Spin Configurations and Quasiparticle Fractional Charge of Fractional-Quantum-Hall-Effect Ground States in the $N=0$ Landau Level

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Spin configurations of fractions $1<\nu<2$ are examined by angular- and n-dependent activation studies. Energy gaps $\Delta(\theta)$ and intercepts $\sigma_{xx}^{c}(\theta)$ that probe the quasiparticle charge e^* quantify a dramatic difference between $\frac{4}{3}$ and $\frac{5}{3}$ states consistent with assignments $\frac{4}{3}\uparrow\downarrow$ and $\frac{5}{3}\uparrow\uparrow$ ($\uparrow\downarrow,\uparrow\uparrow$=zero, maximum polarization). A field-induced phase transition $\frac{4}{3}\uparrow\downarrow\rightarrow\frac{4}{3}\updownarrow$ (partial polarization) in which e^* changes from $e/3$ to $e/5$ is mapped out. The $\frac{7}{5}$ state formed from $e/3$ quasiparticles is destroyed at the $\frac{4}{3}$ transition. High-order assignments $\frac{7}{5}$ ($\uparrow\uparrow$ or $\updownarrow$), $\frac{8}{5}\uparrow\downarrow$, $\frac{10}{7}\uparrow\downarrow$, and $\frac{11}{7}\updownarrow$ are consistent with experiment.

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Following the discovery and initial explanation of the fractional quantum Hall effect[1,2] (FQHE), it was suggested that electrons in some FQHE ground states at Landau-level filling $\nu<1$ in GaAs-GaAlAs structures may have reversed spins due to the small GaAs g factor.[3] Subsequently, there has been mounting theoretical evidence for spin-unpolarized ground states in the $N=0$ Landau level[4,5] and spin-reversed quasiparticle excitations.[6-8] This work has had little impact, as it is believed that the Zeeman energy at magnetic fields B of fractional states $\nu=p/q$ ($B=nh/e\nu$) forces maximum polarization of both the ground and excited states. In this Letter, we show that for the hierarchy of fractions[9] $1<\nu<2$ which occur at *low B*, our activation data are consistent with the following spin assignments:

$$\begin{array}{ccccc} \uparrow\uparrow\frac{5}{3} & & & \frac{4}{3}\uparrow\downarrow & \\ & / & & \backslash & \\ \uparrow\downarrow\frac{8}{5} & & & & \frac{7}{5}\uparrow\uparrow \text{ or } \updownarrow \\ / & & & & \backslash \\ \updownarrow\frac{11}{7} & & & & \frac{10}{7}\uparrow\downarrow\,, \end{array} \tag{1}$$

where $\uparrow\downarrow$, $\updownarrow$, and $\uparrow\uparrow$ signify zero (unpolarized), partial, and maximum polarization. Equation (1) is in exact agreement with extensive finite-size calculations by Maksym[10] centered on these experiments.

