Resonant quasiparticles in plasma turbulence

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A general view is proposed on wave propagation in fluids and plasmas where the resonant interaction of monochromatic waves with quasiparticles is considered. A kinetic equation for quasiparticles is used to describe the broadband turbulence interacting with monochromatic waves. Resonant interactions occur when the phase velocity of the long wavelength monochromatic wave is nearly equal to the group velocity of short wavelength wave packets, or quasiparticles, associated with the turbulent spectrum. It is shown that quasiparticle Landau damping can take place, as well as quasiparticle beam instabilities, thus establishing a direct link between short and large wavelength perturbations of the medium. This link is distinct from the usual picture of direct and inverse energy cascades, and it can be used as a different paradigm for the fluid and plasma turbulence theories.

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I. INTRODUCTION

Plasmas and fluids are usually in a turbulent state. One of the main physical problems is then to understand how monochromatic waves or large scale structures can develop, propagate, or be damped by the turbulence. Here, we propose a global approach where the broadband turbulence is described as a gas of quasiparticles. Such an approach is valid as long as we can identify two space and time scale ranges; a short one associated with the internal oscillations of these quasiparticles and a long one associated with the individual monochromatic waves or large scale structures.

In such a picture of plasma and fluid turbulence, the main physical processes are those leading to a direct coupling between large scale structures (the wave) and short scale structures of the medium (the turbulence). This coupling can lead to an energy exchange in opposite directions. The energy transfer from turbulence to large scale structures can be due to kinetic and hydrodynamic instabilities of the quasiparticle gas. This has been improperly described as an inverse cascading process, but it consists on a single step process and not really a cascade.

The opposite case of an energy transfer to the short scale turbulence is due to wave damping, associated with the resonant interaction between a long wavelength wave and the quasiparticles. It extends the concept of a noncollisional Landau damping from the real particles (electrons and ions, in the original concept) to the quasiparticle distributions.

The universality of the resonant processes should be stressed here. Several particular examples can already be found in the literature, and illustrate such universality:

(i) Electron plasma waves propagating in a photon gas [1]; this is relevant to laser-plasma interactions and to astrophysics.

(ii) Ion acoustic waves in a plasmon gas, where the plasmons describe the broadband electron plasma wave turbulence [2,3]; this process can be seen as of an academic interest, which, however, provides a useful basis for different experimental plasma studies.

(iii) Dust lattice waves in a turbulent plasma sheath [4]; these waves can become parametrically unstable due to the turbulent fluctuations of the plasma sheath where dust plasma crystals are formed, and can eventually lead to melting and sublimation of the crystalline structure. The background quasiparticles here are plasma phonons.

(iv) Zonal flows in which one could call a drifton gas [5–7] (which describes a large spectrum of drift wave turbulence); this is a relevant process for anomalous transport in magnetic fusion plasmas. Similar processes can also be found in fluid dynamics, involving zonal flows and Rossby waves in the oceans or in the atmosphere of the planets [8].

But the universality of the physical picture described here is not exhausted by these examples, and a much larger number of relevant examples can be added. As an extreme example, we could consider the excitation of gravitational waves by a photon gas [9], where the fluid is replaced by the vacuum metric field and the photons can be considered as quasiparticles in the sense of short scale electromagnetic wave packets.

II. DISPERSION RELATION

The general form of the dispersion relation of electrostatic waves in a turbulent plasma can be written as

$$\epsilon(\omega,\vec{k}) = -\chi_{\mathrm{qp}}, \tag{1}$$

where χ_{qp} is the quasiparticle susceptibility, if the turbulent state is described as a quasiparticle gas, and $\epsilon(\omega,\vec{k})$ is the dielectric function for a mode with frequency ω and wave vector $\vec{k}$. Usually, we can also write $\epsilon(\omega,\vec{k}) = 1 + \Sigma_\alpha \chi_\alpha$, where the sum is over the different charged particle species present in the plasma. It is important to note that χ_{qp} contains both resonant and nonresonant quasiparticle contributions, and that the resonant part plays an important role in the energy balance of the plasma, as discussed below.

This kind of dispersion relation can be established by starting with an equation of the propagation for the potential ϕ (or, in alternative, for the density perturbation $\tilde{n}$) of the form

$$L(\vec{r},t,N(\vec{k}'))\phi=0, \tag{2}$$

where L is a space-time differential operator, which depends on the quasiparticle number density $N(\vec{k}')$, and $\vec{k}'$ is the quasiparticle momentum. Note that this is a nonlinear dispersion relation, where $N(\vec{k}')$ is also related to the potential ϕ.

By imposing a perturbation of the form $\exp i(\vec{k}\cdot\vec{r}-\omega t)$ on both ϕ and the perturbed distribution $\tilde{N}(\vec{k}')$, we obtain

$$\epsilon(\omega,\vec{k})\phi=\int g(\vec{k},\vec{k}')\tilde{N}(\vec{k}')d\vec{k}'. \tag{3}$$

Obviously, $\epsilon(\omega,\vec{k})=0$ will be the dispersion relation of the slow wave, in the absence of turbulence. But, in general, we need to establish a relation between $\tilde{N}(\vec{k}')$ with ϕ, in order to derive the dispersion relation (1). This is done by using an evolution equation for the turbulence field in the form of a kinetic wave equation for the quasiparticle distribution.

III. KINETIC EQUATION

Before stating the kinetic wave equation, we need to identify an invariant quantity N, valid for a slowly perturbed turbulent state, in the form of an integral over the sixth-dimensional (position and momentum) phase space $(\vec{r},\vec{k}')$:

$$N=\int N(\vec{k}',\vec{r},t)d\vec{r}d\vec{k}'. \tag{4}$$

In most situations, the quantity appearing inside the integral is nothing but the energy density divided by the frequency, and takes the obvious physical meaning of a quasiparticle number density or wave action $N(\vec{k}',\vec{r},t)=W(\vec{k}',\vec{r},t)/\hbar\omega'$, where we can take $\hbar=1$, and where ω' is the energy of the quasiparticles with momentum $\vec{k}'$. We can thus identify this frequency with the function $\omega'=\omega'(\vec{k}')$, as determined by the linear dispersion relation of the short wavelength turbulence modes. This is valid for the electromagnetic turbulence (where the photons in a plasma can be seen as an extreme example of quasiparticles and more appropriately described as dressed particles in the usual sense of the field theory), as well as for the electrostatic turbulence of the electron plasma or the ion acoustic types (plasmons and phonons). This is also valid in the case of the pseudo-three-dimensional drift wave turbulence, which can be called a drifton gas, as described by a modified Hasegawa-Mima equation [6,7].

It should be noticed that, in the present approach, the quasiparticle number density $N(\vec{k}',\vec{r},t)$ is not a function of the frequency (or energy) ω', and only of the wave vector (or momentum) $\vec{k}'$. Of course, we could have used a frequency dependent distribution $N(\vec{k}',\omega';\vec{r},t)$, as it can also be found in the literature [11]. If the frequency is uniquely determined by a given dispersion relation, we can simply write $N(\vec{k}',\vec{r},t)=2\pi N(\vec{k}',\omega';\vec{r},t)\delta(\omega'-\omega'(\vec{k}'))$.

A consistent definition of the quasiparticle number density, where the relation between these two possible definitions is clarified, can be given in terms of the Wigner function for the electrostatic or the electromagnetic field [10,12]. Once identified the quantity describing the density distribution $N(\vec{k}',\vec{r},t)$ in the quasiparticle phase space, we can establish the corresponding Liouville's theorem that states its total time invariance:

$$\frac{d}{dt}N(\vec{k}')=\left(\frac{\partial}{\partial t}+\vec{v}'\cdot\frac{\partial}{\partial\vec{r}}+\vec{F}'\cdot\frac{\partial}{\partial\vec{k}'}\right)N(\vec{k}')=0, \tag{5}$$

where, for simplicity, we use $N(\vec{k}')=N(\vec{k}',\vec{r},t)$. Here, $\vec{v}'=\partial\omega'/\partial\vec{k}'$ is the quasiparticle velocity (or equivalently, the group velocity of the turbulence wave packets) and $\vec{F}'=d\vec{k}'/dt$ is the force acting on these quasiparticles due to large scale perturbations of the medium. This force term includes refraction effects (which maintain the value of ω'), as well as quasiparticle acceleration (implying the variation of the energy, or frequency ω'). Quasiparticle kinetic equations similar to Eq. (5) have been widely used in the past, but the importance of the force term has only recently been recognized. It will be shown in this paper that this term plays an essential role in plasma turbulence.

We now consider $N(\vec{k}')=N_0(\vec{k}')+\tilde{N}(\vec{k}')$, and linearize the kinetic wave equation around the unperturbed state $N_0(\vec{k}')$, by assuming a slow perturbation with the frequency ω and the wave vector $\vec{k}$. The result is

$$\tilde{N}(\vec{k}')=-i\frac{\vec{F}'\cdot\partial N_0/\partial\vec{k}'}{(\omega-\vec{k}\cdot\vec{v}')}. \tag{6}$$

Noting that we can write the equivalent force as

$$\vec{F}'=-\frac{\partial\omega'}{\partial\vec{r}}=-i\vec{k}f'(\vec{k}')\phi, \tag{7}$$

one can relate the density perturbation $\tilde{N}(\vec{k}')$ with ϕ. Thus, returning to Eq. (3), we obtain

$$\epsilon(\omega,\vec{k})=-\int f(\vec{k},\vec{k}')\frac{\vec{k}\cdot\partial N_0/\partial\vec{k}'}{(\omega-\vec{k}\cdot\vec{v}')}d\vec{k}', \tag{8}$$

where we have used $f(\vec{k},\vec{k}')=f'(\vec{k}')g(\vec{k},\vec{k}')$. Comparing this equation with Eq. (1), we see that the integral in the right-hand side is nothing but the quasiparticle susceptibility χ_{qp}.

IV. RESONANT CONTRIBUTIONS

In order to identify the resonant and nonresonant contributions to χ_{qp}, we consider the parallel and perpendicular motions of the turbulence quasiparticles with respect to the direction of the propagation of the slow wave:

$$\vec{v}'=u\frac{\vec{k}}{k}+\vec{v}'_\perp, \quad \vec{k}'=p\frac{\vec{k}}{k}+\vec{k}'_\perp. \tag{9}$$

Note that the parallel velocity is a function of the parallel and the perpendicular momenta: $u=u(p,\vec{k}'_\perp)$. Equation (8) shows that resonant interactions of the wave with the quasiparticles can occur when the parallel velocity equals the phase velocity of the slow wave. Then, we can write

$$\chi_{qp}=\int f(\vec{k},\vec{k}'_\perp,p)\frac{\partial N_0/\partial p}{(\omega/k-u)}d\vec{k}'_\perp dp. \tag{10}$$

The function $f(\vec{k},\vec{k}'_\perp,p)$ and the equilibrium distribution N_0 are, in general, continuous and single valued functions, and this integral only has one pole at $p=p_0$, determined by the resonance condition $\omega/k=u$. Developing u around its resonance value $u_0=u(p_0,\vec{k}')=\omega/k$, and introducing a parallel function $G(p)$, defined by

$$G(p)=\int \frac{f(\vec{k},\vec{k}'_\perp,p)}{(\partial u/\partial p)_0}N_0(\vec{k}'_\perp,p)d\vec{k}'_\perp, \tag{11}$$

we can finally write the quasiparticle susceptibility in the form $\chi_{qp}=\chi_r+i\chi_{im}$, where the real and imaginary parts are determined by

$$\chi_r=-\mathrm{P}\int\frac{\partial G(p)/dp}{(p-p_0)}dp, \quad \chi_{im}=-\pi\left(\frac{\partial G}{\partial p}\right)_0, \tag{12}$$

and $\mathrm{P}\int$ represents the principal part of the integral. This shows that the resonant and nonresonant contributions to the quasiparticle susceptibility have distinctive properties. The nonresonant real part leads to a small correction of the linear dispersion relation. However, the resonant imaginary part can lead to the wave damping or growth, according to the sign of the derivative of the parallel function $G(p)$ at $p=p_0$. This qualitatively important effect can thus be identified with quasiparticle Landau damping.

V. QUASIPARTICLE BEAMS

Let us first consider the simple and physically relevant case of a Gaussian beam of quasiparticles, described by

$$G(p)=G_0\exp\left(-\frac{(p-\overline{p})^2}{2\sigma^2}\right), \tag{13}$$

where G_0 represents the beam intensity and σ is the spectral width. Note that this will also correspond to a nearly Gaussian parallel distribution for the number density $N_0(\vec{k}'_\perp,p)$, if $f(\vec{k},\vec{k}'_\perp,p)$ and $(\partial u/\partial p)_0$ are slowly functions of p around $p=\overline{p}$. The maximum value for the resonant part of quasiparticle susceptibility corresponds to $p=\overline{p}\pm\sigma$, and it is equal to

$$\chi_{max}=\chi_{Im}(\overline{p}\pm\sigma)=\mp\frac{\pi}{e\sigma}G_0. \tag{14}$$

We see that it decreases with an increasing spectral width, and it is proportional to the intensity of the beam: $|\chi_{max}|\propto G_0$. This means that the resulting kinetic instabilities will typically have a growth rate proportional to G_0.

Another important case corresponds to the monoenergetic particle beam with a negligible spectral width, $\sigma\sim 0$, such that (for one-dimensional problems) it can be represented by

$$N(\vec{k}')=N_0\delta(\vec{k}'_\perp)\delta(p-\overline{p}). \tag{15}$$

In order to study this case, we can write χ_{qp} in the form

$$\chi_{qp}=f(k,0,\overline{p})\int\frac{\partial N/\partial p}{(\omega/k-u)}dp$$
$$=-f(k,0,\overline{p})\int\frac{N_0\delta(p-\overline{p})}{(\omega/k-u)^2}dp, \tag{16}$$

and finally obtain

$$\chi_{qp}=-\frac{f(k,0,\overline{p})k^2N_0}{(\omega-ku_0)^2}=-\frac{\Omega_{qp}^2}{(\omega-ku_0)^2}. \tag{17}$$

We see that this takes the familiar form of the susceptibility of electron or ion beams with velocity u_0 and density N_0. The frequency Ω_{qp} plays the role of a quasiparticle plasma frequency, proportional to the square root of the beam density. Here, again, the contribution of this term to the total wave dispersion relation will become relevant for nearly resonant conditions, such that $\omega\simeq ku_0$. Then, we can use $\omega=ku_0+\eta$, with $|\eta|\ll ku_0$. The dispersion relation (8) will take the form

$$\epsilon(\eta,k)=\frac{f(k,0,\overline{p})k^2N_0}{\eta^2}. \tag{18}$$

For solutions such that $\mathrm{Im}(\eta)>0$, we will have a hydrodynamic type of beam instability, with growth rates that vary with the beam density, typically between $N_0^{1/2}$ and $N_0^{3/2}$. Thus, these appear much stronger than the kinetic beam instabilities associated with the inverse process of Landau damping [1,3,6,9,13].

VI. QUASIPARTICLE TRAPPING

We can also push forward the quasiparticle concept and consider the motion of individual quasiparticles, as described by the characteristics of the kinetic wave equation (5):

$$\frac{d\vec{r}}{dt}=\vec{v}, \quad \frac{d\vec{k}'}{dt}=\vec{F}'=-f'(\vec{k}')\vec{\nabla}\phi(\vec{r},t), \tag{19}$$

where $\vec{r}$ is the position of the quasiparticles (average position of individual short wavelength wave packets). Obviously, this force $\vec{F}'$, is modulated due to the existence of the monochromatic slow wave. If we assume the propagation of this slow wave along the direction Ox, as defined by $\phi(\vec{r},t)=\phi_0\cos(kx-\omega t)$, we can describe the parallel motion of the quasiparticles by

$$\frac{dx}{dt}=\left(u-\frac{\omega}{k}\right), \quad \frac{dp}{dt}=-kf'(p)\phi_0\sin(kx-\omega t). \tag{20}$$

The perpendicular motion is trivially determined by $\vec{k}'_\perp = \text{constant}$. These equations show the existence of an elliptic fixed point at

$$u(p) = \frac{\omega}{k}, \quad x = \frac{\pi}{k}. \tag{21}$$

This means that, for quasiparticles such that $u(p) \simeq \omega/k$, we will have trapped oscillations at the bottom of the slow wave potential, with small amplitudes $\tilde{x} = x - \pi/k$ around the fixed point. From Eqs. (21), we can then derive

$$\frac{d^2\tilde{x}}{dt^2} = -k^2 f'(p) \frac{\partial u}{\partial p} \phi_0 \tilde{x}, \tag{22}$$

which is the equation for the harmonic oscillator with a frequency

$$\omega_b = k\sqrt{f'(p)(\partial u/\partial p)\phi_0}, \tag{23}$$

which is basically the bounce frequency for deeply trapped oscillations of quasiparticles in the slow wave potential ϕ. The similarities with the electron bounce frequency are striking. Kinetic effects leading to the appearance of the Kruer mode in a quasiparticle gas is then conceivable.

