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LETTER TO THE EDITOR

Quantum conductance of a lateral microconstraint in a magnetic field

L I Glazman†‡ and A V Khaetskii§

† Institute of Theoretical Physics, Chalmers University of Technology, S-41296, Goteborg, Sweden

§ Institute of Problems of Microelectronics Technology and High-Purity Materials, USSR Academy of Sciences, 142432 Chernogolovka Moscow District, USSR

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Abstract. The $G(d)$ function for a microconstraint was studied in the presence of a magnetic field (where G and d are respectively the conductance and diameter of a constraint). It is shown that spin splitting of a quantum step in the $G(d)$ function increases with the step number n; hence the spin effects should be more pronounced at greater numbers. The study revealed that the effect of magnetic field on electron orbits does not destroy the adiabatic propagation of an electron wave through the constraint. The adiabaticity condition is maintained if d is small compared to the curvature radius of a constraint. Accordingly, the step–plateau width ratio is small for all values H under consideration. Each width, however, increases parametrically at $r_c(H) \to d$ (r_c is the cyclotron radius). The conductance properties in question are in agreement with the experimental data.

Experiments [1, 2] have revealed quantisation of the conductance of a ballistic microconstraint, defined in 2D electron gas 2DEG. The diameter d of a constraint was controlled by the gate voltage. The electrostatic method of electron channel formation secured the smoothness of the constraint boundaries which allowed the occurrence of plateaus in the $G(d)$ function to be attributed [3] to adiabatic propagation of an electron wave through the constraint. The properties of a constraint may be controlled by both the gate voltage and the magnetic field.

The magnetic field H affects the orbital motion as well as the spin state of an electron. It is convenient to investigate the spin effects when the magnetic field is oriented in the plane of the 2DEG, and the orbital effects are negligible due to spatial quantisation [4]. The experiment in [2] showed that the magnetic field so oriented splits the steps in the $G(d)$ dependence. The step of a height $e^2/\pi\hbar$ was observed to split into two steps with a height $e^2/2\pi\hbar$. Surprisingly, however, the spin splitting occurred only for steps with large numbers.

The effective electron mass in a GaAs heterostructure is small. Therefore, if a magnetic field H is applied perpendicular to the plane of the 2DEG, then the magnetic field influence is implied mainly by orbital effects. These effects enhance with increasing d. In the experiment [5] this led to considerable widening of both the plateaus and the

‡ Permanent address: Institute of Problems of Microelectronics Technology and High-Purity Materials, USSR Academy of Sciences, 142432 Chernogolovka, Moscow District, USSR.

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steps. The steps, nevertheless, remained clearly shaped when the cyclotron radius r is greater than $d/2$.

This Letter presents a theoretical study of the effect of magnetic field on the $G(d)$ dependence. Firstly, it will be shown that the value of spin splitting of a step (at a given value of H) increases with the step number n faster than its width. Therefore, the spin effects are more clearly defined for large n. Hybridisation of the spatial and the magnetic quantisation terms occurs if the magnetic field H is perpendicular to the plane of the 2DEG, which changes the $G(d)$ function. However, it will then be shown, that the magnetic field does not destroy the adiabatic passing of an electron wave through the constraint. The only condition required is the smallness of d compared with the curvature radius R of the channel. When magnetic field is applied, the step–plateau width ratio remains small if the condition $d/R \ll 1$ is maintained. Yet both widths increase parametrically when r tends to $d/2$.

Consider the case when magnetic field is oriented in the plane of the 2DEG and does not affect the orbital motion of an electron. The Hamiltonian of a system with the spin–orbit interaction is

$$\hat{\mathcal{H}} = \hat{\mathcal{H}}_0 + \hat{\mathcal{H}}_1 \qquad \mathcal{H}_0 = p^2/2m \qquad \hat{\mathcal{H}}_1 = \alpha\boldsymbol{\sigma}[\boldsymbol{np}] - \tfrac{1}{2}|g|\mu_B\boldsymbol{\sigma H} \tag{1}$$