The spin configuration of fractional states is probed by changing the field at which they occur, by tilting B at an angle θ to the sample plane normal (increasing its absolute value at *fixed* n, $B_\perp$), and by increasing the electron concentration n (at $\theta=0°$) so that the fraction occurs at higher $B_\perp$. As B is increased, spin-unpolarized states undergo transitions $\uparrow\downarrow\rightarrow\updownarrow\rightarrow\uparrow\uparrow$, whereas configurations with maximum polarization are unchanged. This means of identifying spin assignments is quantified using our probe $\sigma_{xx}^{c}=\sigma_{xx}(1/T=0)=c(e/q)^2/h$ (c is a constant close to 1) of the quasiparticle charge $e^*=\pm e/q$ (Ref. 11) obtained from extrapolated values of ρ_{xx}^{c} defined by $\rho_{xx}=\rho_{xx}^{c}e^{-\Delta/kT}$, where Δ is the energy gap of the fractional state $\nu=p/q$. Our main observations for the evolution of fractional structure on rotating our samples *in situ* from $\theta=0°$ to 70° are as follows: We map out a "destruction" of the proposed $\frac{4}{3}\uparrow\downarrow$ state [$=(\frac{2}{3}\uparrow,\frac{2}{3}\downarrow)$; partial filling factors $\nu\uparrow$ and $\nu\downarrow$ are defined in Ref. 4] characterized by $\Delta_{4/3}(\theta=0°)=340$ mK, $\sigma_{xx}^{c}=0.99(e/3)^2/h$, in which $\Delta_{4/3}$ falls to zero followed by a reemergence with $\Delta_{4/3}=80$ mK, $\sigma_{xx}^{c}=1.1(e/5)^2/h$ for $\theta>60°$. In contrast, the $\frac{5}{3}$ state characterized by $\Delta_{5/3}(\theta=0°)=330$ mK, $\sigma_{xx}^{c}=0.81(e/3)^2/h$ is relatively unaffected over the entire angular range, consistent with a $\frac{5}{3}\uparrow\uparrow[=(\frac{3}{3}\uparrow,\frac{2}{3}\downarrow)]$ assignment, i.e., maximum possible polarization. The $\frac{4}{3}$ results are interpreted as a phase transition from $\uparrow\downarrow$ with $e^*=e/3$ to a partially polarized ($\updownarrow$) state with $e^*=e/5$; the $\updownarrow$ configuration is speculative but $\frac{4}{3}\updownarrow=(\frac{14}{15}\uparrow,\frac{2}{5}\downarrow)$ would be consistent with the σ_{xx}^{c} data. This is supported by the $\frac{7}{5}$ development; $\Delta_{7/5}\sim70$ mK, $\sigma_{xx}^{c}=0.98(e/5)^2/h$ which show little variation with θ up to the $\frac{4}{3}$ transition consistent with a $\frac{7}{5}\uparrow\uparrow=(\frac{5}{5}\uparrow,\frac{2}{5}\downarrow)$ [or $\frac{7}{5}\updownarrow=(\frac{4}{5}\uparrow,\frac{3}{5}\downarrow)$] assignment, at which point it is destroyed. In the hierarchical model, the daughter $\frac{7}{5}$ state forms from $e/3$ charged quasiparticles of the $\frac{4}{3}$ parent. The fundamental change in e^* from $e/3$ to $e/5$ at the $\frac{4}{3}$ transition, as measured by σ_{xx}^{c}, destroys the basis of the $\frac{7}{5}$ state and provides support for the interdependence of high-order quantum fluids of fractionally charged quasiparticles.[9]

We present a tilted-field study of GaAs-$Ga_{0.68}Al_{0.32}As$ heterojunctions G156A and G156B (800-Å spacer layer) with precise Hall bar geometry, at saturated concentration $n_{sat}=1.6\times10^{11}$ cm^{-2}, $\mu=1.6\times10^{6}$ cm^2/V s. Sample G156A was rotated *in situ* on a pivoted platform attached to the laminated Cu cold finger of a dilution refrigerator in a 16.5-T hybrid magnet. Using a long glass mixing chamber tail, G156B was mounted in the dilute phase at field center of a 20-T hybrid magnet at fixed angles $\theta=0°$, 47°, and 60°. The reproducibility of these

