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VELOCITY-ANISOTROPY INSTABILITIES

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In an earlier session Marshall Rosenbluth presented an outline of our present theoretical knowledge of plasma instabilities. It is the pleasant task of some of the other invited speakers to fill in some of the details. Albert Simon has already done this for the two-stream-type instabilities. I will now discuss a class of microinstabilities; namely, those which are due to anisotropies of the velocity distributions of the ions and electrons of the plasma.

Let us assume that the space and time dependence of the electric and magnetic fields is given by

$$\vec{E},\ \vec{B} \sim e^{i(\vec{k}\cdot\vec{x} - \omega t)} .$$

Then, one can solve time linearized Vlasov equations for the perturbed distribution functions, substitute into the charge and current density terms in Maxwell's equations and finally obtain (Reference 1, Chapter 1)

$$\left(k^2 \overleftrightarrow{1} - \vec{k}\vec{k} - \frac{\omega^2}{c^2} \overleftrightarrow{\epsilon}(\vec{k},\ \omega)\right) \cdot \vec{E} = 0 . \qquad (1)$$

Here all the properties of the plasma are contained in the dielectric tensor $\overleftrightarrow{\epsilon}$. It is a function of the wave vector $\vec{k}$ and frequency ω and is a functional of the distribution functions $f_{so}(v)$ (s = electrons, ions) of the unperturbed plasma.

Eqs. (1) are three homogeneous equations for the components of $\vec{E}$. Setting the determinant of the coefficients equal to zero gives a relation between ω and $\vec{k}$ -- the plasma dispersion relation. What is of interest here is the conditions under which the dispersion relation will have a solution with a positive imaginary part of ω. Such a solution represents a wave which grows exponentially in time.

Eq. (1) is rather difficult to treat in general. We shall simplify it by making the so-called quasi-electrostatic approximation; that is, we assume

$$\vec{E} = -\nabla\phi = -i\vec{k}\phi \ . \tag{2}$$

Then Eq. (1) gives

$$\epsilon(\vec{k}, \omega)\ \phi(\vec{k}, \omega) = 0 \tag{3}$$

where

$$\epsilon(\vec{k}, \omega) = \frac{1}{k^2}\ \vec{k} \cdot \overleftrightarrow{\epsilon}(\vec{k}, \omega) \cdot \vec{k} \tag{4}$$

is called the dielectric function of the plasma. Actually, it is not possible to decouple longitudinal and transverse waves except in special cases, so Eq. (2) is not a rigorously valid assumption. It is approximately valid if $\beta = 8\pi P/B^2 \ll 1$.

For an infinite uniform plasma in a magnetic field $\vec{B} = \vec{e}_z\ B_o$, the dielectric function has the form (Reference 1, Chapter 7)

$$\epsilon(\vec{k}, \omega) = 1 + \sum_s \frac{\omega_{ps}^2}{k^2} \sum_{n=-\infty}^{+\infty} \int d^3v\ \frac{J_n^2\left(\frac{k_\perp\ v_\perp}{\omega_{cs}}\right)}{(\omega - k_z v_z - n\omega_{cs})}$$

$$\left[\frac{n\omega_{cs}}{v_\perp} \frac{\partial\ f_{os}}{\partial\ v_\perp} + k_z \frac{\partial\ f_{os}}{\partial\ v_z}\right] . \tag{5}$$

Here ω_{ps} and ω_{cs} are plasma and cyclotron frequencies for species s, and $f_{os}(v_\perp, v_z)$ are the unperturbed velocity distribution functions. J_n is a Bessel function or order n. In doing the velocity space integral it must first be assumed that ω has a positive imaginary part (this is essentially the Landau prescription for treating the singularity); then $\epsilon(\vec{k}, \omega)$ is analytically continued into the rest of the complex ω-plane. Finally, $\epsilon(\vec{k}, \omega) = 0$ gives the dispersion relation.

