ENERGY SPECTRUM AND ELECTRON SCATTERING PROCESSES IN INVERSION LAYERS

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The carrier plasma in an inversion channel is investigated for quantizing band bending. The screening of the charged-impurity field, the dispersion of the plasma oscillations, the electron spectrum, and the relaxation time for scattering by volume and surface ions are determined.

AN inversion layer is produced near the surface of a doped semiconductor in the presence of a strong electric field directed normal to the surface. If the field is sufficiently strong and corresponds in direction to repulsion of the majority carriers from the surface, then the bottom of the conduction band (in the case of a p-type semiconductor) or the top of the valence band (in the n-type case) can intersect the Fermi level near the surface. Then the near-surface layer in an n-type semiconductor is characterized by p-type conductivity (and vice versa in the case of a p-type semiconductor).

Measurements of the Shubnikov–de Haas effect in inversion channels[1] indicate that at the attainable band bending the quantization of the motion of the electron in a direction normal to the surface becomes significant. This is very important for the interpretation of other experiments (for example, the conductivity, the Hall effect in near-surface layers), and for a general understanding of electronic processes in metal-dielectric-semiconductor layered structures.

The difficulties of constructing the theory of an electron plasma in an inversion layer are connected primarily with the inhomogeneity of the system with respect to the coordinate normal to the surface. Even the well-known summation of Gell-Mann and Brueckner, which leads to a renormalization of the electron-electron and electron-ion vertices, is realizable only for spatially homogeneous systems, in which the Green's functions depend on the difference of the arguments. However, as will be shown below, the situation is facilitated in the limiting case of a very thin near-surface layer, and the theory of a weakly-non-ideal "two-dimensional" plasma can be constructed after the model of the known three-dimensional problem (see, for example,[2]). We shall be interested in the criterion for the aforementioned two-dimensional character, the renormalization of the Coulomb interaction and the spectrum of the plasma oscillations, the electron spectrum, and the relaxation time for scattering by ions.

We start with a consideration of the nonrenormalized electron-electron and electron-ion vertices. Since we are dealing with charges in a semiconductor with dielectric constant ϵ_1, bordering on a dielectric (ϵ_2), it is necessary to take into account the electrostatic image forces. Obviously, the interaction energy of two electrons located at the points r_1 and r_2 depends on the differences $x_1 - x_2$ and $y_1 - y_2$ (the z axis is perpendicular to the surface). The Fourier component of this interaction with respect to the coordinates along the surface is equal to

$$\varphi_{ee}(\mathbf{k}; z_1, z_2) = \frac{2\pi e^2}{\varepsilon_1 k}\left[e^{-k|z_1-z_2|} + \frac{\varepsilon_1-\varepsilon_2}{\varepsilon_1+\varepsilon_2}e^{-k(z_1+z_2)}\right], \qquad (1)$$

$\mathbf{k}$ is a two-dimensional vector and the semiconductor occupies the region $z > 0$. We expand φ_{ee} in terms of the eigenfunctions of the transverse motion of the electron $\psi_n(z)$. As will be shown subsequently, the significant values of k are of the order of e^2/ϵ_1 (we use a system of units with $\hbar = m^* = 1$), i.e., k^{-1} is of the order of the effective Bohr radius a_0^* in the volume of the semiconductor. If the thickness of the layer in which the electrons are mainly concentrated is much smaller, then $\varphi_{ee}(\mathbf{k}; z_1, z_2)$ does not have time to change appreciably over distances characteristic of the wave functions $\psi_n(z)$. In this case the principal role is played by the diagonal elements of φ_{ee} with respect to the quantum numbers of the transverse motion

$$\langle n'm'|\varphi_{ee}(\mathbf{k})|nm\rangle \approx \frac{2\pi\tilde{e}^2}{k}\delta_{nn'}\delta_{mm'}, \quad \tilde{e}^2 = \frac{2e^2}{\varepsilon_1+\varepsilon_2}. \qquad (2)$$

The indicated requirement is apparently not exceedingly stringent if it is recognized that the channel thickness can reach 30–40 Å, and a_0^* amounts to ~ 100 Å in GaAs and ~ 500 Å in InSb. Thus, the scattering of electrons with transitions between the subbands of transverse quantization ($n \neq n'$, $m = m'$) is weakened with respect to the parameter $a/a_0^* \ll 1$, where a is the effective thickness of the conducting layer. This parameter determines the criterion of two-dimensionality of the problem and will be used in subsequent calculations.