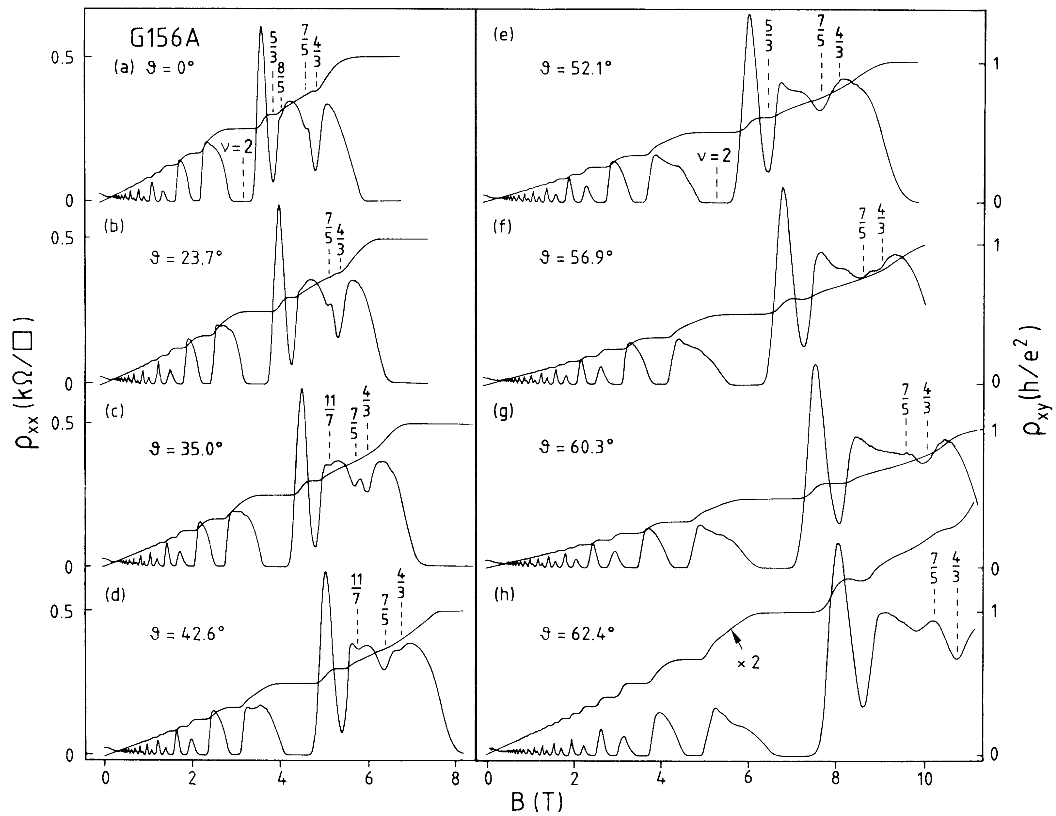

FIG. 1. Angular development of G156A $1<\nu<2$ fractional structure at $T=120$ mK ($n=1.6\times10^{11}$ cm^{-2}).

data was checked by successive cooldowns from room temperature. The "in-phase" sample current was 40.0 nA using an ac lock-in technique and thermometry was provided by a Speer resistor calibrated and corrected for magnetoresistance, mounted identically and next to the sample at field center in both G156A and G156B assemblies. The G156 data are compared with a $\theta=0°$ n-dependent study of a higher-density sample G71 in which $1<\nu<2$ fractions can be moved to significantly higher fields by persistent photoexcitation using a red light-emitting diode.

ρ_{xx} and ρ_{xy} data for sample G156A at 120 mK are shown in Figs. 1(a)–1(h) for eight selected angles. For $1<\nu<2$, at $\theta=0°$ strong $\frac{4}{3}$ and $\frac{5}{3}$ "parent" ρ_{xx} structure with well-defined Hall plateaus and weaker $\frac{7}{5}$ and $\frac{8}{5}$ daughter states are observed. As θ is increased to $\sim45°$, the $\frac{4}{3}$ (and $\frac{8}{5}$) ρ_{xx} and ρ_{xy} structure is systematically destroyed whereas the $\frac{5}{3}$ state is unaffected, the $\frac{7}{5}$ ρ_{xx} minimum which remains at approximately the same depth *apparently* strengthens due to the increasing edge on the $\frac{4}{3}$ side, and an $\frac{11}{7}$ state is brought out. At $\theta=52°$, the $\frac{4}{3}$ state is completely absent and the $\frac{11}{7}$ ρ_{xx} minimum begins to disappear while the well-defined $\frac{5}{3}$ and $\frac{7}{5}$ structure persists. Beyond this angle a weaker $\frac{4}{3}$ state reemerges coincident with a destruction of the $\frac{7}{5}$ state, and for $\theta>60°$ only the parent $\frac{5}{3}$ and reemergent $\frac{4}{3}$ structure remains, with a shallow ρ_{xx} minimum close to $\frac{3}{2}$. The ρ_{xy} scale is doubled in trace 1(h) to show the recovery of a plateau at $\nu=\frac{4}{3}$, although this feature and a now distorted $\frac{5}{3}$ plateau are higher than $\frac{3}{4}$ and $\frac{3}{5}$ times h/e^2, respectively, possibly due to contact problems at high angle. Similar $1<\nu<2$ data are obtained

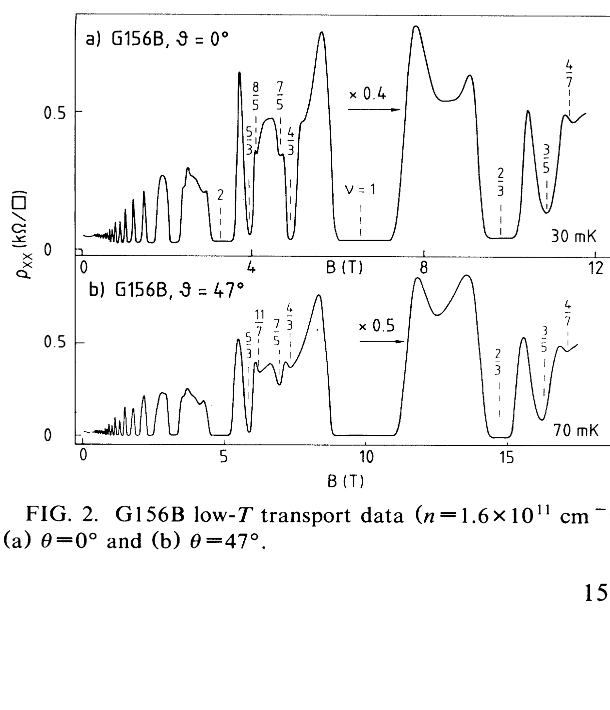

FIG. 2. G156B low-T transport data ($n=1.6\times10^{11}$ cm^{-2}); (a) $\theta=0°$ and (b) $\theta=47°$.