where $\boldsymbol{n}$ is the normal to the 2DEG plane, $\boldsymbol{p}$ is the two-dimensional momentum of the electron, μ_B is the Bohr magneton, $\boldsymbol{\sigma}$ are Pauli matrices; for GaAs the g-factor is $g = -0.5$ [6]. Equation (1) should be completed by the boundary conditions for the wave function $\psi_\sigma(x, y)$, namely:

$$\psi_\sigma(x, y)|_{y=\pm d(x)/2} = 0, \tag{2}$$

where the x axis is directed along the channel, $d(x)$ is the channel width in the x section. The effective field corresponding to the characteristic energy of the spin-orbit interaction is $H_0 = 2\alpha p_F/\mu_B g$. In the case of GaAs the value H is of order of 1 T (we have used the value $k_F = 10^6\ \mathrm{cm}^{-1}$ from [1, 2] and $\alpha = 10^5\ \mathrm{cm\ s}^{-1}$ [7]). Since the electron motion transverse to the channel axis differs qualitatively from that along the channel, this would be expected to give rise to the anisotropy associated with the direction of external magnetic field H for $H \leqslant H_0$. It is easy to prove, however, that no anisotropy arises in the H-dependence of the step position and the spin–orbit interaction may be neglected. Let us use the perturbation theory assuming that $\hat{\mathcal{H}}_1$ is small. To zeroth order the wavefunctions found in the adiabatic approximation are of the form [3]:

$$\psi_{n\sigma}(x, y) = (2/d(x))^{1/2}\sin\{\pi n[y + d(x)/2]/\mathrm{d}(x)\}\psi(x)\chi_\sigma. \tag{3}$$

Calculating the matrix elements of the perturbation operator, we obtain the adiabatic Hamiltonian for the longitudinal motion of an electron. This Hamiltonian can be represented as the operator in spin space

$$\epsilon_n(p_x, x, \hat{\sigma}) = (\pi\hbar n)^2/2md^2(x) + p_x^2/2m + \alpha p_x\hat{\sigma}_y - |g|\mu_B\boldsymbol{\sigma H}/2). \tag{4}$$

The diagonalisation of (4) is straightforward. For example, when $\boldsymbol{H} \parallel x$ we obtain

$$\epsilon_{n\pm}(p_x, x) = (\pi\hbar n)^2/2md^2(x) + p_x^2/2m \pm [(\alpha p_x)^2 + (g\mu_B H/2)^2]^{1/2} \tag{5}$$

The step in the $G(d)$ dependence occurs when the mode $\epsilon_{n\pm}$ is switched on. Switching takes place when $E_F = \min_{p_x} \epsilon_{n\pm}(p_x, 0)$. This means, that the longitudinal velocity of an electron with the Fermi energy vanishes at the point of maximal constriction. It follows from (5) that the threshold values of $z = k_F d(0)/\pi$ are determined by the equation

$$E_F = E_F(n^2/z^2) \pm |g|\mu_B H/2 \tag{6}$$

where the terms of the order of $m\alpha^2$ are neglected. Condition (6) is valid for all the orientations of the field H in the 2DEG plane within the same accuracy. Hence, the energy scale of the spin–orbit interaction in the step position problem is $m\alpha^2$ (but not αp_F). This scale is negligibly small ($\sim 10^{-2}$ K in the case of GaAs). The perturbation theory used here is valid under the condition $d(0) \ll \hbar/m\alpha$; it is well satisfied even for $d(0)$ as large as $d(0) \sim 2 \times 10^3$ Å (this corresponds to $n = 10$). Equation (6) leads to the following value of the spin splitting for the step n:

$$\Delta z_s = |g|\mu_B Hn/2E_F. \tag{7}$$

The quantum width δz of a step [3] increases with its number n as $(n^{1/2}/\pi^2)(\lambda_F/R)^{1/2}$, where λ_F is the Fermi wavelength of an electron. Therefore, spin splitting is first manifested for the relatively high steps with numbers $n > n_S$:

$$n_S = (2/\pi^2)^2(\lambda_F/R)(E_F/g\mu_B H)^2. \tag{8}$$

For such n the splitting Δz_S exceeds the width δz.

We shall now deal with orbital effects for the magnetic field perpendicular to the plane of the 2DEG. Consider the condition of the adiabatic propagation of an electron wave through the constraint. The sharpness of steps in the $G(z)$ dependence at $H = 0$ is caused by the large value of the ratio $R/d(0)$ (R is the curvature radius of channel boundaries in the region of the constriction). Now our aim is to find the relation between the step width δZ and R in the presence of a magnetic field. This is reduced to the problem of the adiabatic motion of an electron in the presence of magnetic field.

We first determine the adiabatic terms in the presence of a magnetic field. Under the Landau gauge the orbital part of the Hamiltonian has the form†

$$\mathscr{H} = -(\hbar^2/2m)[\partial^2/\partial y^2 + (\partial/\partial x - \mathrm{i}\, eH_y/c\hbar)^2]. \tag{9}$$

The smooth variations of the channel width lead to the slow dependence of the longitudinal motion momentum p on x. Therefore, the wave functions of the adiabatic terms are of the form

$$\psi_l(x, y) = \exp\left(\frac{\mathrm{i}}{\hbar}\int_{\varphi_l}^{x} p_l(x_1)\,\mathrm{d}x_1\right)\varphi_l(y, p_l(x), d(x)) \tag{10}$$

where the functions φ_l are the normalised solutions of the one dimensional boundary problem

$$-(\hbar^2/2m)[\partial^2/\partial y^2 - (p/\hbar - (eH/c\hbar)y)^2]\varphi_l = E\varphi_l \qquad \varphi_l|_{y=\pm d(x)/2} = 0. \tag{11}$$

The boundary problem (11) gives a set of terms $E_l(p, d(x))$. The function $p_l(x)$ in (10) for an electron with energy E is determined by the solution of the equation $E_l(p, d(x)) = E$. Note that for each l there exist two solutions, $p_l^+(x)$ and $p_l^-(x)$, which differ in sign.

† In this section we do not take into account the small spin splitting.

To study the transitions between the adiabatic terms we represent the solution of the Schrödinger equation in the form

$$\psi(x, y) = \sum_r \sum_l c_l^r(x)\psi_l(x, y) \tag{12}$$

where the index $r = \pm$ marks the states with different p_l signs. The function $c_l^r(x)$ varies on the spatial scale that is obviously to coincide with the scale of the function $d(x)$. This scale is equal to $(Rd(0))^{1/2} \gg d(0)$. We determine the amplitude of the transition to the term (l, r) from the initial term (l_0, r_0): $c_l^r(x \to -\infty) = \delta_{ll_0}\delta_{rr_0}$. For this purpose we substitute (12) into the Schrödinger equation with Hamiltonian (9). Using the smoothness of the dependence of functions $c_l^r(x)$, $\varphi_l(x, y)$ upon the x coordinate, at $(l, r) \neq (l_0, r_0)$ we obtain

$$c_l^r(x) = \int_{-\infty}^{x} \mathrm{d}x_1 f_{ll_0}^{rr_0}(x_1) \exp\left(\frac{\mathrm{i}}{\hbar}\int^{x_1} \mathrm{d}x_2[p_{l_0}^{r_0}(x_2) - p_l^r(x_2)]\right) \tag{13}$$

where the function $f(x)$ is composed of scalar products of the type $(\varphi_l(y, x), (\partial/\partial x)\varphi_{l_0}(y, x))$, i.e. the spatial scale of variation of $f(x)$ is $(Rd(0))^{1/2}$ which is large compared to $d(0)$. It is obvious from (13) that small excesses $\delta\varepsilon = E - E_{l_0}(0, d(0))$ of an electron energy E over the threshold value $E_{l_0}(0, d(0))$ are most dangerous for adiabatic propagation. The maximal transition amplitude corresponds to $l = l_0$, $r = r_0$ due to the slowest variation of the exponential factor in (13). The $p_{l_0}(x)$ dependence in (13) can be determined by expansion of $E_{l_0}(p, d(x))$ near the extreme point

$$\delta\varepsilon = -|E_{xx}^{l_0}|x^2/2 + E_{pp}^{l_0}(p^2/2). \tag{14}$$

Here $E_{xx}^{l_0} = \partial^2 E_{l_0}/\partial x^2|_{p,x=0} < 0$, $E_{pp}^{l_0} = \partial^2 E_{l_0}/\partial p^2|_{p,x=0}$. In the case under consideration $p_l^r = -p_{l_0}^{r_0}$; with the help of (14) integral (13) is reduced to the form

$$\int \mathrm{d}\kappa\, F(\kappa) \exp\{(\mathrm{i}/\hbar)[E_{pp}^{l_0}E_{xx}^{l_0}]^{-1/2}\delta\varepsilon[\sinh(2\kappa) + 2\kappa]\}.$$