The polarization of the medium also leads to the occurrence of single-electron terms in the potential energy; these terms describe attraction or repulsion (depending on the sign of $\epsilon_1 - \epsilon_2$) of the electron at the surface. These terms can be regarded as a renormalization of the external field, leading to a bending of the bands, and need not be considered separately. In the present article we shall not be interested in the concrete form of the near-surface potential well and the wave functions $\psi_n(z)$, since all the results will be expressed in terms of the number of electrons per unit area of the inversion layer, which can be measured directly in the experiment.

To find the electron-ion vertex part φ_{ei} we start again from expression (1), in which z_2 is regarded as the coordinate of the ion. We shall assume the ions to be infinitely heavy impurities and average the results over the impurity configurations. The quantity $\langle n'|\varphi_{ei}|n\rangle$ for the ion located at the point ρ_j, z_j is equal to

$$\langle n'|\varphi_{ei}(\mathbf{k}; \rho_j, z_j|n\rangle \approx 2\pi\tilde{e}^2k^{-1}\exp(i\mathbf{k}\rho_j - kz_j)\delta_{nn'}. \qquad (3)$$

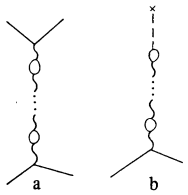

FIG. 1

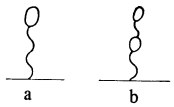

FIG. 2

In calculating the matrix element in (3) we take into account the fact that the wave functions of the electrons cut off the integral at distances on the order of $a \ll k^{-1}$, and the ion coordinate z_j can have arbitrary values.

We now proceed to calculate the renormalized electron-electron and electron-ion vertices Γ_{ee} and Γ_{ei}. We assume that the plasma is weakly nonideal, i.e., the effect of the Coulomb interaction is relatively small. It is well known that to this end the electron gas must be sufficiently dense, so as to satisfy the condition $\tilde{e}^2 \ll p_0$, where p_0 is the Fermi momentum. Since the limiting Fermi momentum is not altered by the particle interaction, it can be calculated by assuming the electron gas to be ideal. Then p_0 is determined from the equations

$$\frac{1}{2}p_0^2 = \mu|_{T=0}, \quad T\sum_m \ln(1+e^{(\mu-\epsilon_m)/T}) = \nu\pi,$$

where ϵ_m are the energy levels of the transverse motion, ν is the number of electrons in the layer per unit area and T the temperature in energy units, and the sum is taken over all the subbands of transverse quantization.

If only one subband is filled, then $p_0^2 = 2\pi\nu$, and the foregoing condition of weak nonideality of the plasma can be written in the form $\nu/\tilde{e}^4 = \nu a_0^{*2} \gg 1$ (i.e., the number of surface electrons on the area of the Bohr orbit should be large).

For a weakly nonideal plasma, the most important processes are those of scattering with a small transfer of momentum $\mathbf{k}$, as seen from (2) and (3). We shall therefore calculate the renormalization of the Coulomb interaction retaining only the principal terms in $\tilde{e}^2/k$. It is known[2] that this problem reduces to a summation of a chain of electron loops strung on the line of the Coulomb interaction (see Fig. 1). Figure 1a pertains to Γ_{ee}, and Fig. 1b to Γ_{ei}. The cross denotes an ion, the dashed line corresponds to the factor φ_{ei}, and the wavy line to Γ_{ee}. The dashed line does not carry frequencies (the ions are infinitely heavy).

Before we calculate Γ_{ee} and Γ_{ei}, let us see how to eliminate in our problem the difficulty connected with divergent diagrams for G-functions of the type shown in Fig. 2. For the spatially-homogeneous problem, the contribution of this diagram vanishes because of the plasma electroneutrality condition (see[2]). In the case considered by us, the situation is somewhat different. The sum of all the irreducible self-energy parts of the type under consideration is equal to

$$\Sigma'(\mathbf{r}-\mathbf{r}', t-t') = 2\tilde{e}^2\int [G_e(\mathbf{r}_1, 0; \mathbf{r}_1, -0) + G_i(0, -0)]$$
$$\times \varphi(\mathbf{r}, \mathbf{r}_1)\, d\mathbf{r}_1 \delta(\mathbf{r}-\mathbf{r}')\delta(t-t') \equiv U(\mathbf{r})\delta(\mathbf{r}-\mathbf{r}')\delta(t-t').$$

The quantity in the square brackets under the integral sign is obviously the total charge density of the ions and electrons, with the latter depending on the coordinates because of the action of the external field. $U(\mathbf{r})$ is thus equal to the additional field produced at the point $\mathbf{r}$ by the inhomogeneous distribution of the electrons. The structure of the quantity $\Sigma'(\mathbf{r}-\mathbf{r}', t-t')$, and namely its proportionality to $\delta(\mathbf{r}-\mathbf{r}')\delta(t-t')$, signifies that in the Schrödinger equation for one electron the quantity $U(\mathbf{r})$ is added to the external potential, i.e., a self-consistent field is produced. Consequently, diagrams of type of Fig. 2 can be disregarded if it is assumed that the single-electron levels and wave functions are known from a solution of the self-consistent problem. This is precisely what we shall assume from now on.

The calculation of Γ_{ei} is in essence the problem of screening of the field of a static impurity. The series represented by the diagram of Fig. 1b can be summed with the aid of the temperature-diagram technique when $\langle n|\varphi_{ee}|m\rangle$ is proportional to δ_{nm}, i.e., in the limit when $a \ll a_0^*$. In this case it reduces to a geometrical progression, and we obtain (for $k \ll p_0$)

$$\Gamma_{ei}(\mathbf{k}; nn', j) = 2\pi\tilde{e}^2\delta_{nn'}\exp(i\mathbf{k}\rho_j - kz_j)\left[k + 2\sum_m \tilde{e}^2 f(\epsilon_m)\right]^{-1}, \quad (4)$$

where $f(\epsilon_m) = (1+e^{(\epsilon_m-\mu)/T})^{-1}$. At $T = 0$ the screening constant in the denominator of (4) is equal to $2\tilde{e}^2N(\mu)$, where $N(\mu)$ is the number of levels ϵ_m lying below the Fermi level. Thus, the screening radius experiences jumps when the new subband begins to be filled[1]. For a nondegenerate gas filling only the first subband, the reciprocal screening radius is equal to $2\pi\tilde{e}^2\nu/T$. It is seen from (4) that the renormalization affects momenta $k \lesssim \tilde{e}^2$; this justifies the assumptions $k \ll p_0$ and $ka \ll 1$ on which the derivation of (2)–(4) is based.

Let us now find the renormalized electron-electron vertex $\Gamma_{ee}(\omega, \mathbf{k})$. For a strongly degenerate electron gas, under the condition that one subband is filled, we obtain

$$\Gamma_{ee}(\omega, k) = 2\pi\tilde{e}^2/[k + 2\tilde{e}^2\Pi(\omega, k)],$$
$$\Pi(\omega, k) = 1 - |\omega|\left[\frac{\Theta(\omega^2 - k^2p_0^2)}{\sqrt{\omega^2 - k^2p_0^2}} + i\frac{\Theta(k^2p_0^2 - \omega^2)}{\sqrt{k^2p_0^2 - \omega^2}}\right], \quad (5)$$

where $\Theta(x) = 1$ for $x > 0$ and $\Theta(x) = 0$ for $x < 0$.

Neglecting the retardation of the interaction, i.e., for $\omega \ll kp_0$, we obtain from (5) the already known result $\Gamma_{ee} = 2\pi\tilde{e}^2/(k + 2\tilde{e}^2)$. The opposite limiting case is characteristic of plasma oscillations. The pole of the quantity Γ_{ee} determines the dispersion law of the plasmons in the two-dimensional plasma of the inversion layer

$$\omega(k) = \sqrt{\tilde{e}^2p_0^2k}(1 + k/2\tilde{e}^2)(1 + k/4\tilde{e}^2)^{-1/2},$$
$$\omega > kp_0, \quad k \ll p_0. \quad (6)$$

At small k ($k \ll \tilde{e}^2$) we obtain $\omega \approx \sqrt{\tilde{e}^2p_0^2k}$, i.e., the dispersion of the plasmons has no gap, just as in the three-dimensional case. For $k \gg \tilde{e}^2$ we obtain $\omega \approx kp_0$. The characteristic plasma frequency is of the order of $p_0\tilde{e}^2$, which amounts to approximately 10^{13} Hz for $\nu \sim 10^{12}$ cm^{-2} and $(\epsilon_1 + \epsilon_2)/2 \sim 10$.