A tremendous amount of work has been done in solving $\epsilon = 0$ when ϵ is given by Eq. (5) and various assumptions are made about $f_{os}(v_\perp, v_z)$. We shall not attempt to summarize the results but rather to distill from this work some insights into the nature of the instabilities. We shall be particularly interested in conditions near the threshold for instability; that is, for small values of the imaginary part of ω. Then, we may use

$$\frac{1}{\omega - k_z v_z - n\omega_{cs}} \xrightarrow[\mathrm{Im}\,\omega \longrightarrow 0]{} P \frac{1}{\omega - k_z v_z - n\omega_{cs}} - i\pi\delta(\omega - k_z v_z - n\omega_{cs}) \quad (6)$$

to write

$$\epsilon(\vec{k}, \omega) = \epsilon_1(\vec{k}, \omega) + i\epsilon_2(\vec{k}, \omega) \quad (7)$$

where ϵ_1 and ϵ_2 are the real and imaginary parts of ϵ when ω is real. If we write $\omega = \Omega + i\gamma$ and assume that $|\gamma| \ll |\Omega|$ and $|\epsilon_2| \ll |\epsilon_1|$, then we find as an approximation that the real part of the frequency is a solution of

$$\epsilon_1(\vec{k}, \Omega_k) = 0 \quad (8)$$

and the imaginary part of the frequency is given by

$$\gamma_{\vec{k}} = -\frac{\epsilon_2(\vec{k}, \Omega_{\vec{k}})}{\left(\frac{\partial\epsilon_1}{\partial\epsilon}\right)_{\Omega_{\vec{k}}}} = -\frac{P}{2W} = \frac{-R_e \vec{E}^* \cdot \vec{J}}{\frac{1}{4\pi}|\vec{E}|^2 \left(\frac{\partial}{\partial\omega}\omega\epsilon_1\right)_{\Omega_k}} . \quad (9)$$

The last part of Eq. (9) follows from $\vec{J} = \sigma\vec{E}$ where

$$\sigma = \frac{i\omega}{4\pi}(1 - \epsilon) \tag{10}$$

is the plasma conductivity. We may interpret $\gamma_{\vec{k}}$ as the ratio of the power dissipated by the wave to twice the energy of the wave. The energy of the wave is given by the electric field energy $|\vec{E}|^2/8\pi$ times the factor $\partial\omega\epsilon_1/\partial\omega$ which corrects for the energy of the particles which move in response to the wave (Reference 1, Chapter 1).

If the energy of the wave is positive, then instability occurs (i.e., $\gamma_k > 0$) when the dissipation P is negative. From Eq. (6) it is seen that the particles which contribute to ϵ_2 are those for which the Doppler shifted frequency $\omega - k_z v_z$ is just equal to a harmonic of the cyclotron frequency $n\omega_{cs}$. These particles are in resonance with the wave and hence strongly absorb energy from and emit energy into the wave. If more particles emit than absorb then the wave grows; if more absorb than emit then the wave is damped.

The energy of a wave can be negative if $\partial\epsilon_1/\partial\omega < 0$. In this case a positive dissipation removes energy from the wave making it more negative and hence increasing its amplitude.

We shall now specialize to the case of the two-temperature Maxwell-Boltzmann distribution given by

$$f_{os}(v_\perp, v_z) \sim e^{-\frac{m}{2}(v_\perp^2/T_\perp + v_z^2/T_\parallel)} \; . \tag{11}$$

We shall also assume that $T_\perp$ and $T_\parallel$ are very small for the electrons. Then, ϵ_1 is determined almost entirely by the electrons and

$$e_1(\vec{k}, \omega) \simeq 1 - \frac{\omega_{pe}^2}{\omega^2}\frac{k_z^2}{k^2} \; . \tag{12}$$

Solving $\epsilon_1 = 0$ for the frequency gives

$$\Omega_{\vec{k}} = \omega_{pe} \frac{k_z}{k} \; . \tag{13}$$

These are plasma oscillations in a magnetic field. The frequency depends on the cosine of the angle between $\vec{k}$ and the field and ranges between zero and ω_{pe}.