Application of the stationary phase method leads to the following result for the probability $R = |c_l^{-r_0}|^2$ of the transition $(l, r_0) \to (l, -r_0)$:

$$R_{l\to l} \sim \exp(-\delta\varepsilon/\Delta_l) \qquad \Delta_l = (\hbar/2\pi)(E_{pp}^l|E_{xx}^l|)^{1/2}. \tag{15}$$

This transition corresponds to the reflection of an electron from the constriction into the same mode. The probabilities of transitions into other modes are parametrically smaller than $R_{l\to l}$. Therefore, the transmission coefficient is equal to $T_l = 1 - R_{l\to l}$ and the decrement Δ_l determines the energy dependence of $T_l(\delta\varepsilon)$. Thus, the width of the steps that occur in the conductance with increasing d is equal to $\delta d_l = \Delta_l[\partial E_l(0, d)/\partial d]^{-1}$. Using the expansion $d(x) = d(0) + x^2/R$ and (15), it is easy to obtain

$$\delta z_l = (\hbar/\pi(2\pi)^{1/2})(E_{pp}^l\, \partial z/\partial E_{xx}^l)^{1/2}(k_\mathrm{F}/R)^{1/2} \qquad z = k_\mathrm{F}d(0)/\pi. \tag{16}$$

The above problem is equivalent to that of an electron passing over a parabolic barrier. However, contrary to the case of $H = 0$, the field H affects both the carrier effective mass and the barrier parameters. If $\delta\varepsilon \gg \Delta_l$, then the electron motion is described by the equations of classical mechanics with the Hamilton function $E_l(p, d(x))$. The classical description for the Fermi electron fails in narrow regions of the parameter $\delta\varepsilon < \delta z_l$. These regions correspond to switching of new modes.

We next obtain the dependence of G on z in a weak magnetic field. Let us assume that the cyclotron radius of the Fermi electron $r_c \gg d(0)$, and that the magnetic field can be regarded as perturbation ($r_c = v_F/\omega_c$, $\omega_c = |l|H/cm$ is the cyclotron frequency). In the zeroth approximation the functions φ_n, being the solutions of the boundary problem (11), coincide with the functions $\sin\{\pi n[y + d(x)/2]/d(x)\}$ from (3). Using these functions to calculate the necessary matrix elements of the H-dependent operators in (11), we obtain

$$E_n(0, d(0)) = (\pi\hbar n)^2/2md^2(0) + \tfrac{1}{24}m\omega_c^2 d^2(0), \tag{17a}$$

$$E^n_{pp} = (1/m)(1 + 8d^2(0)/\pi^4 r_c^2). \tag{17b}$$

to an accuracy of $\sim H^2$.

These formulae are valid for $n \gg 1$. To determine the step positions in the $G(z)$ dependence, the $E_n(0, d(0))$ should be set equal to the Fermi energy. The position of the step n is determined by

$$z_n = n + \tfrac{1}{96}\pi^2 n^3(\hbar\omega_c/E_F)^2. \tag{18}$$

Thus, the width of the plateau n is equal to

$$z_{n+1} - z_n = 1 + \tfrac{1}{32}\pi^2 n^2 (H/H^*)^2. \tag{19}$$

We have introduced a characteristic value $H^* = c\hbar k_F^2/2e$ by way of $\hbar\omega_c/E_F = H/H^*$. Under the conditions of validity of the perturbation theory the changes of the plateau widths are small since the second term in (19) is of order of $(d(0)/r_c)^2$. However, accumulation of these changes leads to a considerable shift in the position of steps. The values $z_n - n$ exceed unity for

$$z_n > 6(r_c k_F/3\pi)^{2/3}. \tag{20}$$

The last inequality becomes true even in the weak field region ($z_n \ll r_c k_F$, or $H \ll H^*$) since $r_c k_F \sim E_F/\hbar\omega_c \gg 1$.