VII. QUASIPARTICLE DIFFUSION

The above kinetic approach can be extended to a quasilinear theory [1,3] if, instead of considering a single monochromatic slow wave, with the frequency ω and the wave vector $\vec{k}$, we consider a larger spectrum of slow waves interacting with the high-frequency turbulence. In this case, we can make a time average over a time scale much longer than the typical period $1/\omega$, and obtain an equation for nearly unperturbed quasiparticle number density

$$\frac{\partial N_0}{\partial t} = -\int \vec{F}'^*_k \cdot \frac{\partial}{\partial \vec{k}'} \tilde{N}_k(\vec{k}') d\vec{k}, \tag{24}$$

where $\vec{F}'^*_k$ is the complex conjugate of the force acting on the quasiparticles with momentum $\vec{k}'$, and due to the slow component $(\omega, \vec{k})$, and $\tilde{N}_k(\vec{k}')$ is the perturbed number density varying as $\exp(i\vec{k}\cdot\vec{r} - i\omega t)$. The integral is obviously taken over the slow wave spectrum. Using Eqs. (6) and (7), we can write

$$\vec{F}'^*_k = i\vec{k} f'^*(\vec{k}')\phi^*_k, \tag{25}$$

and

$$\frac{\partial}{\partial \vec{k}'} \tilde{N}_k(\vec{k}') = -\frac{\partial}{\partial \vec{k}'} f'(\vec{k}')\phi_k \frac{\vec{k}\cdot\partial N_0/\partial \vec{k}'}{(\omega - \vec{k}\cdot\vec{v}')}. \tag{26}$$

Assuming that $f'(\vec{k}')$ is a very slow function of $\vec{k}'$, then we can write Eq. (24) in the form of a diffusion equation

$$\frac{\partial}{\partial t} N_0(\vec{k}') = \frac{\partial}{\partial \vec{k}'} \cdot \mathbf{D}(\vec{k},\vec{k}') \cdot \frac{\partial}{\partial \vec{k}'} N_0(\vec{k}'), \tag{27}$$

where the diffusion tensor in the quasiparticle momentum space is determined by

$$\mathbf{D}(\vec{k},\vec{k}') \simeq i|f'(\vec{k}')|^2 \int \vec{k}\vec{k} \frac{|\phi_k|^2}{(\omega - \vec{k}\cdot\vec{v}')} d\vec{k}. \tag{28}$$

Such a diffusion is a statistical consequence of the acceleration and deceleration of the individual quasiparticles (or high-frequency wave packets), resonantly interacting with the different components $(\omega, \vec{k})$ of the slow wave spectrum.

VIII. SPECIFIC EXAMPLES

We illustrate the above general formalism by considering two specific examples. The first one corresponds to the ion acoustic waves moving in a plasmon gas. More specifically, we assume an isotropic plasma with a broadband electron plasma wave turbulence [3]. In this case, we have the usual dielectric function

$$\epsilon(\omega, \vec{k}) = 1 - \frac{k^2 v_s^2}{\omega^2(1 + k^2\lambda_D^2)}, \tag{29}$$

where v_s is the ion acoustic velocity and λ_D is the electron Debye length. The plasmon gas can be characterized by the following expressions for the velocity and force

$$\vec{v} = 3v_{\text{the}}^2 \frac{\vec{k}'}{\omega'}, \quad \vec{F}' = -\frac{1}{2\omega'} \frac{e^2}{\epsilon_0 m_e} \frac{\partial \tilde{n}_e}{\partial \vec{r}}. \tag{30}$$

Here, $\tilde{n}_e$ is the perturbed electron plasma number density which, for the ion acoustic waves, can be related with the potential perturbation by $\tilde{n}_e = \epsilon_0 \phi/(e\lambda_D^2)$. This leads to a force given by Eq. (7) with

$$f'(\vec{k}') = \frac{1}{2\omega'} \frac{e}{m_e \lambda_D^2}, \tag{31}$$

where $\omega' \simeq \omega_{\text{pe}}$. On the other hand, the calculation of χ_{qp} leads to

$$f(\vec{k},\vec{k}') = \frac{k^2\omega_{\text{pe}}^2}{\omega^2 n_0 m_i (1 + k^2\lambda_D^2)^2 \omega'^2}. \tag{32}$$

This completely characterizes the problem. For example, the bounce frequency ω_b of a trapped plasmon in the potential well of the ion acoustic waves will be given by

$$\omega_b \simeq k \frac{\omega_{\text{pe}}}{\omega'} \left(\frac{e\phi_0}{2m_e}\right)^{1/2}. \tag{33}$$

Comparing this with the electron bounce frequency $\omega_{\text{be}} = (e\phi_0/m_e)^{1/2}$, we conclude that the plasmon behaves as a particle with the electron charge and an effective mass equal to $(2m_e/k^2)$, where we have assumed $\omega' = \omega_{\text{pe}}$.

Our second example will be that of a zonal flow in a drifton gas (or a broadband drift wave turbulence). In this case, we have a pseudo-three-dimensional model in the plan (r,θ) perpendicular to the toroidal magnetic field B_0, which can be described by a modified Hasegawa-Mima equation [6,7]. In this case, we have $\epsilon(\omega,\vec{k}_\perp)=1$, which means that there is no linear dispersion relation for the zonal flows. However, the quasiparticle susceptibility χ_{qp} associated with the drifton gas allows for the existence of a nonlinear dispersion relation determined by the function

$$f(\vec{k},\vec{k}'_\perp)=-\frac{k^2v_s^2}{\omega}\frac{k_\theta'^2\rho_s^2k_r'}{(1+k'^2\rho_s^2)},\tag{34}$$

where $\rho_s=(v_s/\omega_{ci})$ is the ion acoustic Larmor radius. The electron drift wave packets, or driftons in our quasiparticle description, can be characterized by the dispersion relation

$$\omega'=k'_\theta\left(V_0+\frac{V_*}{1+k_\perp'^2\rho_s^2}\right),\tag{35}$$

where $V_0=ck\phi_0/B_0$ and V_* is the electron diamagnetic drift velocity. This means that the force acting on the driftons is

$$\vec{F}'=-\frac{\partial\omega'}{\partial\vec{r}_\perp}=-f'(\vec{k}'_\perp)\frac{\partial}{\partial\vec{r}_\perp}\phi,\tag{36}$$

where $\phi(\vec{r}_\perp,t)$ is the slow potential perturbation associated with the zonal flow and

$$f'(\vec{k}_\perp)=\frac{ck}{B_0}k'_\theta.\tag{37}$$

Here, again, the functions $f(\vec{k},\vec{k}'_\perp)$ and $f'(\vec{k}_\perp)$ will completely characterize our problem and allow us to study the various kinetic and hydrodynamic beam instabilities, as well as to establish the values of the bounce frequency and the quasiparticle diffusion.

IX. CONCLUSIONS

The definition of a quasiparticle susceptibility, and the existence of resonant interactions between quasiparticles and slow waves, leads to a general view of the plasma and fluid turbulence where long wavelength perturbations can be Landau damped by the turbulent medium. Examples of this type of general behavior are the cases of the photon Landau damping of relativistic electron plasma waves, the plasmon Landau damping of the ion acoustic waves, the phonon Landau damping of dust lattice waves in dusty plasmas or in ordinary liquids, or the drifton Landau damping of zonal flows in magnetic fusion plasmas. The existence of this wide variety of examples results from the universality of the resonant coupling between large scale and small scale structures in fluids and plasmas, and it can be seen as different manifestations of anomalous viscosity associated with the turbulent state.

On the other hand, for appropriate turbulent spectra (or quasiparticle distributions), resonant damping can be replaced by resonant amplification, and consequently lead to instabilities. This could be described as an anomalous negative viscosity. The excitation of large scale structures by small scale turbulence in fluids has been sometimes described as in terms of a negative viscosity. But here we prefer to relate it to the existence of quasiparticle distributions. The concept of quasiparticle is indeed one of the basic concepts of the field theory and its importance to plasma physics is stressed here. The resonant interactions between these particles and wave perturbations are efficient channels for the energy exchange between small and large scale events in a fluid without the need for energy cascading. Their equations of motion describe the energy exchange of these quasiparticles with the medium, and show that quasiparticle acceleration and trapping by waves and by moving large scale perturbations of the medium can eventually take place. Photon acceleration by electron plasma waves and by relativistic ionization fronts [14] are nothing but a particular and somewhat spectacular example of this much larger concept.

Furthermore, beams of nearly monoenergetic quasiparticles, with distributions of the type $N(\vec{k}')=N_0\delta(\vec{k}'-\vec{k}'_0)$, can excite waves as these move through the medium, in the same way as electron beams or laser pulses (photon beams) can excite electron plasma waves. This can be seen as quasiparticle beam instabilities, and several examples have been studied [1,3,6]. Finally, mention should be made to the derivation of quasilinear equations describing diffusion in quasiparticle phase space, which can answer the question of the global energy transfer between the large and the small scales of plasma perturbations, for given initial conditions. Particular examples are the expressions for the photon diffusion coefficient when interacting with an electron plasma wave spectrum [1], or the plasmon diffusion coefficient in the plasmon wave number space, expressed in terms of the phonon distribution [3]. Similar expressions could be derived for different kinds of drift wave turbulence interacting with a spectrum of shear flows or zonal flows, as shown by the above general formalism.

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