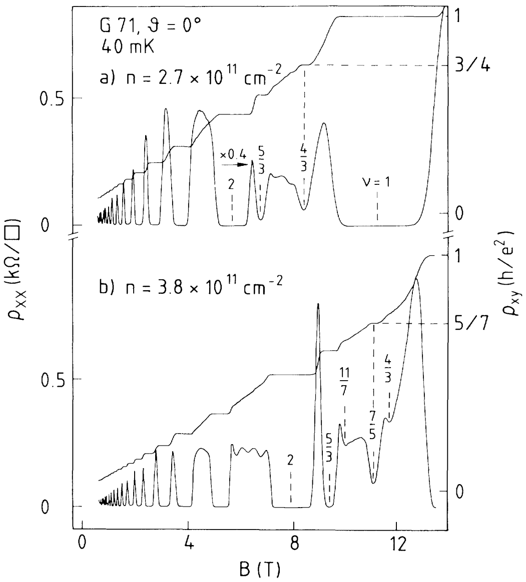

FIG. 3. G71 low-T transport data ($\theta=0°$); (a) $n=2.7\times10^{11}$ and (b) $n=3.8\times10^{11}$ cm^{-2}.

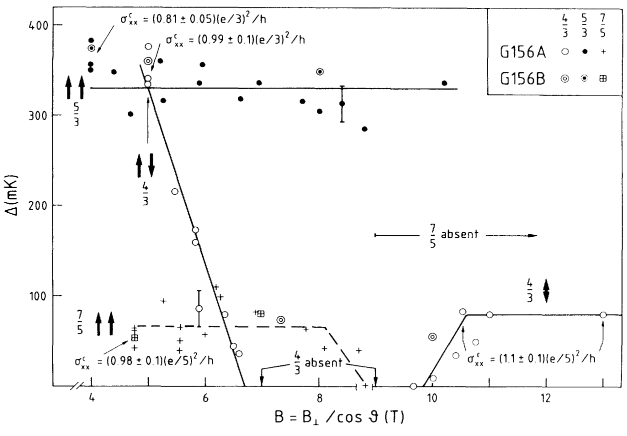

FIG. 4. G156A and G156B $\nu=\frac{4}{3}$, $\frac{5}{3}$, and $\frac{7}{5}$ angular-dependent activation data $\Delta(\theta)$, $\sigma_{xx}^c(\theta)$.

for sample G156B where a $\theta=0°$ and 47° ρ_{xx} comparison is extended to $\nu<1$ structure at high $B_\perp$, which in contrast displays no significant differences as shown in Figs. 2(a) and 2(b).

Almost identical effects are induced by increasing n with $\theta=0°$, as exemplified by data for sample G71 in Fig. 3, taken from a study where n was increased in small steps. The G71 results have been confirmed in several samples,[12] although the full implication of these early data was not realized. Corresponding n-dependent activation data are very similar to tilted-field studies, which indicates that for unpolarized states the effect of the increased total field on the spin configuration is the important factor, of more significance than the effect of $B_\parallel$ on the z extent of the electron wave function in the quasi-2D layer. The destruction of the G71 $\frac{4}{3}$ plateau concomitant with the emergence of a strong $\frac{7}{5}$ plateau presents severe difficulties for the hierarchical model if spin reversal is ignored.

