Temperature effects on edge-state properties in the integer quantum Hall regime

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Abstract

The edge and bulk structure of Landau levels (LLs) in a wide channel at the $\nu = 1$ quantum Hall regime is calculated for not-too-low temperatures, $\hbar \omega_c \gg k_B T \gg \hbar v_g/2 \ell_0$, where $v_g$ is the group velocity of the edge states and $\ell_0 = \sqrt{\hbar c/e|B|}$ is the magnetic length. Edge-states correlations essentially modify the spatial behavior of the lowest spin-up LL, which is occupied, compared to the lowest spin-down LL, which is empty. The influence of many-body interactions on the spatially inhomogeneous spin-splitting between the two lowest LLs is studied within the generalized local density approximation.
Temperature effects on the enhanced spin-splitting, the position of the Fermi level within the exchange enhanced gap and the renormalization of edge-states group velocity by edge states screening are considered. It is shown that the maximum activation energy $G$ in the bulk of the channel is determined by the gap between the Fermi level and the bottom of the spin-down LL, because the gap between the Fermi level and the spin-up LL is much larger. For the maximum value of $G$, it is shown that the renormalized group velocity $v_g \propto T$ for $T \to 0$ and, in particular, the condition of not-too-low $T$ can be satisfied for $4.2 \gtrsim T \gtrsim 0.3$ K. In other words, the regime of not-too-low temperatures regime can be achieved even for rather low $T$. 

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A number of works have been devoted to study the influence of Coulomb interactions on the edge states of a two-dimensional electron system (2DES) in the presence of a strong magnetic field. Quite recently, the spatial behavior of the Landau levels (LL’s) for a wide channel at the filling factor \( \nu = 1 \) was studied in the limit of very low-temperatures, \( T \ll \hbar v_g(0)/k_B\ell_0 \), using the generalized local density approximation (GLDA) in which exchange and correlation effects, especially due to the edge states, are included. \( v_g(0) \) is the group velocity of the edge states and the magnetic length \( \ell_0 = \sqrt{\hbar/m^*\omega_c} \) with \( \omega_c = |e|B/m^*c \). Many-body effects were calculated in the screened Hartree-Fock approximation. It was shown that the maximum activation energy in the bulk of the channel corresponds to a highly asymmetric position of the Fermi level within the gap between spin-down and spin-up LL’s. The renormalized edge-state group velocity \( v_g(0) \) independent on \( T \) exhibits a strong decrease due to correlation effects.

In this work, we extend the previous approach of Ref. [7] for not-too-low temperatures, where \( \hbar\omega_c \gg k_B T \gg \hbar v_g/2\ell_0 \) and the characteristic length \( \ell_T = \ell_0^2 k_B T/\hbar v_g \gg \ell_0/2 \) becomes relevant. In particular, we study here the temperature effects on the enhanced spin-splitting for the 2DES in a wide channel within the \( \nu = 1 \) quantum Hall regime, the position of the Fermi level within the pertinent exchange enhanced gap and the renormalized edge-states group velocity \( v_g \), which can be strongly dependent on \( T \).

We consider a wide symmetric channel along the \( x \) axis in strict two dimensions with a strong uniform magnetic field \( B \) in the \( z \) direction. In the absence of exchange and correlation effects, we take the confining potential \( V_w(y) = 0 \) in the inner part of the channel and \( V_w(y) = m^*\Omega^2 (y-y_r)^2/2 \) at the right edge, \( y \geq y_r \), and \( \Omega \) is the confining frequency. We assume that \( V_w(y) \) is smooth on the scale of the magnetic length \( \ell_0 \); i.e., \( \Omega \ll \omega_c \). Extending the GLDA for \( T = 0 \) of Ref. [7] on not-too-low \( T \), in particular, by using some results of Ref. [8], it follows that the energy spectrum \( E_{0,k_x,1} \) of lowest spin-up LL \( (n = 0, \sigma = 1) \) of the interacting 2DES for \( r_0 = e^2/(\varepsilon\ell_0\hbar\omega_c) \lesssim 1 \) can be obtained from the solution of the single-particle Schrödinger equation with the Hamiltonian \( \mathcal{H} = \mathcal{H}^0 + V_{xc}(y) \), where a self-consistent exchange-correlation potential, smooth on \( \ell_0 \) scale, is given by
\[ V_{xc}(y) = E_{0,y}/\ell_{0,1} - \left( \frac{\hbar \omega_c}{2} - \frac{|g_0| \mu_B B}{2} + V_w(y) \right). \] (1)

