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Spontaneous generation of plasmons by ballistic electrons

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A beam of ballistic electrons moving with a velocity of about twice the Fermi velocity ($\sim 10^7$ cm/s) with respect to a stationary electron gas is shown to lead to a spontaneous generation of plasmons.

Recent advances in semiconductor technology have led to systems where ballistic propagation of electrons has been achieved over distances up to 100 μm in very-high-mobility ($\sim 9\times 10^6$ cm^2/V s), modulation-doped heterojunctions, at relatively high areal electron densities ($\sim 2\times 10^{11}$ cm^{-2}).[1] In such systems, electrons propagate over a large distance without incurring a single scattering event, and phenomena analogous to the traditional geometrical optics, such as refraction, focusing, etc., can occur.[1] In this paper we investigate the possibility of a spontaneous generation of plasmons in such systems. This effect, which might occur at sufficiently high drift velocities of electrons (of the order of the Fermi velocity), would destroy the ballistic motion, by opening an effective electron-plasmon scattering channel. This phenomenon, in turn, should be experimentally detectable in the type of experiments of Ref. 1.

Generation of plasma waves by a current, through current-driven plasma instabilities, is well known in gaseous plasmas.[2,3] It has been shown that analogous plasma instabilities can exist in solid-state systems.[4-8] The drift velocities of carriers needed for this phenomenon to occur in these systems are usually very high, of the order of the Fermi velocity v_F. Although the high-mobility (modulation-doped) semiconductor heterostructures can provide sufficiently large velocities of carriers (at high carrier concentrations), the effectiveness of the current–plasma-wave energy transfer is limited due to an intrinsic absorption associated with the carrier-phonon or carrier-impurity scatterings. On the other hand, these scatterings do not occur, of course, in the ballistic mode of operation, and therefore the current-driven plasma instabilities might occur in such systems more readily.

To study current-driven plasma instabilities in a ballistic system, we consider a stream of ballistic electrons, modeled by a two-dimensional (2D) electron-gas layer with the susceptibility $\chi^b(n_b,q,\omega)$ moving against a uniform 2D background of stationary electrons, with susceptibility $\chi^{\rm st}(n_{\rm st},q,\omega)$. n_b, $n_{\rm st}$, q, and ω are, respectively, electron densities, and the wave vector and frequency of the electromagnetic radiation. In the nonretarded ($c\to\infty$) and long-wavelength ($q\to 0$) limits, the dielectric function of the system is simply given by

$$\epsilon(\mathbf{q},\omega)=1+4\pi\chi^b(n_b,q,\omega)+4\pi\chi^{\rm st}(n_{\rm st},q,\omega)\ , \tag{1}$$

where the stationary susceptibility is given in the random-phase approximation by

$$\chi^{\rm st}=-\frac{e^2}{\kappa q}\int\frac{d\mathbf{p}}{(2\pi)^2}\,\frac{f(\mathbf{p}+\mathbf{q})-f(\mathbf{p})}{(2\mathbf{p}\cdot\mathbf{q}+q^2)\hbar^2/2m-\hbar\omega}\ , \tag{2}$$

$\mathbf{p}$ is the electron wave vector, and m is the effective mass of the electron, and $f(\mathbf{p})$ is the electron distribution function. We note that, in the long-wavelength limit ($q\to 0$), $\epsilon(\mathbf{q},\omega)$ does not depend on the separation between the two layers.

In the scattering-free environment, the electron distribution functions of the stationary and uniformly moving electrons [$f^{\rm st}(\mathbf{p})$ and $f^b(\mathbf{p})$, respectively] are related by $f^b(\mathbf{p})=f^{\rm st}(\mathbf{p}-m\mathbf{v}_{\rm dr}/\hbar)$, where $\mathbf{v}_{\rm dr}$ is the drift velocity of the electrons. This, in turn, leads to a Doppler shift in the susceptibility,[9] $\chi^b(n_b,q,\omega)=\chi^{\rm st}(n_b,q,\ \omega-qv_{\rm dr})$. Therefore Eq. (1) becomes

$$\epsilon(\mathbf{q},\omega)=1+4\pi\chi^{\rm st}(n_b,q,\ \omega-qv_{\rm dr})+4\pi\chi^{\rm st}(n_{\rm st},q,\omega)\ . \tag{3}$$

For $T=0$, the susceptibility of the stationary plasma is given by the well-known Stern formula:[10]

$$\chi^{\rm st}=\frac{e^2nk_F}{2\varepsilon_F\kappa q^2}\left\{\frac{q}{k_F}-\left[\frac{1}{4}\left[\frac{2\omega}{qv_F}+\frac{q}{k_F}\right]^2-1\right]^{1/2}+\left[\frac{1}{4}\left[\frac{2\omega}{qv_F}-\frac{q}{k_F}\right]^2-1\right]^{1/2}\right\}, \tag{4}$$

where n is the surface density of electrons, κ the dielectric constant, and k_F, v_F, and ε_F are, respectively, the Fermi wave vector, velocity, and energy. The plasma-mode structure of the system is given by the usual condition

$$\epsilon(q,\omega)=0\ . \tag{5}$$

In the driftless case ($v_{\rm dr}=0$), Eqs. (3)–(5) lead to the well-known 2D-plasmon-mode dispersion relation ($q\ll k_F$),[4,10]

$$\omega^2=2\pi ne^2q/m\kappa\quad\text{with } n=n_b+n_{\rm st}\ . \tag{6}$$

Although this is the only solution of Eq. (5), in this case the real part of $\epsilon(q,\omega)$ vanishes also in the vicinity of the line $\omega=qv_F$, i.e., at the edge of the single-particle continuum where the susceptibility [Eq. (4)] experiences rapid variation and changes sign. The imaginary part of $\epsilon(q,\omega)$ at this "spurious mode" is positive, reflecting the absorption associated with the single-particle excitations.

If, however, v_{dr} is nonzero, the electron distribution function of the moving electrons shifts correspondingly, leading to a population inversion and subsequent downward single-particle transitions. These transitions for a sufficiently large drift velocity (threshold drift velocity) produce an energy flow which can overcome the usual absorption due to upward single-particle transitions. Then, Eq. (5) has another solution for which $\mathrm{Re}(\omega)\approx qv_F$ and $\mathrm{Im}(\omega)>0$. This can be analytically illustrated by considering the case of $n_b = n_{\mathrm{st}}$. Then it is straightforward to show that both the real and the imaginary parts of the dielectric function [Eq. (5)] vanish if

$$\omega = qv_F - \hbar q^2/2m + \hbar^2 q^3/2m^2 v_F + O(q^4) \tag{7}$$

and

$$v_{\mathrm{dr}} = 2\omega/q = 2v_F - \hbar q/m + \hbar^2 q^2/m^2 v_F + O(q^3)\ . \tag{8}$$

In this case, Eq. (7) is the dispersion relation of a potentially unstable, plasma mode. For small q this mode is an acoustic one with its phase velocity almost equal to the Fermi velocity. At the threshold drift, ω given by Eq. (7) is real and the amplitude of the mode neither increases nor decreases in time, i.e., the mode is at the onset of instability. For drift velocities larger than the threshold velocities given by Eq. (8), the mode frequency becomes complex with $\mathrm{Im}(\omega)>0$, and the mode amplitude grows in time.

The instability domain can be studied by numerically solving Eq. (5). A typical, complex dispersion for an unstable mode (real and imaginary parts of frequency versus q), for a fixed V_{dr}, is shown in Fig. 1. Here $v_{\mathrm{dr}} = 2.5v_F$, $n_{\mathrm{st}} = n_b = 2\times 10^{11}$ cm^{-2}, and other material parameters are chosen to be those for GaAs: $m = 0.0665m_e$, $\kappa = 13.1$. While the real part of the mode varies linearly with q (with the slope equal to $v_{\mathrm{dr}}/2$), the imaginary part (gain) has a maximum approximately in the middle of the instability domain. The strength of this maximum is a measure of the efficiency of the energy transfer between electrons and plasmons. We have also examined various cases with unequal densities. The highest gains occur for $n_{\mathrm{st}} \approx n_b$.

In the instability domain, electrons are scattered (through plasmon generation) and can no longer be considered ballistic, i.e., their mean free path L rapidly decreases. In fact, one can estimate this decrease by first noting that the effective scattering frequency

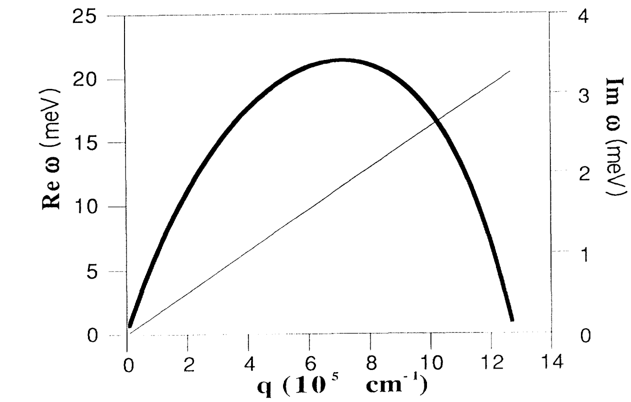

FIG. 1. Complex solution of the dispersion relation [Eq. (3)], in the instability domain. Parameters are $v_{\mathrm{dr}} = 2.5v_F$, $n_{\mathrm{st}} = n_b = 2\times 10^{11}$ cm^{-2}, $m = 0.0665m_e$, $\kappa = 13.1$. The real part of the frequency is shown as a thin line, while the imaginary part of the frequency is represented by a bold line.

$v_{\text{el-pl}} \approx \max\{\mathrm{Im}(\omega)/2\pi\}$. Then from Fig. 1 we find $v_{\text{el-pl}} \approx 10^{11}$ s^{-1}, which corresponds to the mobility 2.7×10^5 cm^2/V s, and $L \approx 4$ μm.

The mean free path in the case of Ref. 1 is 64 μm, and the motion between electrodes (separated by ≈ 25 μm) is ballistic. If an additional layer of stationary electrons is provided in the arrangement of Ref. 1, plasmon generation through the above-mentioned instability mechanism would become feasible. Once the plasmon generation begins, the mean free path shrinks dramatically (to ≈ 4 μm) and ballistic propagation would not be detectable in this configuration, providing a test of this phenomenon.

This basic phenomenon of spontaneous generation of plasmons by ballistic, relative motion of carriers could also be realized in other physical systems in which the threshold condition can be met. This process, when it occurs, can be detected in other ways as well, e.g., through a conversion of the plasma-wave energy to electromagnetic radiation by placing a grating (or other coupler) to quench the excess momentum. We note that such an arrangement has potential device applications as a radiation source.

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