The situation is quantified by activation studies to determine Δ and σ_{xx}^c in the Fig. 1 evolution. Results for samples G156A and G156B are summarized in Fig. 4. The disparity between $\frac{4}{3}$ and $\frac{5}{3}$ states, believed to be identical if they occur at the same field in the (fully polarized) hierarchical framework, is evident. $\Delta_{4/3}$ collapses from 340 mK at $\theta=0°$ ($B_\perp=5$ T) to zero over the range $B=B_\perp/\cos\theta=5$–6.7 T; the $\frac{4}{3}$ state is absent from $B=7$–9 T and reemerges with a constant $\Delta\sim80$ mK for $B>10.5$ T (verified up to 16 T). In contrast, $\Delta_{5/3}$ remains essentially unchanged at ~330 mK over this entire angular range. Within errors, $\Delta_{7/5}$ is scattered about an average ~70 mK which falls to zero between 8 and 9 T when the $\frac{4}{3}$ reemergent state forms. The $\frac{7}{5}$ state is then absent beyond 9 T. This development is shown by the full and dashed lines in Fig. 4, which are guides to the eye. In our Δ analysis, straight-line fits were made to the raw activation data without hopping subtraction. While low-T deviation from straight-line behavior of the ρ_{xx} vs $1/T$ plots was not so significant for the $\frac{4}{3}$ data and the $\frac{5}{3}$ plots at small angles, the $\frac{5}{3}$ hopping component systematically increased with θ. Although this can be subtracted, such fits are model dependent. The effect will be to increase $\Delta_{5/3}$ with θ, making the contrast between $\frac{5}{3}$ and $\frac{4}{3}$ behavior even more striking, and this aspect, of relevance to the nature of the gap, will be addressed separately following studies of the hopping regime at lower temperatures. Extrapolation of the raw data does *not* affect the σ_{xx}^c intercept as discussed in Ref. 11. The G156B data are similar to the comprehensive G156A results, except that $\Delta_{5/3}=530$ mK was obtained at $\theta=47°$, which is not shown.

The σ_{xx}^c data are also summarized in Fig. 4. In our previous study of a low-density sample, G139,[11] we found that $\sigma_{xx}^c=c(e/q)^2/h$, where $c=0.91\pm0.11$ for thirteen fractional states $\nu=p/q$, providing compelling evidence that σ_{xx}^c is an experimental probe of $e^*=\pm e/q$. This conclusion is confirmed for sample G156A at $\theta=0°$, $\sigma_{xx}^c(\frac{4}{3})=(0.99\pm0.1)(e/3)^2/h$, $\sigma_{xx}^c(\frac{5}{3})=(0.81\pm0.05)(e/3)^2/h$, and $\sigma_{xx}^c(\frac{7}{5})=(0.98\pm0.1)(e/5)^2/h$. Over the angular range in which $\Delta_{4/3}$ collapses to zero, $\sigma_{xx}^c(\frac{4}{3})$ falls between $(e/3)^2/h$ and $(e/5)^2/h$. In the reemergent $\frac{4}{3}$ phase beyond $B=10$ T, however, $\sigma_{xx}^c(\frac{4}{3})=(1.1\pm0.1)(e/5)^2/h$ and the new $\frac{4}{3}$ phase takes on the characteristics of a $\frac{7}{5}$ state. As θ is increased to 70°, $\sigma_{xx}^c(\frac{5}{3})$ falls by $\sim20\%$ and the average of all angles is $0.69(e/3)^2/h$. However, with increasing angle the temperature range of the activated $\frac{5}{3}$ ρ_{xx} data is progressively restricted due to the onset of hopping at higher temperatures which leads to extrapolation errors that might account for this falloff. Over the angular range up to the onset of the reemergent $\frac{4}{3}$ phase which correlates with the disappearance of the $\frac{7}{5}$ state, $\sigma^c_{xx}(\frac{7}{5})$ remains constant at $0.98(e/5)^2/h$. The σ^c_{xx} results for $1<\nu<2$ are very similar in sample G156B.[13] In contrast, for the high-field fractions $\frac{2}{3}$, $\frac{3}{5}$, and $\frac{4}{7}$ studied in G156B no such effects are induced by the tilted field; the activation energies of these states simply strengthen and for $\theta=\{0°,47°\}$, $\sigma^c_{xx}(\frac{2}{3})=\{1.02,1.02\}(e/3)^2/h$, $\sigma^c_{xx}(\frac{3}{5})=\{0.7,0.84\}(e/5)^2/h$, and $\sigma^c_{xx}(\frac{4}{7})=\{0.73,0.7\}(e/7)^2/h$. Arrhenius plots for sample G156B are presented in Ref. 13.

The $1<\nu<2$ assignments of Eq. (1) are therefore based on the following conclusions.

(i) At low B ($\theta=0°$), $\Delta_{4/3}\sim\Delta_{5/3}$ and $e^*=e/3$ for both states. As B is increased, the $\frac{4}{3}$ state is destroyed whereas the $\frac{5}{3}$ state is largely unaffected and with hopping subtraction $\Delta_{5/3}$ will increase with θ. