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Spin-flip excitations in the fractional-quantum-Hall-effect regime studied by polarized photoluminescence of charged excitons

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We report on a low-temperature magneto-optical study of the spin-singlet charged exciton (X_s^-) transitions in GaAs/$Al_{0.1}Ga_{0.9}As$ modulation-doped multiple quantum wells with an optically tuned two-dimensional electron gas (2DEG) density corresponding to a filling factor in the range $0<\nu\leqslant 1$. We find strong evidence for the effect of 2DEG magnetic correlations on the X_s^- spin-resolved transitions. The energy difference between the σ^-- and σ^+-polarized X_s^- photoluminescence (PL) intensity maxima is found to oscillate with varying ν. The numerical derivative of these oscillations with respect to ν shows minima at odd denominator rational fractions and maxima at even denominator fractions, similar to the longitudinal resistivity measured in the fractional quantum Hall effect. At $\nu=1$, the observed σ^--polarized low-energy PL tail is interpreted (by a line-shape analysis) to be due to a finite-**k** 2DEG spin-wave emission coupled to the optical recombination of the X_s^-. [S0163-1829(97)51144-1]

Magnetooptical spectroscopy of the two-dimensional electron gas (2DEG) formed in high-quality semiconductor structures is used for studying the properties of the correlated electronic system and its interaction with the photoexcited hole. In single heterojunction structures, where the 2DEG is spatially separated from the hole by an internal electric field, the correlated electron states are weakly perturbed by photoexcitation. Then, the electron magnetic correlations which are manifest in magnetotransport as the fractional quantum Hall effect[1,2] (FQHE), are observed in the electron-hole bound state (neutral exciton) radiative recombination.[3–5] The electron FQHE states that are optically detected in quantum well (QW) structures, having a co-planar arrangement of the holes and the 2DEG, are strongly affected by the hole binding.[6] Recent investigations[7–14] of QW's with a low density (n_e) 2DEG have revealed the existence of charged excitons (X^-)—two electron and one (heavy) hole bound complexes. Their appearance was induced either by reducing n_e (Refs. 8–10) or by the application of a sufficiently strong, perpendicular magnetic field[7,14] (B). In both cases the X^- formation results from the electron phase-space depletion and is related to the 2DEG localization.[8,14] In the magnetically quantized 2DEG, the X^- states were shown[14] to form at filling factors $\nu=(hcn_e/eB)<1$ in QW areas occupied by the 2DEG, that result from ionized donor potential fluctuations. In modulation doped QW's with wide spacers, the size of these fluctuations ($\sim$1000 Å) is significantly larger than the characteristic magnetic length $[\lambda=(\hbar c/eB)^{1/2}$ $\sim$100 Å]. This implies that a photocreated electron-hole pair binds one additional electron out of confined electrons that may still maintain long-range correlated FQHE states. In this case, despite the disorder present, the X^- transitions might be affected by the 2DEG collective states that arise from electron-electron interactions.

The present study demonstrates that 2DEG magnetic correlations indeed have a pronounced effect on the X^- spin states. As one can expect, the observed perturbation ($\sim$100 μeV) of the X^- transition energy, due to this effect, is found to be much smaller than the photoexcited hole binding energy ($\simeq$7 meV). The observed spectroscopic features of X^- transitions evidence their coupling to collective spin-flip (SF) excitations of the 2DEG in the extreme magnetic quantum limit. These are shown to be related to the particular 2DEG ground state as determined by the filling factor. We study the photoluminescence (PL) spectra of the spin-singlet (X_s^-) transitions that are split by a magnetic field ($0<B<7$ T) applied perpendicularly to the 2DEG plane and at low temperatures ($1.5<T<3$ K). The PL is resonantly excited just above the $e1$-$hh1$ band gap and is observed parallel to B. It is decomposed into σ^-- and σ^+-polarized spectra and monitored by a double spectrometer (photon energy resolution of $\sim$10 μeV). The studied structure[14,15] is a 300-Å-wide GaAs multiple QW with 3600-Å-thick $Al_{0.1}Ga_{0.9}As$ barriers, symmetrically δ doped in order to obtain a high 2DEG mobility ($\sim 4\times10^6$ cm^2/Vs) and a low 2DEG density of $n_e^0=7.1\times10^{10}$ cm^{-2}. The very same sample exhibits[15] clear FQHE features in low-temperature magnetotransport measurements. n_e is further reduced[14] by an additional above-barrier band-gap photoexcitation down to $\cong 0.8\times10^{10}$ cm^{-2}. The variable 2DEG density allows us to tune the filling factor in the required range $0<\nu\leqslant 1$ under a fixed value of B, which sets the single-particle magnetic energies.

Figure 1 presents the PL spectra of charged and neutral exciton transitions measured for σ^- and σ^+ polarizations.

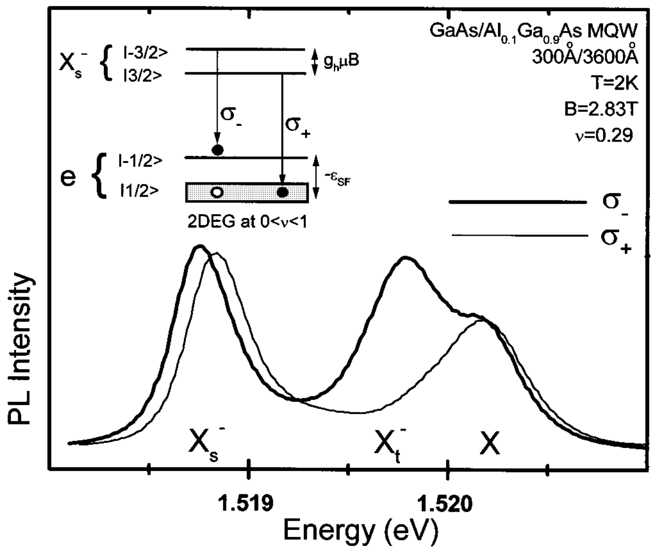

FIG. 1. The σ^-- and σ^+-polarized PL spectra of charged (X^-) and neutral (X) excitons, that are typical of $0<\nu<1$. The inset shows schematically the optical selection rules for the lowest X_s^- transitions.

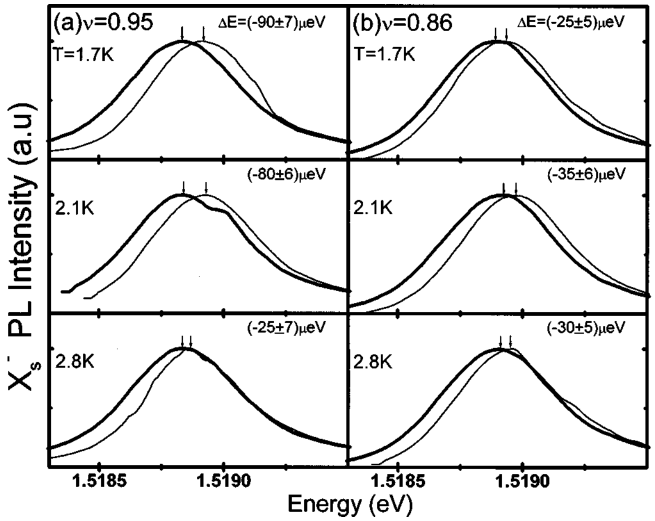

FIG. 2. The σ^-- (thick line) and σ^+-polarized (thin line) X_s^- PL spectra observed for (a) $\nu=0.95$ and (b) $\nu=0.86$, at a fixed $B=2.83$ T and for various temperatures. Each spectrum is normalized to its maximum intensity. The measured values of the energy difference, $\Delta E \equiv E_{\sigma^-}-E_{\sigma^+}$, between the spectral maxima (denoted by arrows) are indicated.