We proceed now to calculate the Green's function and the energy spectrum of the electrons, with allowance for their Coulomb interaction. The problem consists of cal-

[1] A similar problem of screening of a chared impurity in a thin film was solved by Rytova for the case when one subband is filled [3]. Formula (4) with $N(\mu) = 1$ agrees with Rytova's results.

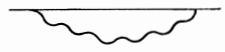

FIG. 3

culating the self-energy part represented by the diagram of Fig. 3. Without allowance for the renormalization of the Coulomb interaction, this quantity is equal to

$$\Sigma_0(\mathbf{p}) = i2\pi\tilde{e}^2 \int \frac{G(\omega_1, \mathbf{p}_1)e^{i\omega_1 t}}{|p-p_1|}\frac{d\mathbf{p}_1 d\omega_1}{(2\pi)^3} \qquad (7)$$

$$= \begin{cases} -\dfrac{2\tilde{e}^2 p_0}{\pi} E\left(\dfrac{p}{p_0}\right) & \text{for } p < p_0, \\ -\dfrac{2\tilde{e}^2 p}{\pi}\left[E\left(\dfrac{p_0}{p}\right) - \left(1-\dfrac{p_0^2}{p^2}\right)K\left(\dfrac{p_0}{p}\right)\right] & \text{for } p > p_0, \end{cases}$$

where K and E are respectively complete elliptic integrals of the first and second kinds. The value of $\Sigma_0(p)$ at $p = p_0$ determines the renormalization of the chemical potential $\Delta\mu = -2\tilde{e}^2 p_0/\pi$. Expression (7) must be corrected in the region $(p - p_0) \lesssim \tilde{e}^2$, for in this case small momentum transfers in the vertex become significant. It is necessary to calculate $\Sigma(p)$ using formula (5) for Γ_{ee}. Omitting the rather cumbersome manipulations, we present the final result:

$$\mathrm{Re}\,\Sigma(p) = \Sigma_0(p) - \frac{2\tilde{e}^4}{\pi}\mathrm{sign}\,\xi(p)\left[\pi\ln\left(1+\frac{v}{2}\right)\right.$$

$$\left. + \int_v^\infty \left(\mathrm{arctg}\frac{v}{\sqrt{x^2-v^2}} + \mathrm{arctg}\frac{v}{(1+x)\sqrt{x^2-v^2}}\right)\frac{dx}{x+2}\right],$$

$$\mathrm{Im}\,\Sigma(p) = \frac{2\tilde{e}^4}{\pi}\mathrm{sign}\,\xi(p)\left[\int_0^v \frac{\ln(1+x)}{x+2}dx\right. \qquad (8)$$

$$\left. - \frac{1}{2}\int_v^\infty \ln\left(1-\frac{v^2(x+2)}{x(x+1)^2}\right)\frac{dx}{x+2}\right],$$

$$\xi(p) \equiv \frac{1}{2}(p^2-p_0^2), \quad v \equiv \frac{|\xi|}{2\tilde{e}^2 p_0}.$$

From this we get the energy spectrum $\epsilon(p)$ and the damping $\gamma(p)$ of the electronic excitations in the region $|\xi| \ll \tilde{e}^2 p_0$:

$$\varepsilon(p) = \xi(p)\left[1+\frac{\tilde{e}^2}{\pi p_0}\left(\ln\frac{2p_0}{\tilde{e}^2}-2\right)\right],$$

$$\gamma(p) = \frac{\xi|\xi|}{4\pi p_0^2}\left(\ln\frac{2p_0\tilde{e}^2}{|\xi|}-\frac{1}{2}\right) \qquad (9a)$$

and in the region $p_0\tilde{e}^2 \ll |\xi| \ll p_0^2$:

$$\varepsilon(p) = \xi(p)\left[1+\frac{2\tilde{e}^2}{\pi p_0}\left(\ln\frac{2\sqrt{2}p_0}{\sqrt{|\xi|}}-\frac{1}{2}\right)\right],$$

$$\gamma(p) = \frac{2\tilde{e}^4}{\pi}\mathrm{sign}\,\xi\left(\ln^2\frac{|\xi|}{2p_0\tilde{e}^2}-1+\frac{\pi^2}{24}\right). \qquad (9b)$$