Here \( E_{0,k_x,1} = \varepsilon_{0,k_x,1} + \varepsilon_{0,k_x,1}^c \), where \( \varepsilon_{n,k_x,\sigma} \) are the eigenvalues of the one-electron Hamiltonian \( \mathcal{H}^0 = [(\hat{p}_x + eBy/c)^2 + \hat{p}_y^2]/2m^* + V_w(y) + g_0 \mu_B \hat{S}_z B/2 \) and \( \varepsilon_{0,k_x,1}^c \) is the exchange and correlation contributions in the screened Hartree-Fock approximation; \( g_0 \) the bare g-factor, \( \mu_B \) the Bohr magneton and \( \hat{S}_z \) is the z component of the spin operator with eigenvalues \( \sigma = 1 \) and \( \sigma = -1 \) for spin up and down; \( \varepsilon \) is the background dielectric constant. Further, Eq. (1) is essentially different from Eq. (17) of Ref. [7] as for \( E_{0,k_x,1} \) now we have a very different expression, \( k_x \geq 0 \) and \( \tilde{v}_g = (\pi \hbar \varepsilon / e^2)v_g \), as

\[
E_{0,k_x,1} = \frac{\hbar \omega_c}{2} - \frac{|g_0| \mu_B B}{2} + \frac{m^* \Omega^2 \ell_0^2}{2}(k_x - k_r)^2 \Theta(k_x - k_r) - \sqrt{\frac{\pi}{2}} \varepsilon \ell_0 k_{\text{bulk}}(r_0) f_{0,k_x,1} + \frac{2e^2}{\pi \varepsilon \tilde{v}_g} f_{0,k_x,1} \times \int_0^\infty dq_x \exp(-q_x^2 \ell_0^2/2) F_1(q_x, k_x - k^{(1)}_{r_0}) F_2(q_x, k_x - k^{(2)}_{r_0}) \left[ 1 + \frac{1}{\tilde{v}_g} \exp(-q_x^2 \ell_0^2/2) r_{00}^s(q_x) \right]^{-1}, \] (2)

where \( f_{0,k_x,1} = 1 / \left[ 1 + \exp(E_{0,k_x,1} - E_F)/k_B T \right] \approx \{1 - \tanh(\ell_0^2 (k_x - k^{(1)}_{r_0}) / 2 \ell_T)\}/2, \) with \( E_F = E_{0,k^{(1)}_{r_0},1} \) being the Fermi energy renormalized both by exchange and correlations and \( \Theta(k_x - k_r) \) is the step function. The edge of the \((n = 0, \sigma = 1)\) LL is denoted by \( y^{(1)}_{r_0} = \ell_0^2 k^{(1)}_{r_0} \), with \( k^{(1)}_{r_0} = k_r + k^0_{x,1} \) and \( k_r = y_r / \ell_0^2 \) and it is assumed that the Fermi wave vectors \( k^{(1)}_{r_0} \) and \( k^0_{x,1} \) are not changed by exchange correlation effects, even though the Fermi energy in the Hartree approximation \( E^H_F \) is essentially different from \( E_F \). The fourth term of Eq. (2) comes from the exchange interaction, where the additional factor \( k_{\text{bulk}}(r_0) \) < 1 takes into account the decreasing of the many-body contribution due to the weak “bulk” screening of the fully occupied LL caused by inter-LL virtual transitions. For assumed conditions \( k_{\text{bulk}}(r_0) \) is well approximated by a factor calculated exactly in Ref. [6], for the bulk of the channel. For instance, we find \( k_{\text{bulk}}(r_0) \) are 0.79, 0.74, 0.70, 0.66 and 0.63 for \( r_0 = 0.6, 0.8, 1.0, 1.2, \) and 1.4, respectively. For \( r_0 < 1 \), \( k_{\text{bulk}}(r_0) \) tends to 1 and the well-known exchange contribution is recovered. In the last term of Eq. (2), the function \( r_{00}^s(q_x) \), dependent on \( \ell_T \), is presented in Ref. [5] and the functions \( F_1 \) and \( F_2 \) are given by

