Probing two-subband systems in a quantizing magnetic field with non-equilibrium phonons

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Abstract. We propose to use phonon absorption spectroscopy to study many-body gaps and phases of two-subband heterostructures in the quantum Hall regime. Implications of the spin-orbit interaction for phonon absorption in this system are considered.

Phonon spectroscopy is a powerful experimental technique that has been successfully used in the study of the properties of a two-dimensional (2D) electron gas in the quantum Hall regime [1]. It allows probing of the states with non-zero momentum, which are not accessible by either cyclotron resonance or photoluminescence methods. We propose to use acoustic-phonon spectroscopy to study many-body effects in a two-subband quasi-2D system, formed at a single semiconductor heterojunction and subjected to a quantizing magnetic field. We show that the formation of many-body gaps and the existence of different spin phases in this system result in a striking difference in the magnetic field dependence of the phonon absorption rate at different filling factors.

Acoustic-phonon scattering of 2D electrons in a quantizing magnetic field is strongly suppressed due to the dispersionless structure of the Landau levels (LLs). In a two-subband system, however, a sharp enhancement of the electron-phonon interaction occurs when the two LLs corresponding to different size-quantization subbands (we consider the second LL of the first subband and the first LL of the second subband only) approach one another, and the separation between these energy levels, $\Delta = \Delta_{12} - \hbar \omega_c$, becomes close to $\hbar s / l$. Here $\hbar \omega_c$ is the cyclotron energy, $\Delta_{12}$ is the intersubband splitting, $s$ is the velocity of sound and $l$ is the magnetic length. This electron-phonon interaction enhancement should manifest itself in the magnetic field dependence of the phonon absorption rate and dissipative conductivity. At filling factor slightly greater than two, when the electron-electron interaction between the particles belonging to either of the intersecting LLs can be neglected, a double-peak structure in the magnetic field dependence of the phonon-mediated transport properties of the two-subband system has been predicted [2]. With increasing electron density, the repulsion between the levels due to electron-electron exchange interactions opens a gap in the interlevel excitation spectra. At filling factor $\nu = 3$ this gap exists at all values of the bare level splitting $\Delta$ and can well exceed $\hbar s / l$. As a result, the double-peak structure transforms into a single-peak one (see Figure 1). By far the most interesting case is when the electron filling factor $\nu = 4$. When the cyclotron energy is close to the intersubband splitting the system can be mapped onto a four-level electron system with an effective filling factor $\nu^* = 2$. The ground state is either a ferromagnetic state or a spin-singlet state, depending on the values of the interlevel splitting and Zeeman energy $\Delta_Z$. The electron-electron interaction renormalizes the phase boundaries, which results in the existence of a ferromagnetic phase even for $\Delta_Z = 0$. The excitation spectrum has a gap for any value of the interlevel splitting $\Delta$ and Zeeman
splitting $\Delta z$. This results in strong suppression of the electron-phonon interaction. The rate of phonon absorption by the considered quasi-2D electron system at $\nu = 4$ again has a double-peak structure as a function of level splitting and a steplike structure as a function of Zeeman splitting (see Figure 2).

In the absence of the spin-orbit interaction, the transitions between levels with different spin indices are forbidden because the electron-phonon interaction conserves the electron spin. The spin-orbit interaction allows such transitions. Matrix elements of the spin-orbit interaction Hamiltonian $H_{SO}$ between levels differing by subband number, LL number, and spin projection on the normal to the structure plane, are given by following expressions [3]:

$$\langle \eta, n, k_x, -1/2 | H_{SO} | \mu, n - 1, k_x, +1/2 \rangle = iC \frac{1}{\hbar \Omega} \sqrt{2n},$$

$$\langle \eta, n, k_x, +1/2 | H_{SO} | \mu, n - 1, k_x, -1/2 \rangle = A \frac{1}{\hbar \Omega} \sqrt{2n},$$

$$\langle \eta, n - 1, k_x, -1/2 | H_{SO} | \mu, n, k_x, +1/2 \rangle = A \frac{1}{\hbar \Omega} \sqrt{2n},$$

$$\langle \eta, n - 1, k_x, +1/2 | H_{SO} | \mu, n, k_x, -1/2 \rangle = -iC \frac{1}{\hbar \Omega} \sqrt{2n},$$

(1)

where $H_{SO} = \left( \alpha \hbar^3 (2m^3 E_g)^{-1/2} \mathbf{k} + \beta [\mathbf{k} \times \nabla] \right) \cdot \mathbf{S}$, $\mathbf{k} = [k_x (k_x^2 - k_y^2), k_y (k_y^2 - k_z^2), k_z (k_z^2 - k_x^2)]$, $\eta$ and $\mu$ numerate subbands, $n$ is the LL number, and $\alpha$ and $\beta$ are phenomenological spin-orbit constants of the semiconductor [4]. The expressions for $A$ and $C$ are given by

$$A = \alpha \hbar^3 (2m^3 E_g)^{-1/2} \int_{-\infty}^{+\infty} \Psi_\mu(z) \frac{d^2 \Psi_\eta(z)}{dz^2} dz,$$

$$C = \beta \int_{-\infty}^{+\infty} \Psi_\mu(z) \Psi_\eta(z) \frac{d V_{QW}}{dz} dz,$$

(2)

where $V_{QW}$ is the quantum-well confinement potential. These matrix elements cause an admixture of states with opposite spin direction to both the initial and final states for the phonon absorption. As a result, the phonon absorption rate at $\nu = 4$ never becomes zero; in the region between peaks, it is of the order of $(M/\Delta g)^2$ times its value at the peak of absorption, where $M$ is the corresponding matrix element of the spin-orbit interaction (Eq.(1)), and $\Delta g$ is the many-body energy gap.

In the absence of many-body effects, the inter-subband spin-orbit coupling results in anticrossing of levels with different spin, subband and LL indices. The resulting energy gaps can be measured by high-resolution spectroscopic techniques (optical or spin/cyclotron resonance spectroscopy), as has been suggested for one-subband systems in tilted magnetic fields [5].

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References

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Figure 1. The phonon absorption rate per phonon, $\omega_0 = \omega_{\text{abs}}/n(q_0)$, as a function of bare level separation, $\Delta = \Delta_{12} - \hbar \omega_c$, at filling factor $\nu = 3$ (solid and dotted lines) and for the case of non-interacting electrons (dashed line). The results for $\nu = 3$ are shown for two different values of the Fang-Howard parameter (inverse quantum well width) $b = 1$ and $b = 1.5$ in units of inverse magnetic length $l^{-1}$. The absorption rate $\omega_0$ is in units of $10^{10} \text{s}^{-1}$; $\Delta$ is in Coulomb units ($\varepsilon_C = e^2/\kappa l$).

Figure 2. The phonon absorption rate $\omega_0$ at filling factor $\nu = 4$ as a function of interlevel separation $\Delta$ for the case of zero Zeeman splitting, $\Delta_Z = 0$, for three values of parameter $b$. The data for $b = 1.0$ and 1.5 are multiplied by 10. The absorption rate $\omega_0$ is in units of $10^{10} \text{s}^{-1}$; $\Delta$ is in units of $\varepsilon_C$. 