In formulas (9), the energy $\epsilon(p)$ is reckoned from the renormalized chemical potential. From (9a) we can find the change of the effective mass of the electron as a result of the interelectron interaction

$$m_{\mathrm{eff}}^{-1} = m^{*-1}\left[1+\frac{\tilde{e}^2}{\pi p_0}\left(\ln\frac{2p_0}{\tilde{e}^2}-2\right)\right],$$

where m* is the usual effective mass in the volume (outside the inversion channel). Thus, the longitudinal effective mass of the electron in the inversion layer is smaller than m* and increases, tending to m*, with increasing p_0, i.e., with increasing transverse electric

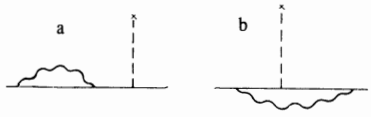

FIG. 4

field. It should be noted that the dependence of m_{eff} on the transverse field may also be connected with another circumstance. The deviation of the nonrenormalized dispersion law $\epsilon_0(p)$ from parabolic also changes the effective mass at the Fermi level with changing field. It is natural to assume that true $\epsilon_0(p)$ curve lies lower than the parabola $p^2/2m^*$. Therefore m_{eff} will increase with increasing p_0, i.e., with increasing transverse field. The order of magnitude of this effect can be estimated as the ratio of p_0 to the reciprocal-lattice vector. For $\nu \sim 10^{12}$ cm^{-2} this ratio is several percent. The parameter $\tilde{e}^2/\pi p_0$ at this value of ν, $(\epsilon_1 + \epsilon_2)/2 \sim 10$, and $m^* \sim 0.1\, m_e$ is equal to 0.06. Thus, both effects are of comparable order of magnitude and act in the same direction; m_{eff} increases with increasing external field.

Finally, let us calculate the time of relaxation of the inversion-layer electrons due to scattering by ions. In the calculation of the scattering matrix element it is necessary to replace the quantity φ_{ei} by the renormalized vertex Γ_{ei}. In addition, the question arises of allowance for the corrections represented by the diagram of Fig. 4. Figure 4a corresponds to replacement of the nonrenormalized G function by the quantity $(G_0^{-1} - \Sigma)^{-1}$. Figure 4b, as can be verified, makes a contribution that is small in terms of the parameter $\tilde{e}^2/p_0$ compared with 4a, owing to the limitations on the angles between the momenta.

The remaining procedure is perfectly analogous to the calculation of the residual resistance of metals with the aid of the technique of Abrikosov and Gorkov[2]. The only difference is that the equation for the Green's function averaged over the positions of the impurities contains the quantity $(G_0^{-1} - \Sigma)^{-1}$ in place of G_0. This means that in addition to impurity damping there is an intrinsic damping in the electronic Green's function, connected with the electron-electron scattering and described by the quantity $\gamma(p)$ in formulas (9). It will be shown below that in the significant region of values of $\xi(p)$ ($\xi\tau_v \sim 1$), where τ_v is the time of relaxation on the ions, the intrinsic damping is small compared with $1/\tau_v$. This enables us to write down immediately the result for the conductivity, for in the case when $\gamma(p) \ll 1/\tau_v$ the problem reduces to that solved in[2]:

$$\frac{1}{\tau_v(p)} = \frac{2\pi}{S}\sum_j |\Gamma_{ei}(\mathbf{p}-\mathbf{p}')|^2\langle\exp(-2|\mathbf{p}-\mathbf{p}'|z_j)\rangle \qquad (10)$$

$$\times(1-\cos\hat{\mathbf{p}}\mathbf{p}')\delta\left(\frac{p'^2}{2}-\frac{p'^2}{2}\right)d\mathbf{p}'.$$