\[
F_1(q_x, k_x - k^{(1)}_{r_0}) = \int_0^\infty dq_y \frac{\exp(-q_y^2 \ell_0^2/4) \cos(q_y (k_x - k^{(1)}_{r_0}) \ell_T^2)}{\sqrt{q_x^2 + q_y^2} \left[ 1 + q_y^2 \ell_T^2 \right]}, \] (3)
\[
F_2(q_x, k_x - k_{v,0}^{(1)}) = \int_0^\infty dq_y \exp(-q_y^2 \ell_0^2 / 4) \sqrt{q_x^2 + q_y^2} \cos\{q_y (k_x - k_{v,0}^{(1)}) \ell_0^2 \} F(q_y),
\]

where \( F(q_y) = \pi |q_y| \ell_T / \sinh(\pi |q_y| \ell_T). \) The group velocity of the edge states \( v_g = (\partial E_{0,k_x,1}/\hbar \partial k_x)_{k_x = k_{v,0}^{(1)}} \) renormalized by exchange and correlations, is determined by a positive solution of the equation

\[
\tilde{v}_g = \tilde{v}_g^H + \frac{\pi \ell_0}{4 \ell_T} \left\{ \sqrt{\frac{\pi}{2}} k_{\text{bulk}}(r_0) - \frac{2\ell_0}{\pi} \int_0^\infty dq_x \exp(-q_x^2 \ell_0^2) \frac{F_1(q_x, 0) F_2(q_x, 0) \ell_0}{\tilde{v}_g + \exp(-q_x^2 \ell_0^2) r_{00}^s(q_x)} \right\},
\]

where \( \tilde{v}_g^H = (\pi \hbar \varepsilon / e^2) v_g^H \) and \( v_g^H = v_{g,0}^{(1)} \), with \( v_{g,0}^{(1)} = c \mathcal{E}_0^{(1)} / \hbar \) being the group velocity in the Hartree approximation. Here \( \mathcal{E}_0^{(1)} = \Omega \sqrt{2m^* \Delta_{F_0}^{(1)}/|e|} \) is the electric field associated with the confining potential \( V_w(y) \) at \( y = (1)_{0}^{(1)} \) and \( \Delta_{F_0}^{(1)} = E_F^H - \hbar \omega_c/2 - g_0 \mu_B B / 2 \). Equation (5) was evaluated using \( [\partial F_{1,2}(q_x, k_x - k_{v,0}^{(1)}/\partial k_x)_{k_x = k_{v,0}^{(1)}} = 0 \) and the relation \( [-\partial f_{0,k_x,1}/\partial k_x]_{k_x = k_{v,0}^{(1)}} = \ell_0^2 / 4 \ell_T \). Note that since \( \ell_T = \ell_0^2 k_B T / \hbar v_g \), we have that \( \ell_T \propto v_g^{-1} \) in Eqs. (2)-(5). Equations (4) and (3) provide the self-consistent scheme to calculate the LL spectrum in the GLDA for not too-low temperatures. We observe that for \( \tilde{v}_g^H \to 0 \), the proper solution \( \tilde{v}_g \) tends to a finite value. This result is quite different from the case of very low temperatures in which \( v_g \propto \sqrt{v_g^H} \), for \( v_g^H \to 0 \). If correlations due to edge-states screening are neglected (excluding the term with integral) in Eq. (3), then it follows that for many experimentally realistic conditions Eq. (3) does not have any positive solution, i.e., the condition \( \sqrt{\pi / 32} r_0 k_{\text{bulk}}(r_0)[\hbar \omega_c / k_B T] < 1 \) is not satisfied.