We now determine the step widths. Calculating by means of (17a) the quantity $\partial^2 E_n(0, d(0))/\partial z$ and using expressions (16) and (17b) we obtain

$$\delta z_n(H)/\delta z_n(H = 0) = [1 + \tfrac{1}{16}\pi^2(\tfrac{1}{6} + 16/\pi^4)n^2(H/H^*)^2] \qquad H^* = c\hbar k_F^2/2|e|. \tag{21}$$

The magnetic field causes step broadening which is small for the weak field H due to both the small parameter $nH/H^* \sim d/2r_c \ll 1$ and the small numerical coefficient $\tfrac{1}{16}\pi^2(\tfrac{1}{6} + 16/\pi^4) \simeq 0.2$. Calculations for small n show that the corresponding coefficient is negative only for $n = 1$ (and is also small). The sharpening of the $n = 1$ step results from the magnetic field induced enhancement of the effective mass which occurs for the ground term only.

We now study the $G(z)$ function for the range of strong fields, when $2r_c - d(0) \ll 2r_c$. We assume, however, that the condition $E_F \gg \hbar\omega_c$ remains valid (the field is non-quantising). In the case under consideration large values of n are of interest. Therefore, the terms are determined by means of the semiclassical quantisation procedure. Since of particular interest is the switching process for a mode, the turning points for transverse motion should coincide with the boundaries of the channel, $y = \pm d(x)/2$. Using the Bohr–Zommerfeld rule we obtain

$$(1 - 2y_0/d(x))\{\varepsilon - [1 - 2y_0/d(x)]^2\}^{1/2} + [1 + 2y_0/d(x)]/\{\varepsilon - [1 + 2y_0/d(x)]^2\}^{1/2}$$
$$+ \varepsilon\{\sin^{-1}[1 - 2y_0/d(x)/\varepsilon^{1/2}] + \sin^{-1}[1 + 2y_0/d(x)/\varepsilon^{1/2}]\} = 4\varepsilon_0^{1/2} \tag{22}$$

where $y_0 = -cp/|e|H$; $\varepsilon = 8E_n(p, d(x))/m\omega_c^2 d^2(x)$, $\varepsilon_0 = 8E_n^{(0)}/m\omega_c^2 d^2(x)$ and $E_n^{(0)} = (\pi\hbar n)^2/2md^2(x)$. If the quantity $d(0)$ approaches the $2r_c$ value, then $m\omega_c^2 d^2(0)/8 \simeq E_F$ and $\varepsilon - 1 \ll 1$. Therefore, from (22) we obtain

$$E_n(0, d(0)) = \hbar\omega_c n + \tfrac{4}{3}(E_F/\pi)[(2r_c/d)^2 - 1]^{3/2} \tag{23a}$$

$$E_{pp}^n = (2/\pi m)[(2r_c/d)^2 - 1]^{-1/2}. \tag{23b}$$

Now let us estimate plateau widths. From the equation $E_n(0, d(0)) = E_F$ it follows that

$$1 - \pi z_n/2r_c k_F = \tfrac{1}{2}(3\pi\hbar\omega_c/4E_F)^{2/3}(E_F/\hbar\omega_c - n)^{2/3}. \tag{24}$$

We now introduce a parameter $n_0 = E_F/\hbar\omega_c \equiv H^*/H$. The integer part of this parameter determines the full number of steps in the $G(z)$ function (at a fixed field value H). At $1 \ll n_0 - n \ll n_0$ equation (24) leads to

$$z_{n+1} - z_n = (4/3\pi)^{1/3}[n_0/(n_0 - n)]^{1/3}. \tag{25}$$

The plateau widths increase parametrically when n tends to n_0.

The width of steps can be found by formulae (16), (23), (24):

$$\delta z_n(H)/\delta z_n(H = 0) = (2/\pi^{1/2})(4/3\pi)^{1/3}[n_0/(n_0 - n)]^{1/3} \qquad n_0 = H^*/H. \tag{26}$$

Thus, the step width also increases parametrically as n tends to n_0. The values δz_n and $z_{n+1} - z_n$ depend upon n according to the same law. The step–plateau width ratio remains small: $\delta z_n/(z_{n+1} - z_n) \sim (r_c/R)^{1/2} \ll 1$, i.e. the step-like form of the function $G(z)$ is maintained in the strong magnetic field limit. It is important that in this limit the relation

$$\delta z_n/(z_{n+1} - z_n) = (1/\pi^2)(d/2R)^{1/2} \tag{27}$$

coincides with the corresponding formula in the case of $H = 0$ with the same numerical factor.

The theory developed above allows explaining the observed [2, 5] effect of the magnetic field on the conductance of a lateral microconstraint.

It has been shown experimentally [2], that for the in-plane magnetic field orientation spin splitting first occurs for steps with large step numbers n. This is in agreement with the conclusions drawn from equation (8). It has been demonstrated, that the value of spin splitting exceeds the step width for steps $n > n_S$ (see (8)). Here we estimate the value of n_S for the conditions of the experiment [2]: $H = 13.6\ T$, $\lambda_F = 400$ Å. To determine the value of R essential for calculations one can use the results of [8] for the electrostatics of the microconstraint. This value is approximately equal to the distance from the lithographic gap edge to the nearest boundary of the electrostatically confined electronic channel. Neglecting the thickness $h \simeq 800$ Å of the layer between the planes of the gate and 2DEG, we obtain $R = [D - d(0)]/2$, where $D = 0.3$–0.7 is the width of the lithographic gap. The channel width $d(0) = n\lambda_F/2$ in the experiment was equal to $d(0) \simeq 0.2\ \mu$m for $n = 10$. Therefore, neglecting of h is reasonable for the largest value of D. For $D \simeq 0.3\ \mu$m the quantity h should be taken into account: $R \simeq \{[D - d(0)]^2/4 + h^2\}^{1/2}$ [8]. Calculation by equation (8) yields $n_S = 3$–10 which is in reasonable agreement with the experimental value $n_S = 10$.

The orbital effects were studied in the experiment [5], where the magnetic field was applied perpendicular to the plane of the heterostructure. It should be noted that

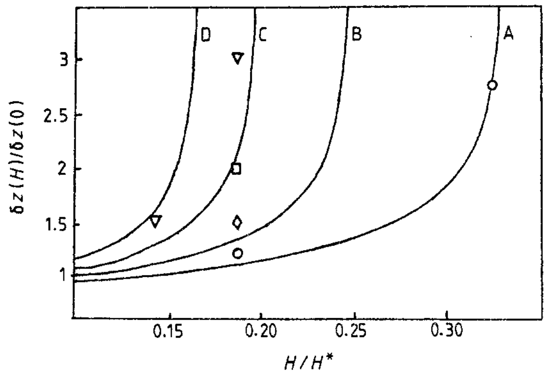

Figure 1. Dependence of normalised step widths on magnetic field H. Curves A, B, C and D correspond to asymptotic dependence (equation (26)) for $n = 3, 4, 5, 6$, respectively. The experimental values were obtained by treating the curves from reference [5] for $H = 0.7$; 1.0; 1.8 T. ○, $n = 3$; ◇, $n = 4$; □, $n = 5$; ▽, $n = 6$.

considerable shifts of step positions were already observed in the region of n, H, where step broadening was negligible. This agrees with the conclusions made with regard to the $G(z)$ dependence in a weak magnetic field (equation (21)). Unfortunately, reference [5] does not contain sufficiently accurate information for comparison with our equations (21) and (26) for step widths. We have attempted, however, to treat the experimental data of [5] for relatively large numbers $n = 3$–6, where the broadening of steps was considerable. The results of the comparison with equation (26) are shown in figure 1. We have chosen the value of $H^* = 5.6$ T, which corresponds to $E_F = 9.4$ meV.

In conclusion, it should be emphasised that in [3] the conductance quantisation has been explained by adiabatic propagation of an electron wave through the constraint. In this work we have generalised the above approach to include the case of an applied external magnetic field. The theory developed is able to account for the experimental results.

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