Here S is the surface area of the sample, the angle brackets denote averaging over the coordinate of the ion, and the summation is over all ions. Substituting in (10) the expression for Γ_{ei} (see (4)) and assuming the ions to be distributed in the volume of the sample with constant density n_v, we obtain

$$\frac{1}{\tau_v(p)} = \frac{\pi\tilde{e}^4 n_v}{p^3}\int_0^{\pi/2}\frac{d\alpha\sin\alpha}{(\sin\alpha+\tilde{e}^2/p)^2} \approx \frac{\pi\tilde{e}^4 n_v}{p^3}\left(\ln\frac{2p}{\tilde{e}^2}-1\right). \qquad (11)$$

The last equation in (11) was obtained under the condi-

tion $\tilde{e}^2 \ll p$, since the conductivity of a strongly degenerate gas is determined by $\tau_V(p)$ at $p = p_0$.

We can now verify the statement made above concerning the intrinsic damping of the electronic excitations. In both limiting cases (9a) and (9b) the ratio of τ_V to $1/\gamma(p)$ in the region $\xi \sim 1/\tau_V$ turns out to be much smaller than unity. The total electric current through the layer, per unit length, is

$$j_v = 2\sqrt{2\pi}\, e^2 \nu^{1/2} F / n_v \tilde{e}^4 \left(\ln \frac{2\sqrt{2\pi\nu}}{\tilde{e}^2} - 1 \right), \qquad (12)$$

where F is the drawing electric field (along the surface). The scattering centers can be located not only in the interior of the semiconductor, but also on the very boundary separating the semiconductor from the dielectric. If the surface concentration of such ions is n_S, then it is necessary to average expression (10), taking the ion density in the form $n_V + \delta(z) n_S$. Then the total relaxation time is determined from the formula

$$\frac{1}{\tau} = \frac{1}{\tau_v} + \frac{1}{\tau_s}, \quad \frac{1}{\tau_s(p)} = \frac{2\pi^2 \tilde{e}^4 n_s}{p^2}. \qquad (13)$$

We note that in the calculations of $\tau_S(p)$ we can neglect the screening in the denominator of Γ_{ei}, since $\tilde{e}^2 \ll p$ and the transport cross section for scattering by a surface center converges even without allowance for the screening. This constitutes the difference from scattering by an ion in the interior, where the transport cross section diverges logarithmically in the small-angle region if screening is neglected. The current over the layer per unit length equals in the case $\tau_S \ll \tau_V$

$$j_s = e^2 \nu^2 F / \pi \tilde{e}^4 n_s. \qquad (14)$$

The relative role of the volume and surface ions depends on the ratio of the quantities $\nu^{3/2}/n_V$ and ν/n_S. For example, for $\nu \sim 10^{12}$ cm^{-2} and $n_V \sim 10^{16}$ cm^{-3}, the surface scattering prevails over volume scattering at $n_S > 10^{10}$ cm^{-2}. In the general case, as seen from (12) and (14), the dependence of the effective surface mobility on ν is intermediate between power-law relations with exponents 1 and 3/2. This conclusion is in qualitative agreement with experiment (see[4]).

All the formulas obtained above pertain to a degenerate plasma. In the Boltzmann case j_V is proportional to

$$\frac{\nu T^{3/2}}{n_v} \left(\ln \frac{T}{\tilde{e}^2} \right)^{-1},$$

and $j_S \sim \nu T / n_S$, i.e., the surface mobility ceases to depend on the strong transverse field. This also agrees qualitatively with the experimental data.

The problem of carrier mobility in inversion quantizing layers was also discussed by Stern and Howard[5]. Stern and Howard made a variational calculation of the self-consistent problem for the determination of the form of the potential well and of the wave functions $\psi_n(z)$. They also calculated the surface mobility as a function of ν but (inasmuch as computer calculations were involved) they represented their results in the form of plots, making a comparison with the present work difficult. As seen from the foregoing, in the case of very thin inversion layers it is possible to obtain rather simple analytic expressions for τ_S and τ_V without knowing the concrete form of the potential and of the wave functions. The results are expressed in terms of the concentration in the inversion layer, which can be determined directly from experiment.

In conclusion, I am grateful to É. G. Batyev for useful discussions and to S. P. Sinitsa for numerous discussions on problems touched upon in the paper and for the opportunity of becoming acquainted with the results of his experiments prior to publication.

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