These spectra are typical of filling factors $0<\nu<1$. The present study concerns the lowest X_s^- transitions governed by the selection rules that are shown schematically in the inset of Fig. 1: The initial X_s^- spin-doublet states are split by the hole Zeeman energy ($g_h \mu B$), since the two electrons have antiparallel spins in this state. For σ^+ polarization, the hole recombines with the spin-$\downarrow$ electron so that the spin-$\uparrow$ electron is left below the 2DEG Fermi surface. For σ^- polarization, the spin-$\downarrow$ electron is left along with a spin-hole (a spin-$\uparrow$ electron vacancy). Therefore, the difference between the two circularly polarized recombination final states is due to a single SF process of the 2DEG, while the energy difference between the initial states is assumed to be unaffected by the presence of the rest of the 2DEG. Thus, we use the X_s^- polarized PL transition line shapes in order to study the collective SF excitations of the 2DEG.

The following three figures display various aspects of the energy difference, $\Delta E \equiv E_{\sigma^-}-E_{\sigma^+}$ (measurable for $B>2.8$ T) between the maxima of σ^-- and σ^+-polarized X_s^- PL spectra. Figure 2 shows the PL spectra for two values of the filling factor (a) $\nu=0.95$ and (b) $\nu=0.86$ and at various temperatures. We find a strong ν dependence of ΔE at $T=1.7$ K. However, it weakens rapidly as the temperature is raised to 2.8 K for which ΔE appears to be nearly independent of ν. The value of ΔE obtained at the lowest temperatures is interpreted in terms of a ν-dependent SF excitation energy (ε_{SF}), namely, $\Delta E(\nu,B) \equiv \varepsilon_{\text{SF}}(\nu,B)+g_h(B)\mu B$. There is a strong nonmonotonous B dependence of the hole g factor[16] (a result of the valence $hh1$-$lh1$ subband mixing). Therefore, in order to study a ν dependence of ε_{SF} from the measured ΔE, it is advantageous to fix B and vary n_e. We do it for $B=2.83$ T and $1.0 \leqslant n_e(10^{10}\ \text{cm}^{-2}) \leqslant 7.1$, and plot ΔE as a function of ν in Fig. 3: Its deep oscillations with varying ν yield in the numerical derivative, $\partial \Delta E/\partial \nu$; local minima at rational fractions with an odd denominator and local maxima at those with an even denominator (as indicated by arrows). The depth of the indicated minima (the weak ones at $\frac{1}{7}$, $\frac{2}{7}$ are not labeled) shows almost perfect symmetry around $\nu=\frac{1}{2}$. This correspondence between the local extrema of $\partial \Delta E/\partial \nu$ (observed here at $T \leqslant 2$ K) and the rational ν values strongly indicates the effect of 2DEG energy gaps that are known from numerous FQHE studies (usually observed at $T<1$ K). Therefore, the results presented in Fig. 3 might imply that the SF excitation energy ε_{SF} of an ideal FQHE system ($T=0$ and no disorder) have discontinuities at rational fractions, similar to other quantities like magnetization or chemical potential.[17,18] In the derivative these discontinuities would yield δ extrema with strength that is proportional to corresponding FQHE gaps. Then, at a nonzero T and/or in the presence of a disorder, the steps of $\varepsilon_{\text{SF}}(\nu)$ disappear, while $d\varepsilon_{\text{SF}}(\nu)/d\nu$ is still expected to have extrema at rational values of ν and with a finite depth

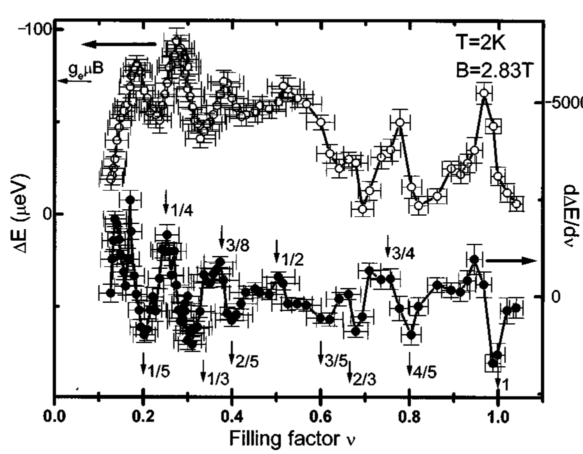

FIG. 3. Top: $\Delta E(\nu)$ (open circles) measured for a fixed value of $B=2.83$ T and varying $1.0<n_e(10^{10}\ \text{cm}^{-2})<7.1$. The value of $g_e \mu B$, with the bare electron g factor $g_e=-0.44$ (GaAs), is indicated. Bottom: The numerical derivative of the measured ΔE with respect to ν (solid circles). Solid lines connect the experimental points.

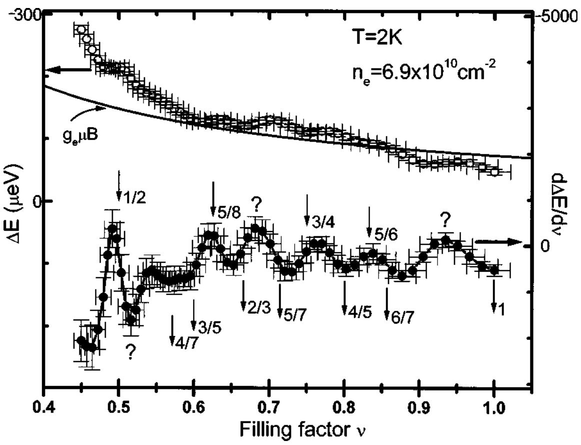

FIG. 4. Top: $\Delta E(\nu)$ (smoothed and shown by open circles) measured for a fixed value of $n_e = 6.9\times10^{10}\ \mathrm{cm}^{-2}$ and varying $2.83\ \mathrm{T}\leqslant B\leqslant 6.40\ \mathrm{T}$, $g_e\mu B$ (with $g_e=-0.44$) is also shown (solid line). Bottom: The numerical derivative of the measured ΔE with respect to ν (solid circles).

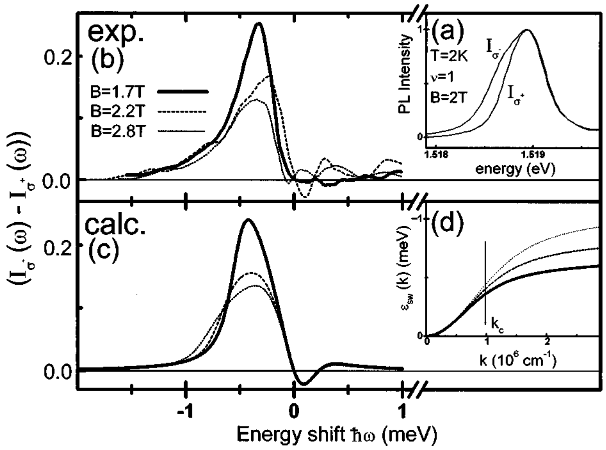

FIG. 5. (a) The σ^-- and σ^+-polarized PL spectra of the X_s^- transitions, measured at $\nu=1$ and normalized to their maximum intensity. (b) $[I_{\sigma^-}(\omega)-I_{\sigma^+}(\omega)]$ extracted from (a) for various values of the magnetic field. The energy shift $\hbar\omega$ is measured with respect to the spectrum energy maximum (virtually the same for both polarizations). (c) $[I_{\sigma^-}(\omega)-I_{\sigma^+}(\omega)]$ calculated using Eq. (1) for the same values of B as in (b). (d) The SW dispersion relations calculated following Ref. 21 for the same values of B as in (b), (c).

directly related to corresponding FQHE gaps.

In Fig. 4 we show another measurement of $\Delta E(\nu)$ carried out for a fixed value of $n_e=6.9\times10^{10}\ \mathrm{cm}^{-2}$ and various $2.83\ \mathrm{T}\leqslant B\leqslant 6.40\ \mathrm{T}$. These results confirm the existence of oscillations with ν, which are superposed on the monotonous increase of ΔE with increasing B. As in Fig. 3, the observed extrema of $\partial\Delta E/\partial\nu$ show a correspondence with rational ν values. However, in contrast to Fig. 3, there are few well-resolved extrema which are not assigned (denoted by a question mark in Fig. 4). Their ν values deviate from the rational fractions with the lowest possible denominator by more than the experimental error range (horizontal bars) of the 2DEG density.[14] Also, in this case, the depth of the observed extrema does not obey the odd denominator hierarchical sequence. These discrepancies may result from a significant contribution to the observed $\Delta E(\nu)$ of the (B-dependent) hole Zeeman energy.

In order to understand qualitatively the relation between $\varepsilon_{\mathrm{SF}}$, introduced by these spectroscopic measurements, and the ground state of the 2DEG at a fractional ν, we consider first the case of $\nu=1$ for which the effect of collective SF excitations is known to be enhanced[19] by the 2DEG exchange interactions. This is done by analyzing the line shape of the X_s^- polarized PL spectra observed for various values of $B\leqslant 2.8$ T and of n_e corresponding to $\nu=1$ (Fig. 5). In this case of a fully occupied spin-↑ level, a single SF process in the final state of the σ^--polarized recombination gives rise to a dispersive spin-wave (SW) collective excitation of the 2DEG. A clear indication of such a cooperative transition involving the X_s^- and a finite **k** SW is the pronounced low-energy σ^--polarized PL tail, observed in the vicinity of $\nu=1$ (Ref. 20) [Fig. 5(a)]. In order to resolve this feature, we take the difference between the $I_{\sigma^-}(\omega)$ and $I_{\sigma^+}(\omega)$ spectra (each one is unity normalized at $\omega=0$), and compare the obtained line shapes [shown in Fig. 5(b)] with those calculated by

$$I_{\sigma^-}(\omega)-I_{\sigma^+}(\omega)=N_-\int d^2\mathbf{k}M(k)\,\delta[\varepsilon_{\mathrm{SW}}(k)-\hbar\omega] -N_+\delta(\hbar\omega). \qquad (1)$$

Here,

$$M(k)=\frac{1}{\pi k_c^2}\exp\left(-\frac{k^2}{k_c^2}\right)$$

is the squared optical matrix element representing in this simplest phenomenological form the overlap between the X_s^- and the SW wave functions. k_c introduces a cutoff in k space and is used as a fitting parameter. The SW dispersion relation[21] is

$$\varepsilon_{\mathrm{SW}}(k)=\Delta\left[1-\exp\left(-\frac{k^2\lambda^2}{4}\right)I_0\left(\frac{k^2\lambda^2}{4}\right)\right],$$

where I_0 stands for the zeroth-order modified Bessel function. The electron and the hole Zeeman energies are neglected for the low values of $B\leqslant 2.8$ T used here. For a numerical evaluation of Eq. (1) the δ functions are replaced by Lorentzians with a half-width $\Gamma\simeq 0.2$ meV that reflects the observed σ^+-polarized transition broadening. $N_\pm$ are the normalization factors.

The assignment of the σ^--polarized PL tail to the X_s^- recombination coupled to an emission of the SW dispersive excitation is confirmed by the line-shape dependence on the magnetic field, which scales both the SW energy and **k**. Indeed, the observed decrease in the $[I_{\sigma^-}(\omega)-I_{\sigma^+}(\omega)]$ intensity maximum with increasing B [Fig. 5(b)] is found to be in agreement with the calculated line shapes [Fig. 5(c)]. This is because, as seen from Eq. (1), I_{σ^-} is inversely proportional (in the limit of $\Gamma\to 0$) to the SW dispersion slope [Fig. 5(d)], which increases with increasing B. The fitted value of $k_c \simeq 1\times10^6\ \mathrm{cm}^{-1}$ determines the energy shift of $\simeq -0.3$ meV of the $[I_{\sigma^-}(\omega)-I_{\sigma^+}(\omega)]$ maximum, (that is nearly independent of B) as well as its intensity variation with B. Also, the calculated line shapes occur to be nearly insensitive of a specific analytic form chosen for the matrix element.