Figure 1 shows the contributions to $\gamma_{\vec{k}}$ as a function of $\Omega_{\vec{k}}$. The terms due to $\partial f_{so}/\partial v_z$ in Eq. (5) are shown as solid lines. The electrons contribute a large but narrow peak near the origin. The ion contributions are centered on the harmonics of ω_{ci}. The ion contribution from the $\partial f_{so}/\partial v_\perp$ terms are drawn as dotted lines; these contributions are always negative. In order for $\gamma_{\vec{k}}$ to be positive the positive contributions must exceed the negative. This can only happen near to (but slightly below) the harmonics of ω_{ci}. Also $\Omega_{\vec{k}}$ must be in the range $0 < \Omega_k \leq \omega_{pe}$, so

$$\omega_{pe} > n\omega_{ci} \tag{14}$$

is necessary for instability of the nth harmonic.[2] In order for the positive contributions from the ions to exceed the negative contributions, it may be shown that

$$\left(\frac{T_\parallel}{T_\perp}\right)_{ions} < \frac{1}{2n} \tag{15}$$

is necessary.[3]

If the electron temperature is increased, the large negative contribution of the electrons shown on Fig. 1 widens and its amplitude will decrease. The first effect of increasing the electron temperature is to stabilize the first few harmonics. If the electron temperature increases until the λ_k - Ω_k plot looks like Fig. 2, then the cold electron assumption which led to Eq. (12) is no longer valid. Instead ϵ_1 becomes

$$\epsilon_1(\vec{k}, \omega) \simeq 1 + \frac{1}{k^2 L_D^2} + \frac{\omega_{pi}^2}{\omega^2} \frac{k_z^2}{k^2} \tag{16}$$

where L_D is the electron Debye length. The second term is the electron contribution and the third term is the n = 0 ion contribution. Solving for the frequency gives

$$\Omega_{\vec{k}} = \frac{\omega_{pi} k_z}{\sqrt{k^2 + \frac{1}{L_D^2}}} \simeq \sqrt{\frac{T_e}{M}}\, k_z \ . \tag{17}$$

These are ion sound waves. Their maximum frequency is ω_{pi}; so Eq. (14) must be replaced by

$$\omega_{pi} > \omega_{ci} \tag{18}$$

as the necessary condition for instability in the region of high electron temperatures. Eq. (15) is unchanged.

We can now see changes in the distribution function that would increase the instability. For the distribution function of Eq. (11) the $\partial f_{os}/\partial v_\perp$ terms in Eq. (5) give the negative contributions (dotted curves) in Figs. 1 and 2. These contributions can be made positive if $\partial f_{os}/\partial v_\perp$ is positive where $J_n^2(k_\perp v_\perp/\omega_{cs})$ is large. What we have in mind is a distribution like that shown in Fig. 3. This is a loss cone distribution and the instability it gives rise to is the loss cone instability.[4] The case $k_z = 0$ is rather interesting. For it Eq. (5) takes the form[5]

$$\epsilon(\vec{k}, \omega) = 1 + \sum_s \frac{\omega_{ps}^2}{k^2} \sum_{n=-\infty}^{+\infty} \frac{n\omega_{cs}}{\omega - n\pi_{cs}} \int d^3v \, \frac{\partial f_{os}}{\partial v_\perp} \frac{1}{v_\perp} J_n^2\left(\frac{k_\perp v_\perp}{\omega_{cs}}\right) . \tag{19}$$

For sufficiently narrow loss cone distributions $\epsilon(\vec{k}, \omega) = 0$ gives instabilities with the real part of ω equal to zero. This would seem to contradict the comment made in Rosenbluth's paper that low frequency velocity space instabilities were excluded by the conservation of $\mu = {v_\perp}^2/B$. Actually, μ is not conserved here because the wavelengths of the unstable modes are shorter than the particles radius of gyration.

I shall now leave the subject of quasi-electrostatic instabilities and say a few brief words about electromagnetic instabilities. If $\vec{k}$ is parallel to the unperturbed magnetic field $\vec{B}_o$, then longitudinal and transverse waves are decoupled. One finds that Eq. (1) has circularly polarized transverse wave solutions with

$$E_y = \pm \, iE_x \tag{20}$$

and

$$\frac{k^2c^2}{\omega^2} = \epsilon_{11} \pm i\epsilon_{12} = 1 + \sum_s \frac{4\pi e_s^2}{m_s\omega^2} \int d^3v \, \frac{v_\perp/2}{\omega - kv_z \pm \omega_{cs}}$$

$$\left[(\omega - kv_z) \frac{\partial f_{os}}{\partial v_\perp} + kv_\perp \frac{\partial f_{os}}{\partial v_z}\right] . \tag{21}$$

An analysis of this dispersion relation yields Alfven, whistlers and light waves; these may be unstable for sufficiently anisotropic velocity distributions. I see in the program that some of these instabilities are being discussed in the contributed papers, so I shall not discuss them further.

Now I shall return to the subject of the quasi-electrostatic instabilities. Theory predicts instabilities with frequencies near harmonics of the ion cyclotron frequency when the ions have a loss cone distribution or a sufficiently anisotropic distribution. Such distributions are expected in magnetic mirror confined plasmas. Indeed, large amplitude oscillations at harmonics of ω_{ci} are observed in such plasmas. Figure 4 shows spectra of oscillations observed in the DCX-2 experiment at the Oak Ridge National Laboratory.[6] Harmonics up to the 100th are observed. In fact, it is rather difficult to explain on the basis of linear stability theory why such high harmonics are unstable. I believe that a more likely explanation is that only the first few harmonics are linearly unstable and that nonlinear processes feed energy into the higher harmonics. The nonlinear process involved here is the three-wave interaction in which two waves with wave vectors and frequencies $\vec{k}_1$, ω, and $\vec{k}_2$, ω_2 combine to form a third wave with $\vec{k}_3$, ω_3 as shown in Fig. 5. For this to be possible it is necessary to have

$$\vec{k}_1 + \vec{k}_2 = \vec{k}_3 \tag{22}$$

and

$$\omega_1 + \omega_2 = \omega_3. \tag{23}$$

This is possible if the frequencies are given by either Eq. (13) or (17). If $\omega_1 = n_1\omega_{ci}$ and $\omega_2 = n_2\omega_{ci}$ then $\omega_3 = (n_1 + n_2)\omega_{ci}$. In this way the higher harmonics obtain their energy from lower harmonics. The exchange of energy between the waves by non-linear intereactions is a complicated process and quantitative predictions are difficult. An attempt to explain some experimental results in these terms has been published elsewhere. Although not as successful as could be desired, it seems to be a step in the right direction.[7]

The linear theory of instabilities in infinite homogeneous plasmas is now rather well developed. In order for theory to be really useful in interpreting experiments, we need to go beyond this and develop the theory of (a) non-linear effects such as those mentioned in the preceeding paragraph, and (b) finite-plasma effects.

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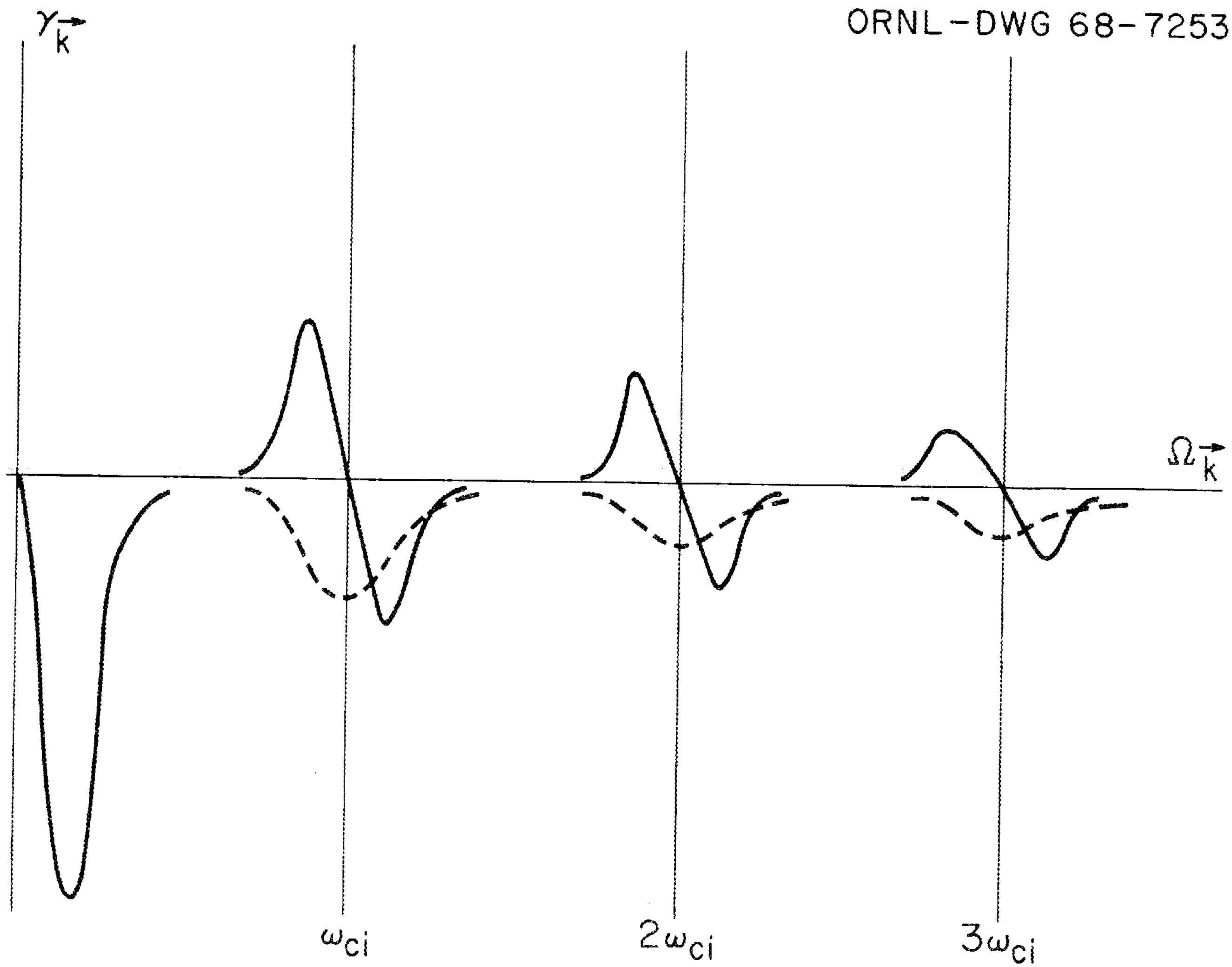

Fig. 1. Contributions to $\gamma_{\vec{k}}$ for a Cold Electron Plasma.

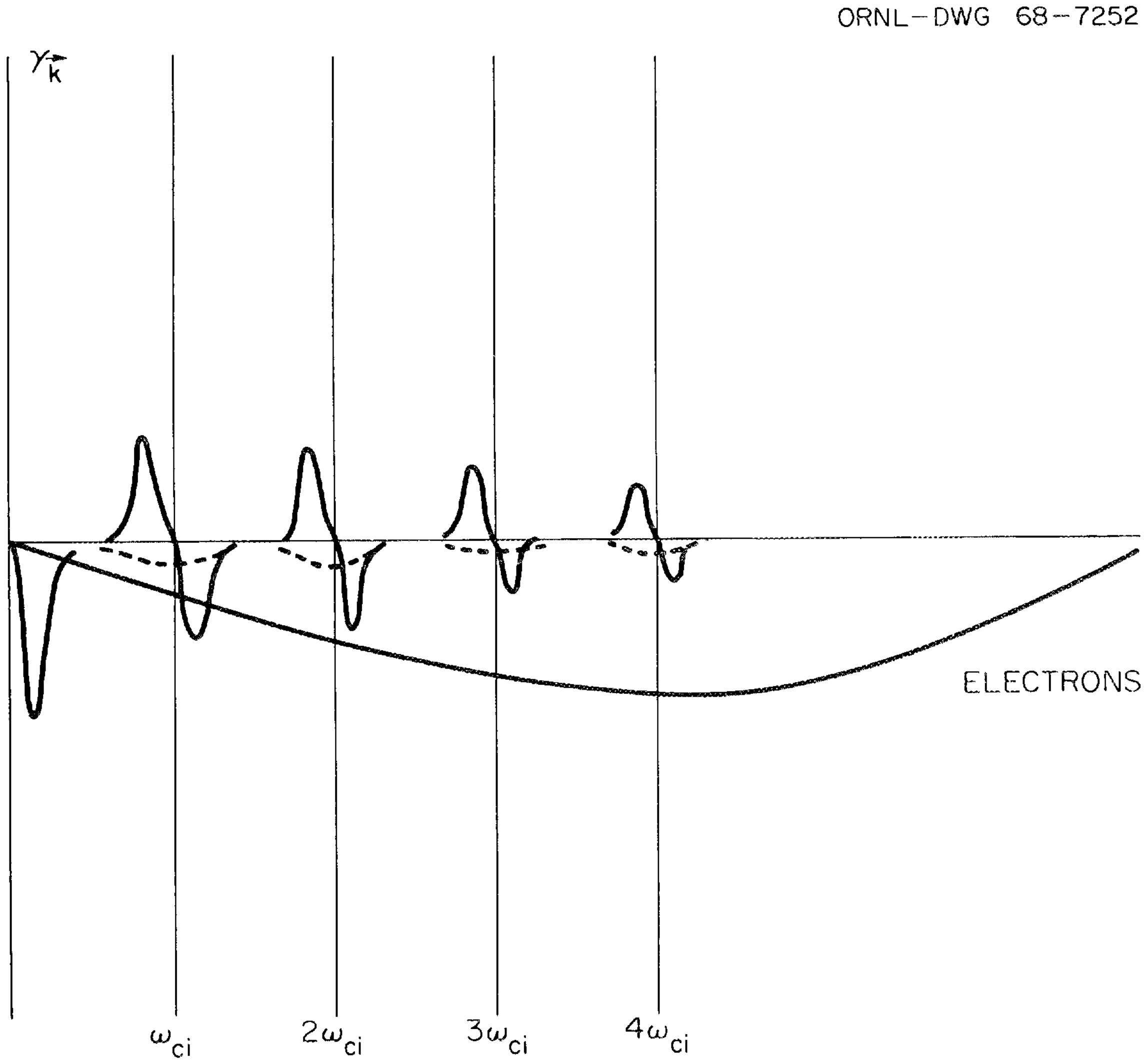

Fig. 2. Contributions to $\gamma_{\vec{k}}$ for a Hot Electron Plasma.

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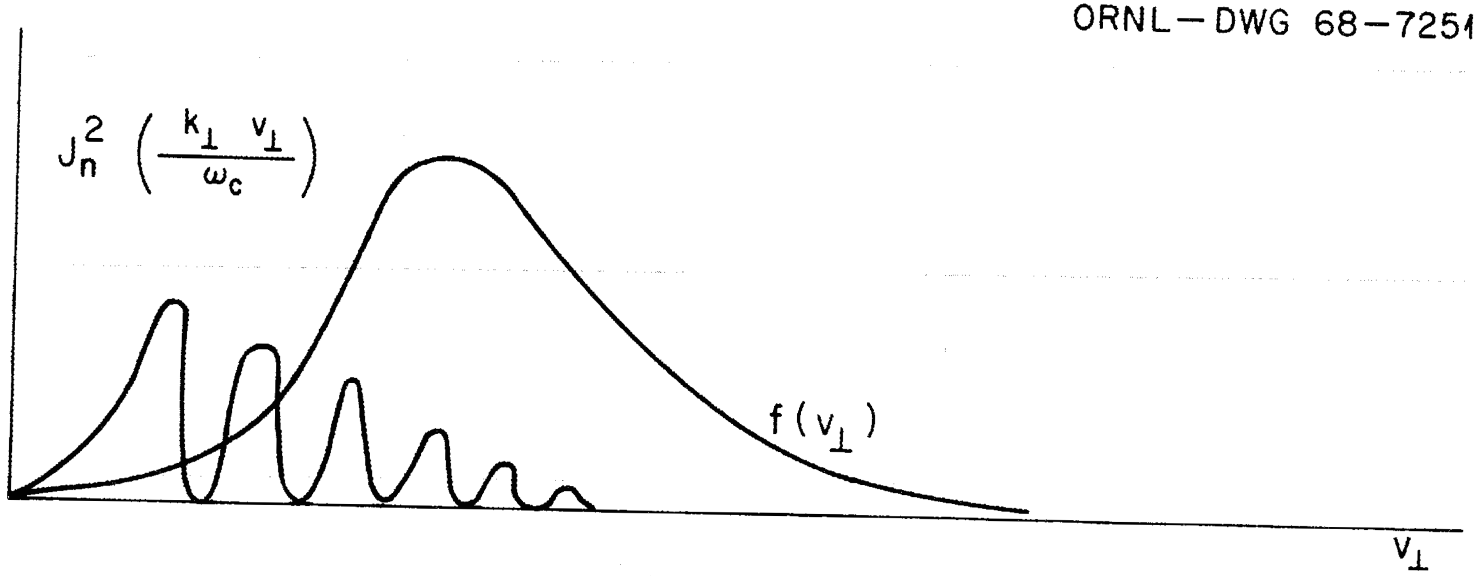

Fig. 3. Loss Cone Distribution.

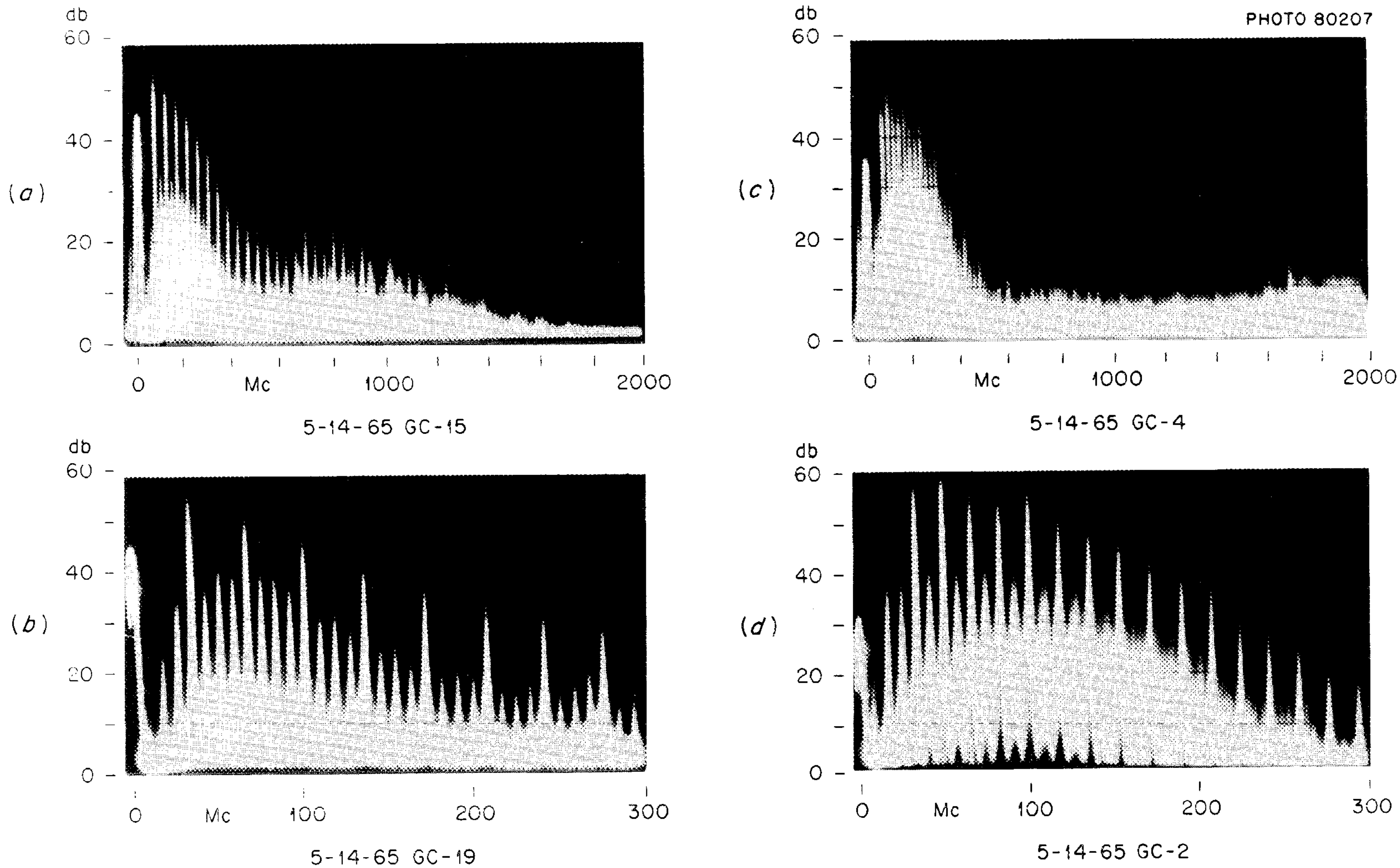

Fig. 4. Comparison of rf Spectra with Gas Dissociation and H_2 Arc Distribution. (a) and (b) are dissociation by hydrogen gas, (c) and (d) are dissociation in the hydrogen arc.