The positive gap between the bottom of the upper spin-split LL and the Fermi level of the interacting 2DES, \( G(v_g^H) = E_{0,0,-1} - E_{0,k_{v,0}^{(1)},1} \), is then written as

\[
G = |g_0| \mu_B B - m^* \omega_c^2 (v_g^H)^2 + \frac{1}{2} \sqrt{\frac{\pi}{2}} \frac{e^2}{\varepsilon \ell_0} k_{\text{bulk}}(r_0) - \frac{e^2}{\pi \varepsilon \tilde{v}_g} \int_0^\infty dq_x \exp(-q_x^2 \ell_0^2) \\
x F_1(q_x, 0) F_2(q_x, 0) \left[ 1 + \frac{1}{\tilde{v}_g} \exp(-q_x^2 \ell_0^2/2) r_{00}^s(q_x) \right]^{-1},
\]

where \( \tilde{v}_g \equiv \tilde{v}_g(v_g^H) \geq 0 \) is the function of \( v_g^H \) obtained from Eq. (3). In addition, in Eq. (3) it was used for the term \( \propto (k_{v,0}^{(1)} - k_r)^2 \) that \( k_{e,0}^{(1)} = (m^* \omega_c^2 / \hbar \Omega^2) v_g^H \). In the bulk of
In the absence of many-body interactions, the maximum value of the dimensionless activation gap $G_a(v_g^H) = G/(\hbar\omega_c/2) \approx 0.015$. We use the parameters of GaAs based samples, in particular $\varepsilon \approx 12.5$, $g_0 = -0.44$ and $m^* = 0.067m_0$. Then the activation gap is enhanced when $G_a > 0.015$. The asymmetry of the Fermi level position within the Fermi gap in the bulk of the channel can be characterized by another dimensionless function $\delta G(v_g^H) = (\bar{G}_{-1,1} - G_a)/G_a$, where $\bar{G}_{-1,1} = G_{-1,1}/(\hbar\omega_c/2)$. Notice that, when $v_g^H \to 0$, $E_H^F$ tends, from the upper side, to the bottom of the $(n = 0, \sigma = 1)$ LL, in the absence of many-body interactions.

In Fig. 1 $\tilde{v}_g$ is depicted as a function of $\tilde{v}_g^H$ using Eq. (5). The solid, dashed and dotted curves correspond to $\hbar\omega_c/k_BT = 10, 15$ and 74.5 or $T \approx 31.3, 20.8$ and 4.2 K respectively for $B = 15.7$ T and electron density $n_s = 3.8 \times 10^{11}$ cm$^{-2}$, $r_0 = 0.66$, $k_{\text{bulk}}(r_0) = 0.77$. From Fig. 1, it is seen that for $\tilde{v}_g^H \to 0$ the renormalized group velocity $\tilde{v}_g$ tend to a finite value in contrast with the regime of very low temperatures. We point out that for these parameters there is no positive solution $v_g$ if edge-state correlations are neglected. Notice that for $\tilde{v}_g^H = 0$, depending on $v_g$, the ratio $\ell_T/\ell_0 = 7.3, 3.6$ and 2.0 for the solid, dashed and dotted curves in Fig. 1. So the assumed condition $2\ell_T/\ell_0 \gg 1$ is well satisfied in Fig. 1.

In Fig. 2, using Eqs. (5) and Eq. (6), we plot $G_a$ and $(\delta G/10)$ as a function of $\tilde{v}_g^H$. In the solid and dashed curves we used the same parameters as in Fig. 1. $G_a$ is represented by the curves with a negative slope while the curves with positive slope represent $(\delta G/10)$. Here we use the value of the confining frequency $\Omega = 8.2 \times 10^{11}$ s$^{-1}$, which is slight larger than $\Omega = 7.8 \times 10^{11}$ s$^{-1}$ obtained in Ref. [9]. In this case $\omega_c/\Omega = 50$. We observe that the maximum of $G_a$ and the corresponding minimum of $\delta G$ are kept for $\tilde{v}_g^H \to 0$. In particular, in the solid curves, these values are $G_a \approx 0.43$ and $\delta G \approx 2.0$. In the dashed curves, the values are $G_a \approx 0.30$ and $\delta G \approx 3.4$ as $\tilde{v}_g^H \to 0$. Notice that for $\tilde{v}_g^H \to 0$, and $\tilde{v}_g$ given by the dotted curve in Fig. 1, we have $G_a \approx 8.7 \times 10^{-2}$ and $\delta G \approx 14.1$, which gives $G \approx 13.5$ K,
i.e. about three times greater than, the liquid helium temperature, \( T = 4.2 \) K. So the Fermi level is substantially closer to the bottom of the upper spin-split LL then to the bottom of the lowest spin-split LL. Such an asymmetry increases as \( T \) decreases. Furthermore, one can see that \( G_a \) drops fast by increasing \( \tilde{v}_g^H \). Note, however, that for sufficiently small \( G_a \), such that \( G = (\hbar \omega_c/2)G_a \leq k_BT \), our formulas, of course, are not valid.

In Fig. 3 \( \tilde{v}_g \) is shown as a function of \( T \) by solid curves for \( \tilde{v}_g^H = 0 \), calculated from Eq. (5). The dashed curves represent \( \ell_T/\ell_0 \) pertinent to the solid curves. The long curves correspond to the parameters \( n_s = 3.8 \times 10^{11} \) cm\(^{-2} \), \( B = 15.7 \) T, \( r_0 = 0.66 \), \( k_{\text{bulk}}(r_0) = 0.77 \). The short curves correspond to \( n_s = 1.26 \times 10^{11} \) cm\(^{-2} \), \( B = 5.2 \) T, \( r_0 = 1.146 \), \( k_{\text{bulk}}(r_0) = 0.67 \). In Fig. 3 the curves are plotted for the temperature range satisfying \( \hbar \omega_c/k_BT \geq 10 \). As \( T \) drops to below 4.2 K (as well for \( T \ll 1 \) K), the upper and lower dashed curves tend to \( 2\ell_T/\ell_0 \approx 5.2 \) and 3.8, respectively. From this result it follows that the regime of not-too-low temperatures, \( k_BT \gg \hbar v_g/2\ell_0 \), can hold even for \( T \lesssim 1 \)K.

In Fig. 4 we show the energy spectra as a function of \( Y = \ell_T^2(k_x - k_r^{(1)})/\ell_T \) for \( \tilde{v}_g^H = 0 \) and the parameters of solid curves in Figs. 1 and 2. In particular, we use \( \hbar \omega_c/k_BT = 10 \), \( B = 15.7 \) T, \( \omega_c/\Omega = 50 \) and \( \tilde{v}_g = 0.065 \); here \( k_r^{(1)} = k_r \) as \( \tilde{v}_g^H = 0 \). Using Eq. (2), \( E_{0,k_x,-1} = \varepsilon_{0,k_x,-1} \) and the dashed horizontal line is the Fermi level of the interacting 2DES.

