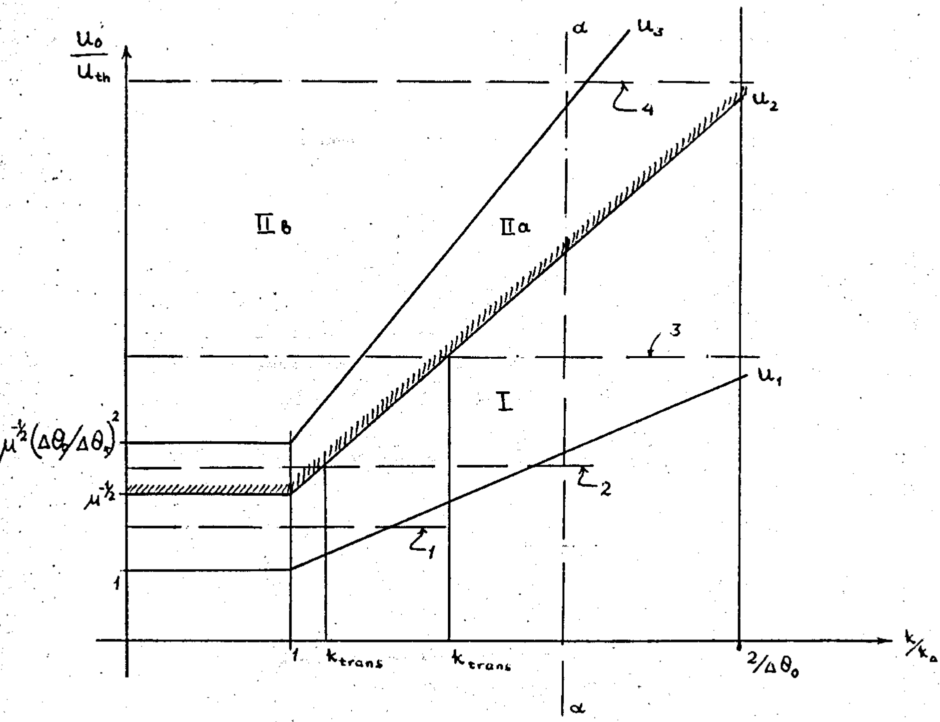
IIIb is an area of dissipative instability $\gamma_k < \omega_s$; $\gamma_k \sim u_0$.

$$u_3 = u_2 \left(\Delta\vartheta_0 / \Delta\vartheta_k \right)^2$$



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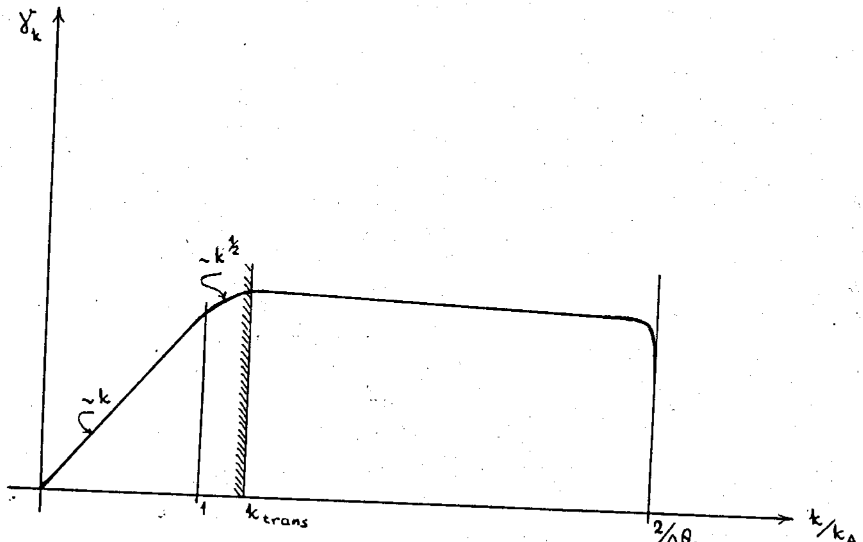

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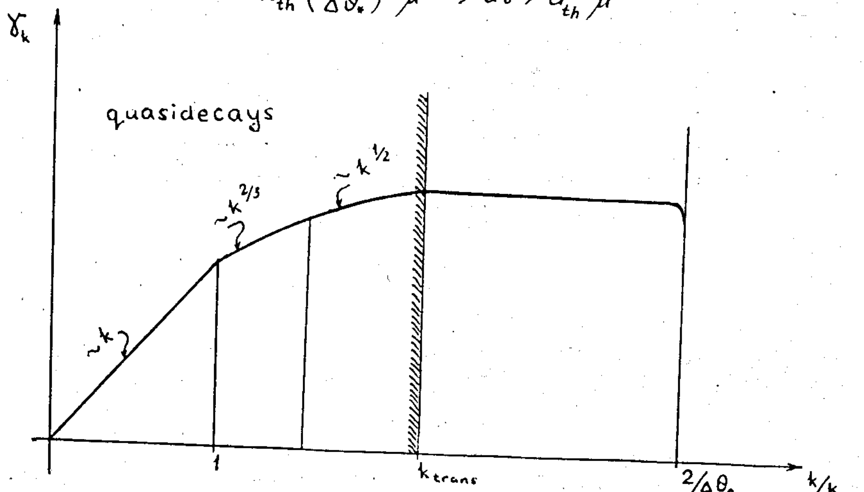
$$U_{th}\left(\frac{\Delta\vartheta_0}{\Delta\vartheta_*}\right)^2\mu^{-1/2} > U_0 > U_{th}\,\mu^{-1/2}$$



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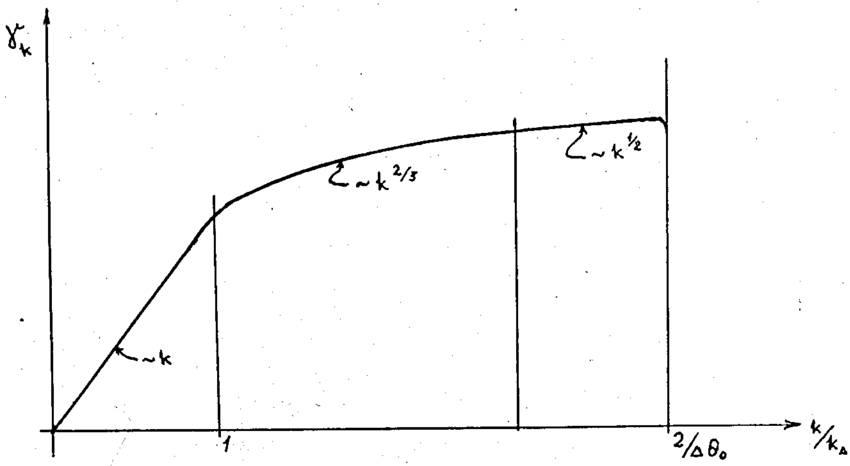
$$U_0 > U_{th}\left(\frac{\Delta\vartheta_0}{\Delta\vartheta_*}\right)^2\mu^{-1/2}$$



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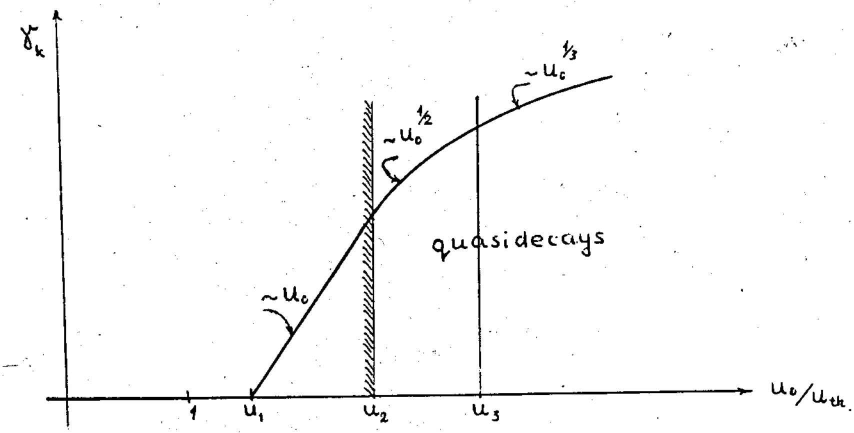

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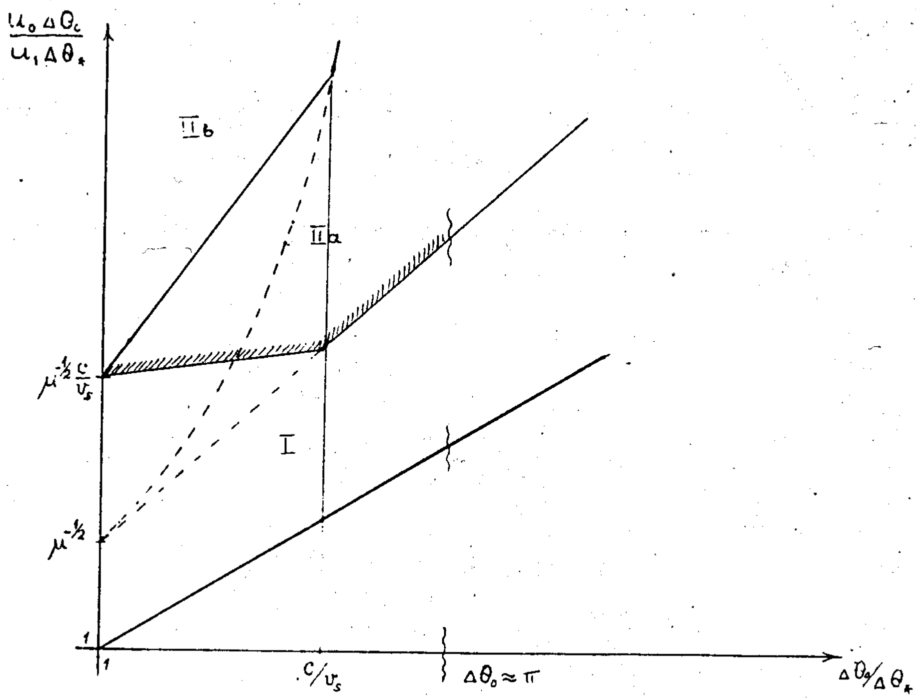

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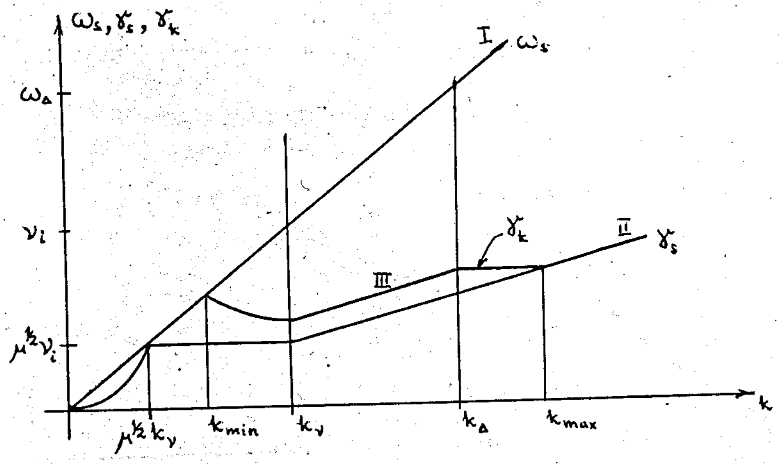

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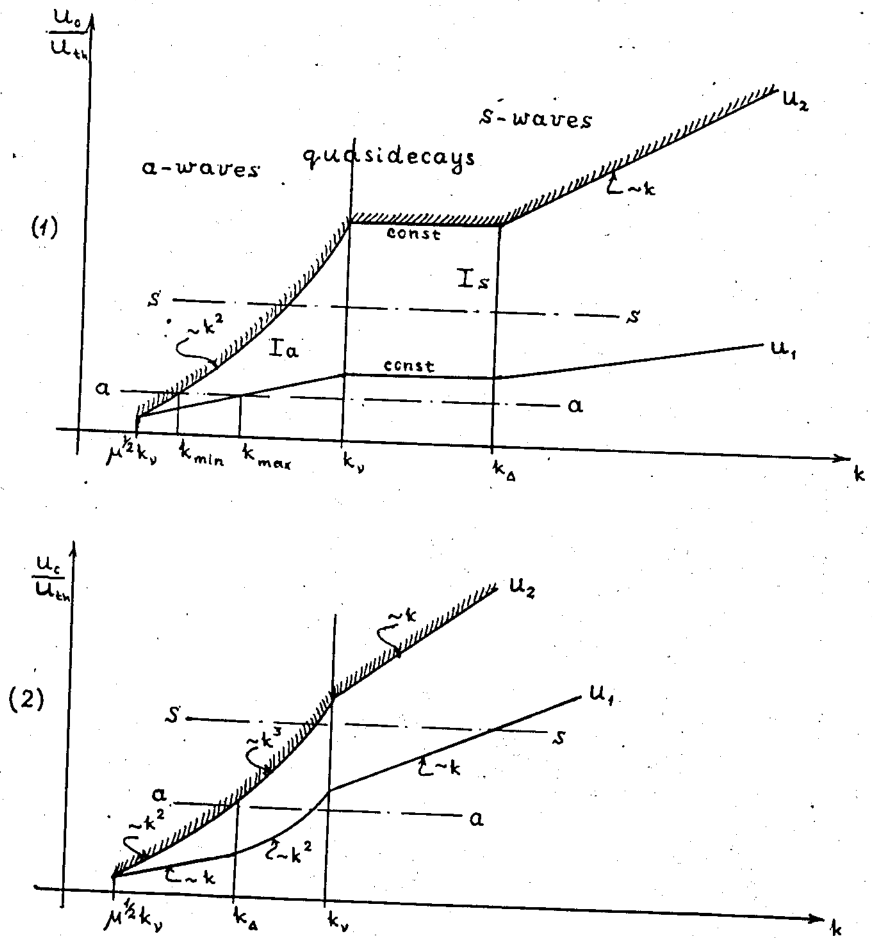

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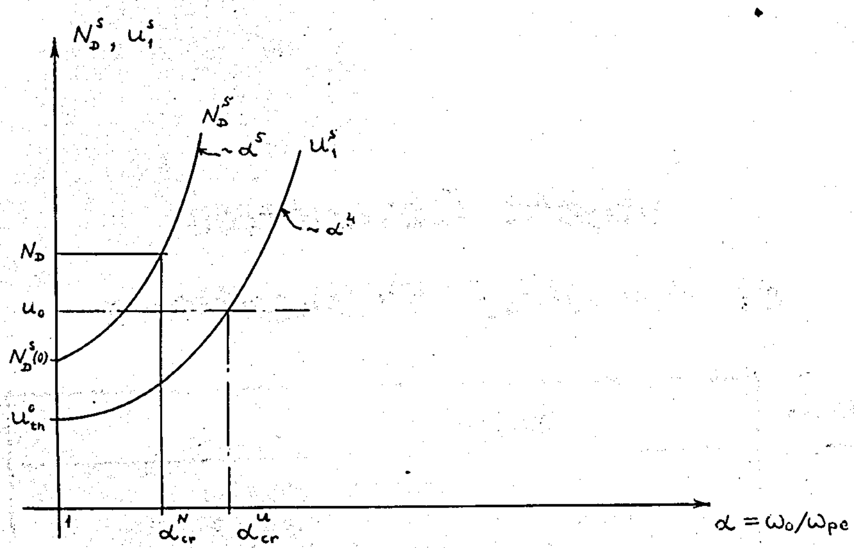

Fig. 11. Dependence of the kinetic S -instability threshold and the critical Debay number N_D^S (17) on $\alpha = \omega_0/\omega_{pe}$:

$$\alpha_{cr}^u = (u_0/u_{th}^0)^{1/4}, \quad u_{th}^0 = u_{th} \quad \omega_0 = \omega_{pe}$$

$$\alpha_{cr}^N = (N_D/N_D^s(0))^{1/5}, \quad N_D^s(0) = N_D^s \quad \omega_0 = \omega_{pe}$$

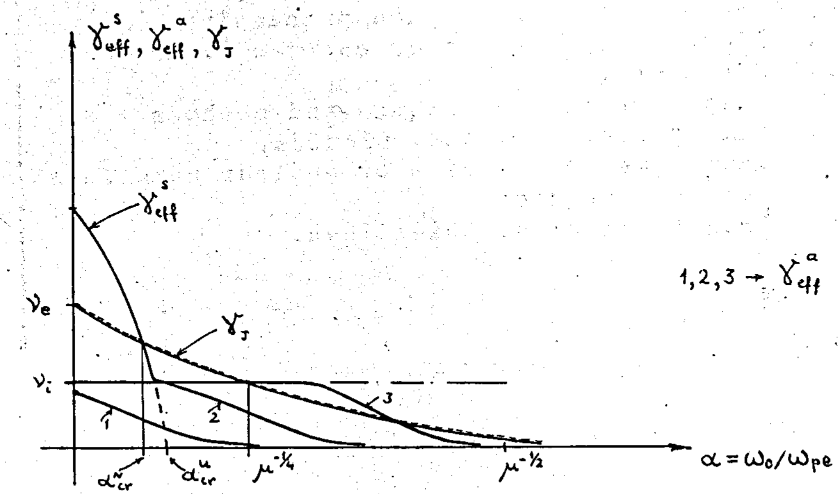

Fig. 12. Dependence of the anomalous heating rate on ω_0/ω_{pe}.

The influence of pair collisions on the generation of L.F. oscillation and consequently on the anomalous dissipation connected with this is generally speaking miserable if k_ν much less than k_Δ where the growth rate of the instability has a maximum i.e.

$$k_\Delta \equiv k_o \Delta \theta_o \gg k_\nu \equiv \nu_e / v_e$$

or

$$N_D \gg N_D(\Delta) \equiv \frac{\omega_{pe}}{\omega_o} \frac{c}{v_e} \frac{1}{\Delta \theta_o} .$$

On the other hand one can see that the growth rate increases with decreasing k being less than ν_i if $k < k_\nu$. This is illustrated in Fig. 9, where thresholds u_1 and u_2 with allowance for the generation of acoustio oscillations are pictured. It is easily seem that ion-sound oscillations are not being excited and acoustic waves become unstable in the region $k_{min} \leq k \leq k_{max}$ when $u_o < u_{th}$. Although the growth rates of a -waves less than ν_i their effect can be quite essential if an interaction of a laser with plasma continues long enough (longer than time of accumulation of a -oscillations in the interaction region, in other words longer than the inverse growth rate of a -wave interaction). Therefore the thresholds of excitation of a-waves are significantly less than those for s -waves.

We shall give the frequencies of the anomalous heating connected with the generation of acoustic waves and range of plasma densities in which these processes are essential.

The influence of pair collisions of plasma particles can in fact be important at yet $k_\Delta \gtrsim k_\nu$. If the impulse of laser radiation is sufficiently short for $\tau_{rad} < 1/\gamma_k$ the accumulation of L.F. oscillations does not occur and the dissipation of laser energy to be of Joule type. This takes place at $\gamma_J > \gamma_k$ and $1/\gamma_k > \tau_{rad} > 1/\gamma_J$.

When $\gamma_J > \gamma_k$ and the impulse of a laser is long the angular relaxation of the laser beam occurs and ν_{eff} decreases to be less or of order γ_k in about time-scale $1/\gamma_k$. If $\gamma_J < \gamma_k$ or

$$N_D > N_{D1}(\gamma) \equiv \mu^{-3/2} \left(\frac{c}{v_e}\right)^2 \frac{\Delta\vartheta_*}{\Delta\vartheta_0} \left(\frac{\omega_{pe}}{\omega_0}\right)^3 \qquad (28)$$

the laser flash energy can be anomalously dissipated.

When

$$N_{D1}(\gamma) > N_D > \mu^{1/2} N_{D1}(\gamma) \qquad (29)$$

the condition $\gamma_k > \gamma_J$ yields

$$u_0 > u_{th}(\nu) \equiv u_{th} \frac{N_{D1}(\gamma)}{N_D} . \qquad (30)$$

When (27), (28) and (30) hold the effective frequency of absorption is determined by (11).

Finally the dissipation becomes that of Joule type at

$$u_{th}(\nu) > u_0 > u_{th} . \qquad (31)$$

Let us now study the range of plasma parameters necessary for the effect of Coulomb collisions to be decisive. It is clear this occurs at $k_\Delta < k_\nu$ or

$$N_D < \frac{\omega_{pe}}{\omega_0} \frac{c}{v_e} \frac{1}{\Delta\vartheta_0} . \qquad (32)$$

Suppose $k_\Delta \approx k_\nu$ and $u_0 > u_{th}$. Then as it follows from Fig. 9

$$\gamma_k^{(s)}(max) = \gamma \equiv \mu^{1/2}\omega_\Delta \frac{u_0}{u_{th}}, \quad \gamma_k^{(a)}(max) = \gamma \frac{k_\nu^2}{k_\Delta k_{min}}$$

and k_{min} is given by equation $\gamma_k^{(a)}(max) = k v_s$ therefore

$$k_{min}^2 = k_\nu^2 \mu^{1/2} \frac{u_0}{u_{th}} \quad and \quad \gamma_k^{(a)}(max) = \gamma \frac{k_\nu}{k_\Delta} \mu^{-1/4} \sqrt{\frac{u_{th}}{u_0}}.$$

By comparing $\gamma_k^{(s)}(max)$ and $\gamma_k^{a}(max)$ we find at

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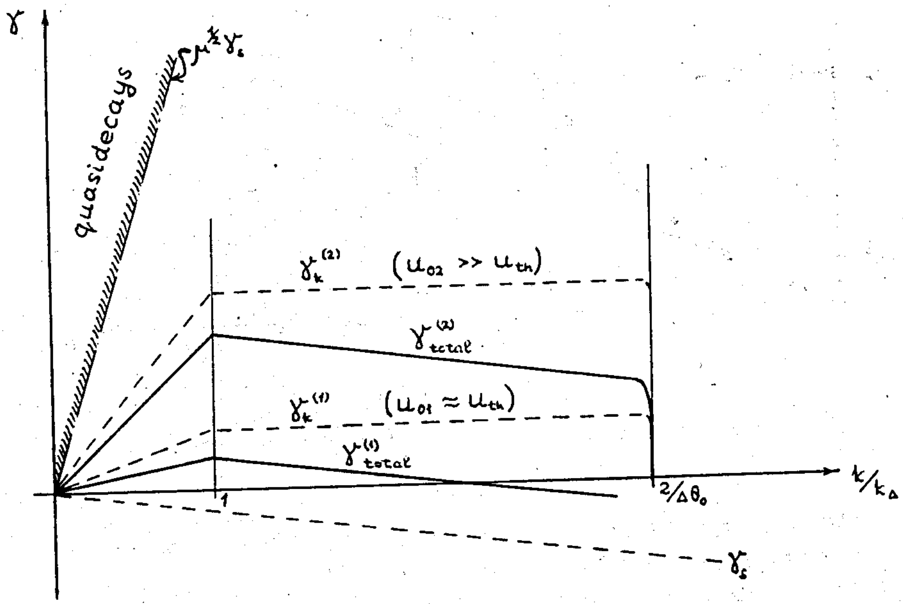

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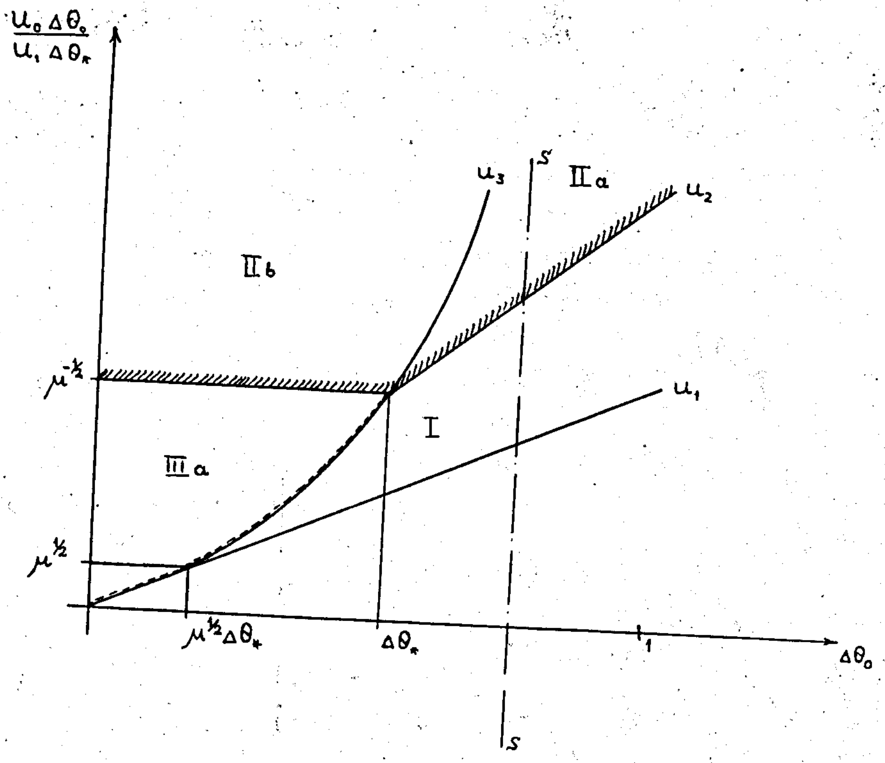

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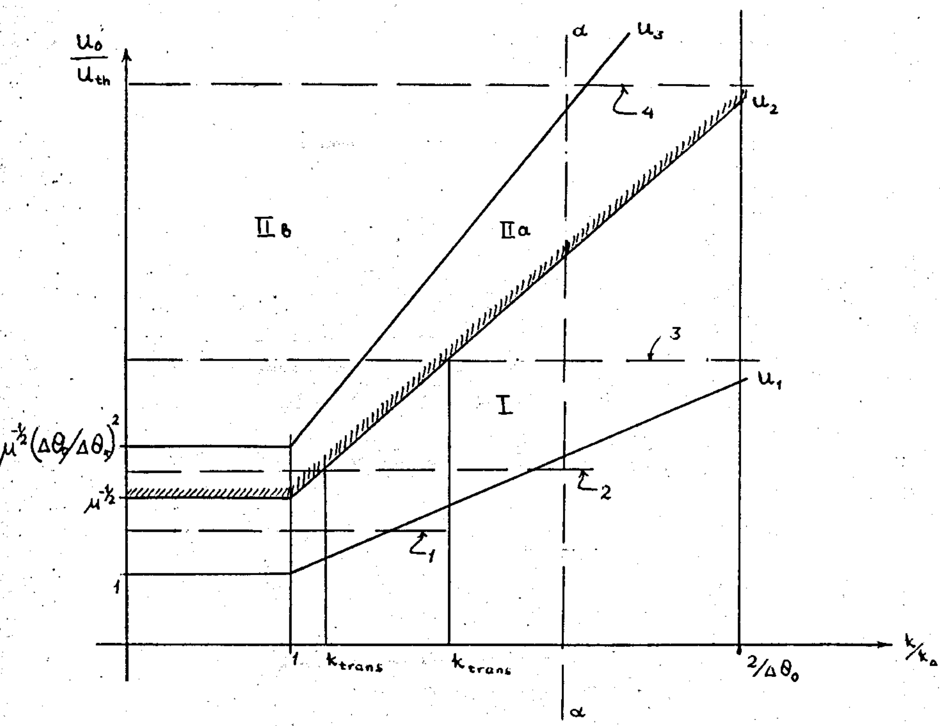

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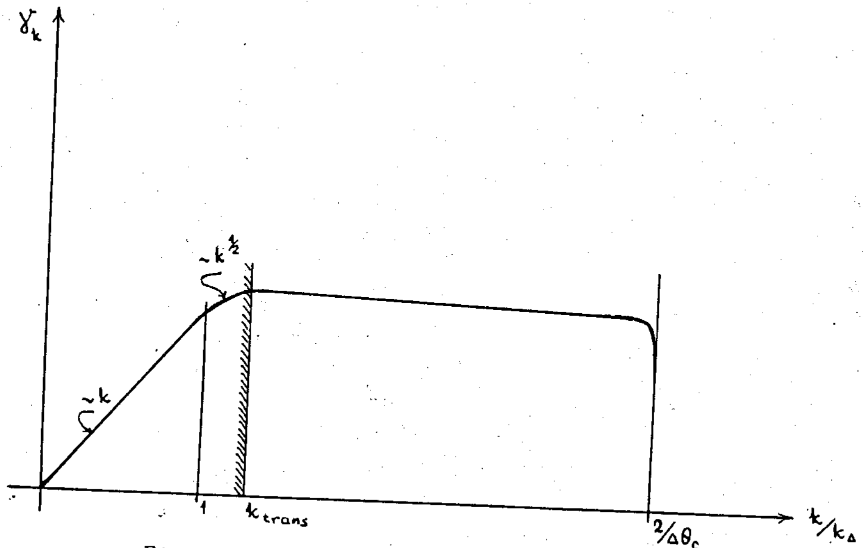

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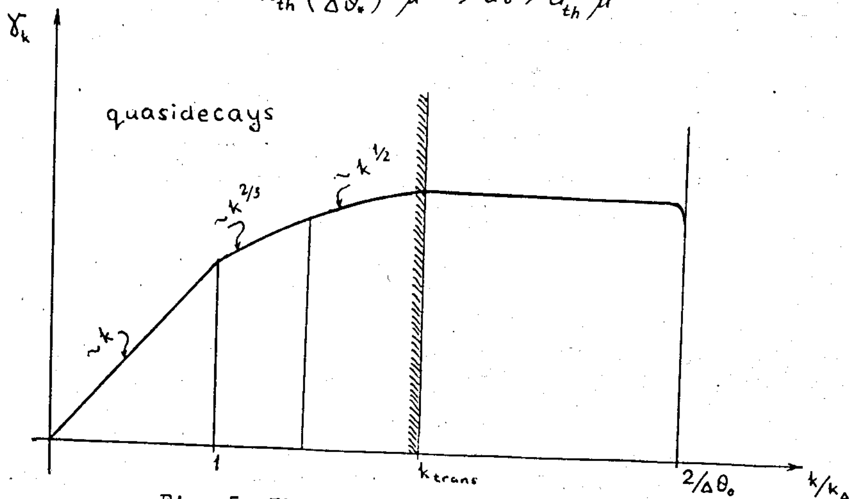

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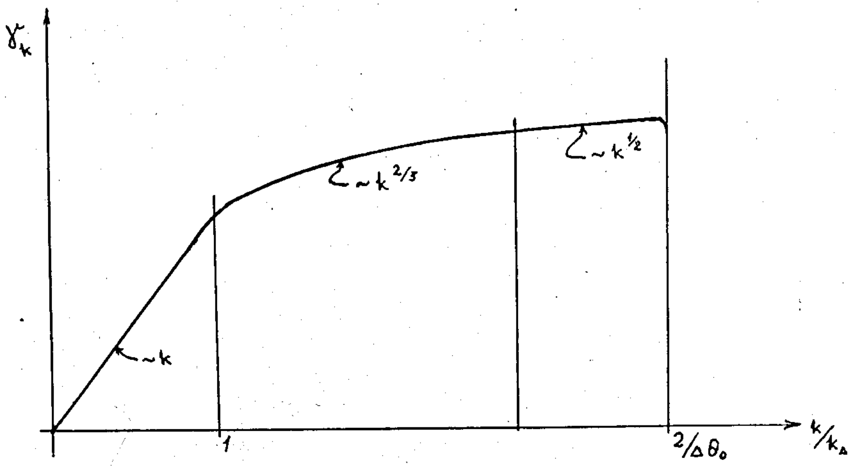

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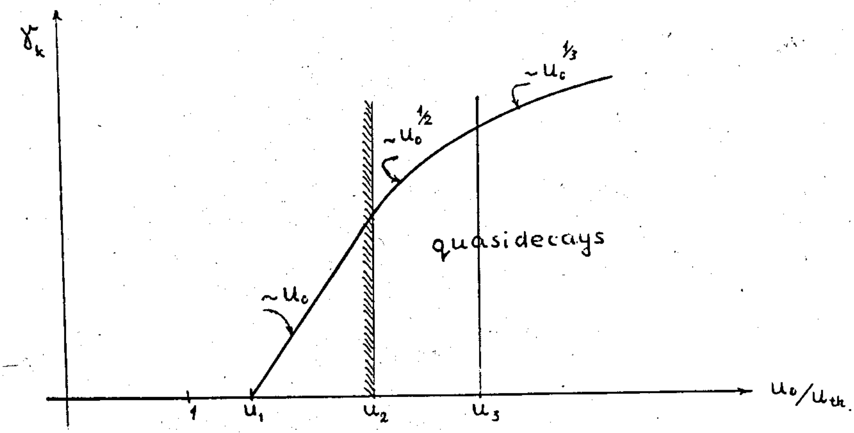

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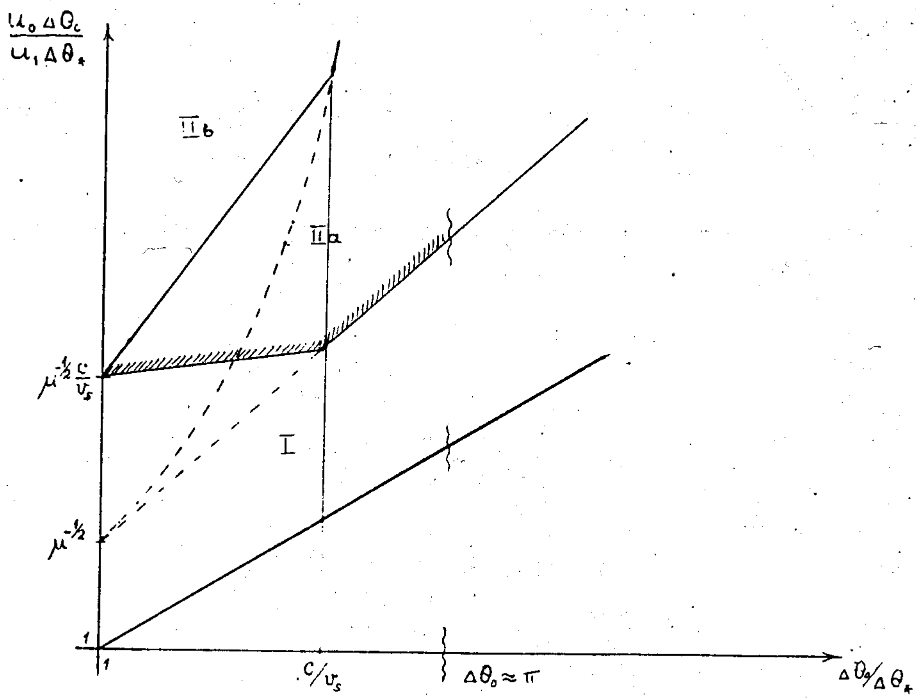

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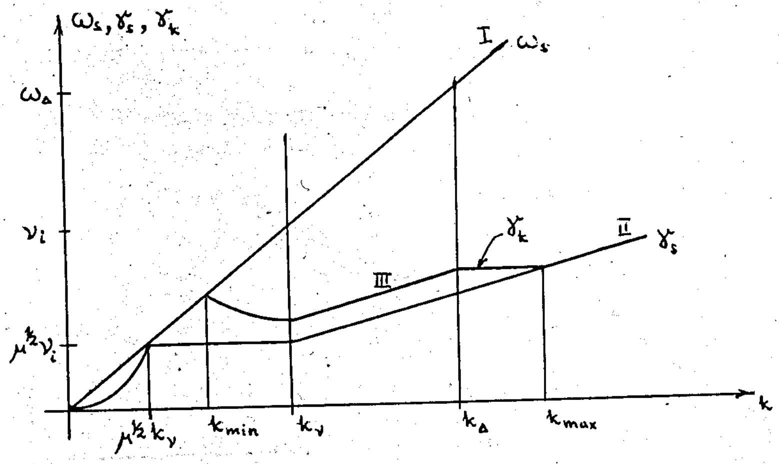

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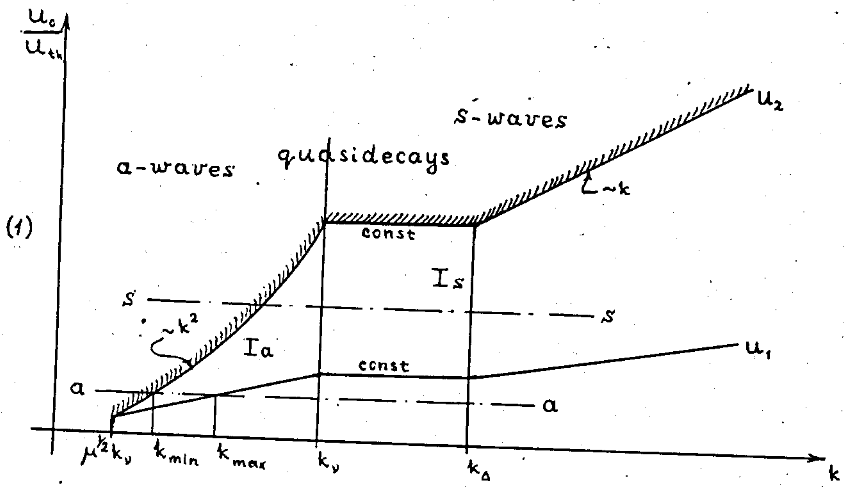

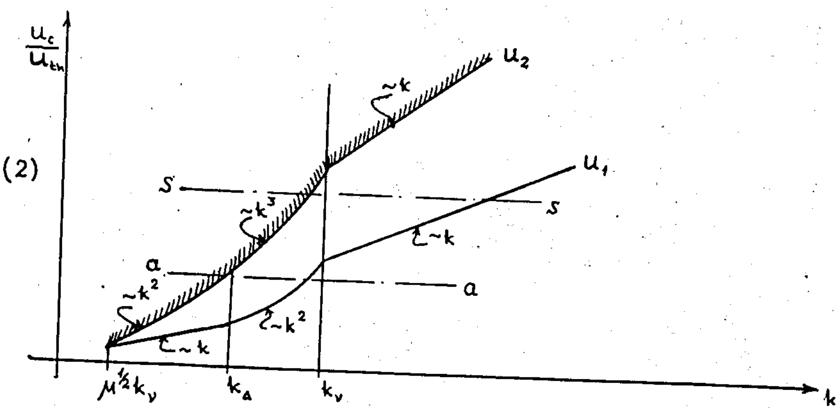

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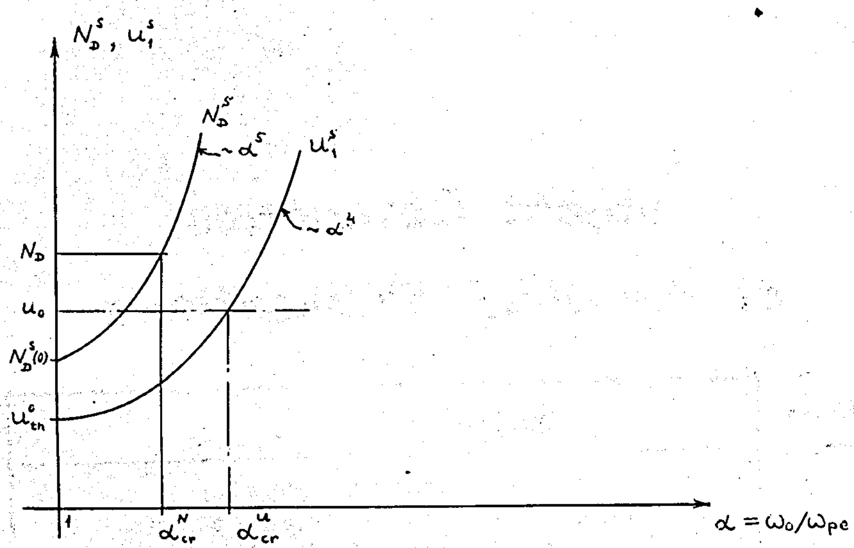

Fig. 11. Dependence of the kinetic S -instability threshold and the critical Debay number N_D^S (17) on $\alpha = \omega_0/\omega_{pe}$:

$$\alpha_{cr}^u = (u_0/u_{th}^0)^{1/4}, \quad u_{th}^0 = u_{th} \quad \omega_0 = \omega_{pe}$$

$$\alpha_{cr}^N = (N_D/N_D^S(0))^{1/5}, \quad N_D^S(0) = N_D^S \quad \omega_0 = \omega_{pe}$$

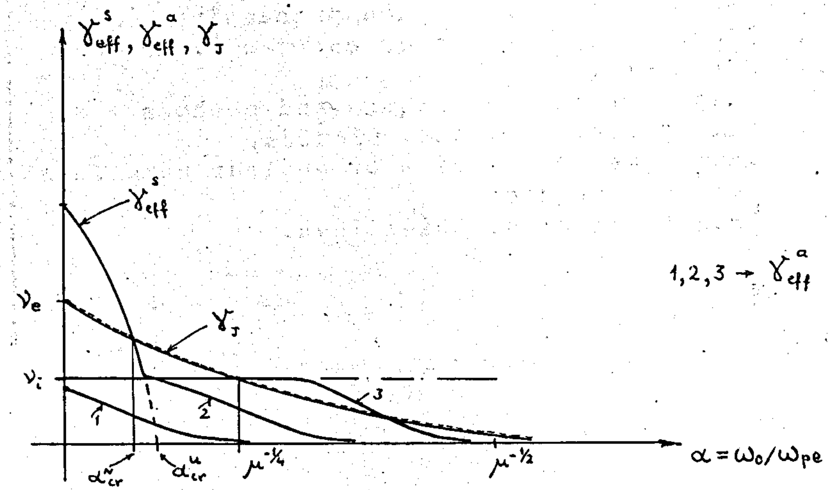

Fig. 12. Dependence of the anomalous heating rate on ω_0/ω_{pe}.