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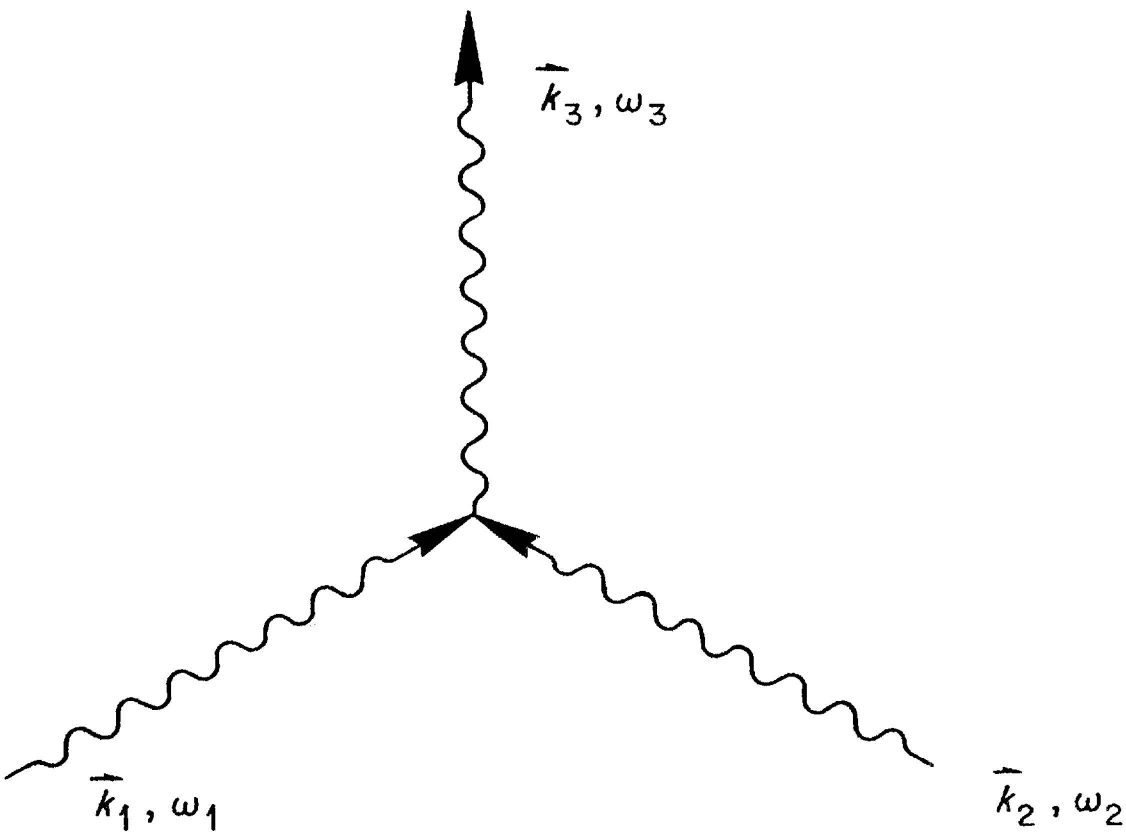

Fig. 5. Non-linear Coupling of Three Waves.

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