In conclusion, we have shown that edge-state correlations modify drastically the \((n = 0, \sigma = 1)\) LL spectrum of the 2DES in the \( \nu = 1 \) quantum Hall regime at not-too-low temperatures, \( \hbar \omega_c \gg k_BT \gg \hbar v_g/2\ell_0 \), in a wide region \( (\gg \ell_T) \) nearby the channel edge. The spectrum is obtained in a well justified self-consistent GLDA. We have shown that the position of the Fermi level \( E_F \) is highly asymmetric within the gap defined by the \((n = 0, \sigma = 1)\) and \((n = 0, \sigma = -1)\) LLs in the bulk of the channel due to such correlations. Hence, the activation gap \( \Delta_{ac} = 2G \) (obtained from the resistivity as \( \rho_{xx} \propto \exp(-\Delta_{ac}/2k_BT) \)) is much smaller than the Fermi gap. In particular, we have obtained the edge-state group velocity \( v_g \), renormalized by exchange and correlations effects, and an analytical expression for \( G \), the gap between the bottom of the upper spin-split LL and \( E_F \). We have obtained
that there is a maximum of $G$ for $\tilde{v}_g^H = 0$, which essentially depends on $T$. Moreover, our self-consistent study shows that for $\tilde{v}_g^H = 0$ the regime of not-too-low temperatures can be achieved for very small temperatures, e.g., $4.2 \gtrsim T \gtrsim 0.3$ K, as for $T \lesssim 1$ it follows that $v_g \propto T$ and $2\ell_T/\ell_0$ becomes temperature independent while it is still large, $2\ell_T/\ell_0 \gg 1$.

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FIGURE CAPTIONS

Fig. 1. Renormalized group velocity $\tilde{v}_g$ as a function of the group velocity within the Hartree approximation $\tilde{v}^H_g$, in units of $e^2/\pi \hbar \epsilon$. The solid, dashed and dotted curves correspond to $\hbar \omega_c/k_B T = 10, 15$ and 74.5, or $T \approx 31.3, 20.8$ and 4.2 K respectively for $B = 15.7$ T and electron density $n_s = 3.8 \times 10^{11}$ cm$^{-2}$ (the $\nu = 1$ quantum Hall regime is considered).

Fig. 2. Many-body enhancement of the activation gap $G_a = G/(\hbar \omega_c/2)$ and fractional difference $\delta G = (G_{-1,1} - G)/G$, where $G_{-1,1} = E_{0,0,1} - E_{0,0,-1}$, as a function of $\tilde{v}^H_g$. $\delta G$ displays the asymmetry of the Fermi level position within the Fermi gap in the bulk of the channel. The solid and dashed curves correspond to the solid and dashed curves in Fig. 1. $G_a$ is represented by curves with negative slope and those with positive slope corresponds to $\delta G$; for $\omega_c/\Omega = 50$. 
Fig. 3. Renormalized group velocity $\tilde{v}_g$ as a function of temperature for $\tilde{v}_g^H = 0$ is represented by solid curves; $\ell_T/\ell_0$ is shown by dashed curves. The long curves correspond to the sample parameters $n_s = 3.8 \times 10^{11} \text{ cm}^{-2}$; $B = 15.7 \text{ T}$, $r_0 = 0.66$, $k_{\text{bulk}}(r_0) = 0.77$, while the short curves correspond to $n_s = 1.26 \times 10^{11} \text{ cm}^{-2}$; $B = 5.2 \text{ T}$, $r_0 = 1.15$, $k_{\text{bulk}}(r_0) = 0.67$. The curves are plotted only for $T \leq \hbar \omega_c/10k_B$. For $T \lesssim 1 \text{ K}$, the upper and lower dashed curves give $2\ell_T/\ell_0 \approx 5.2$ and 3.8, respectively.

Fig. 4. Energy spectra as a function of $Y = \ell_0^2(k_x - k_x^{(1)})/\ell_T$ for $\tilde{v}_g^H = 0$ and same parameters used for the solid curves in Figs. 1 and 2; $\hbar \omega_c/k_BT = 10$, $B = 15.7 \text{ T}$, $\omega_c/\Omega = 50$ and $\tilde{v}_g = 0.065$. Here $k_x^{(1)} = k_x$ as $\tilde{v}_g^H = 0$. The lower solid curve represents $E_{0,k_x,1}$, while the upper solid curve is $E_{0,k_x,-1} = \epsilon_{0,k_x,-1}$. The dashed horizontal line is the Fermi level of the interacting 2DES.
Fig. 1
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