It is worth noticing that the SF excitation with $k=0$ carries no information about electron-electron interactions at any ν, that is a consequence of quite a general (Larmor's) theorem. The coupling mechanism of a finite **k** SW excitation to the final state of the X_s^- recombination might be due to the following: (1) Short-range potential scattering (by QW interface roughness, for instance); (2) magnetic-field-induced large **k** mixing in the process $X_s^- +$spin-hole→SW +photon. The second mechanism requires that the spin-hole in the initial state, resulting from a vacancy of the additional X_s^- electron, is spatially well separated from X_s^-. In this case, the magnetooptical transition virtually occurs from the bound (X_s^-) to free (SW) state and is accompanied by **k** nonconservation. Alternatively, this condition can be formulated in terms of global minimum at a large **k** value for the magnetic energy dispersion of the (electrically neutral) initial state. The magnetically induced **k** mixing is a plausible mechanism since the fitted value of k_c occurs to be close to λ^{-1} for the range of 1.7 T$\leqslant B \leqslant$2.8 T.

Extending this line-shape analysis to the FQHE state with an odd denominator fraction, the collective SF excitations, according to Laughlin's theory,[22] may be considered as carried by quasiparticles with a fractional charge e^*. Then, for fractions $\nu=\frac{1}{3},\frac{2}{3},\frac{1}{5},\frac{4}{5}$, the SF excitations are quasiparticle SW excitations corresponding to $\nu^*=1$, while for $\nu=\frac{2}{5},\frac{3}{5}$ they are quasiparticle spin-triplet excitations with $\nu^*=2$. In terms of these dispersive quasiparticle excitations, $\varepsilon_{\mathrm{SF}}$ represents the SF excitation energy at singular points in **k** space, where $d\varepsilon_{\mathrm{SF}}/dk=0$. Since, to the best of our knowledge, the SF dispersion in the FQHE regime has not been studied theoretically, we assume that such points exist as a result of mixing[23] with Coulomb gap modes.[24,25] An experimental indication for such a mixing has been provided[26–28] for $\nu=\frac{2}{3},\frac{3}{5}$ by FQHE transport measurements in tilted magnetic fields, where SF excitations were shown to contribute appreciably to the measured fractional energy gaps.

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[1] K. von Klitzing, G. Dorda, and M. Pepper, Phys. Rev. Lett. **45**, 494 (1980).
[2] D. C. Tsui, H. L. Störmer, and A. C. Gossard, Phys. Rev. Lett. **48**, 1559 (1982).
[3] A. J. Turberfield, S. R. Haynes, P. A. Wright, R. A. Ford, R. G. Clark, J. F. Ryan, J. J. Harris, and C. T. Foxon, Phys. Rev. Lett. **65**, 637 (1990).
[4] A. J. Turberfield, R. A. Ford, I. N. Harris, J. F. Ryan, C. T. Foxon, and J. J. Harris, Phys. Rev. B **47**, 4794 (1993).
[5] For a review, see I. V. Kukushkin and V. M. Timofeev, Adv. Phys. **45**, 185 (1996).
[6] D. Heiman, B. B. Goldberg, A. Pinczuk, C. W. Tu, A. C. Gossard, and J. H. English, Phys. Rev. Lett. **61**, 605 (1988).
[7] K. Kheng, R. T. Cox, Y. Merle d'Aubigne, Franck Bassani, K. Saminadayar, and S. Tatarenko, Phys. Rev. Lett. **71**, 1752 (1993).
[8] G. Finkelstein, H. Shtrikman, and I. Bar-Joseph, Phys. Rev. Lett. **74**, 976 (1995).
[9] H. Buhmann, L. Mansouri, J. Wang, P. H. Beton, N. Mori, L. Eaves, M. Henini, and M. Potemski, Phys. Rev. B **51**, 7969 (1995).
[10] A. J. Shields, M. Pepper, D. A. Ritchie, M. Y. Simmons, and G. A. C. Jones, Phys. Rev. B **51**, 18 049 (1995); **52**, R5523 (1995).
[11] Arza Ron, H. W. Yoon, M. D. Sturge, A. Manassen, and E. Cohen, Solid State Commun. **97**, 741 (1996).
[12] G. Finkelstein, H. Shtrikman, and I. Bar-Joseph, Phys. Rev. B **53**, R1709 (1996).
[13] A. J. Shields, M. Pepper, M. Y. Simmons, and D. A. Ritchie, Phys. Rev. B **52**, 7841 (1996).
[14] D. Gekhtman, E. Cohen, Arza Ron, and L. N. Pfeiffer, Phys. Rev. B **54**, 10 320 (1996).
[15] L. N. Pfeiffer, K. W. West, J. W. Baldwin, and P. Gammel, Appl. Phys. Lett. **61**, 1211 (1992).
[16] M. J. Snelling, E. Blackwood, C. J. McDonagh, R. T. Harley, and C. T. B. Foxon, Phys. Rev. B **45**, 3922 (1992).
[17] B. I. Halperin, Phys. Rev. Lett. **52**, 1583 (1984).
[18] A. H. McDonald, H. C. A. Oji, and K. L. Liu, Phys. Rev. B **34**, 2681 (1986).
[19] T. Ando and Y. Uemura, J. Phys. Soc. Jpn. **37**, 1044 (1974).
[20] The observed spectroscopy of the X_s^- transition at $\nu=1$ is given by D. Gekhtman, E. Cohen, Arza Ron, and L. N. Pfeiffer, in Proceedings of the 12th International Conference on the Application of High Magnetic Fields, edited by G. Landwehr, Würzburg, 1996 (to be published). At the same conference the theory of PL at $\nu=1$ was presented by N. R. Cooper and D. B. Chklovskii.
[21] The SW dispersion relation for the strong B case is given in C. Kallin and B. I. Halperin, Phys. Rev. B **30**, 5655 (1984). We use this relation with the $k\to\infty$ gap energy Δ taken from the low B transport experiment: A. Usher, R. J. Nicholas, J. J. Harris, and C. T. Foxon, Phys. Rev. B **41**, 1129 (1990). The measured value of the transport gap at $\nu=1$ allows for the effects of higher Landau-level mixing, disorder-induced overlap between spin-↑ and spin-↓ levels, and finite thickness of the QW. A possible contribution of charged spin texture excitations (which are not allowed in a single optical recombination process) to this measured value is neglected.
[22] R. B. Laughlin, Phys. Rev. Lett. **50**, 1395 (1983).
[23] V. M. Apalkov and E. I. Rashba, Phys. Rev. B **46**, 1628 (1992).
[24] S. M. Girvin, A. H. MacDonald, and P. M. Platzman, Phys. Rev. Lett. **54**, 581 (1985).
[25] A. Pinczuk, B. S. Dennis, L. N. Pfeiffer, and K. W. West, Phys. Rev. Lett. **70**, 3983 (1993).
[26] J. P. Eisenstein, H. L. Stormer, L. N. Pfeiffer, and K. W. West, Phys. Rev. B **41**, 7910 (1990).
[27] R. G. Clark, S. R. Haynes, J. V. Branch, A. M. Suckling, P. A. Wright, P. M. W. Oswald, J. J. Harris, and C. T. Foxon, Surf. Sci. **229**, 25 (1990).
[28] L. W. Engel, S. W. Hwang, T. Sajoto, D. C. Tsui, and M. Shayegan, Phys. Rev. B **45**, 3418 (1992).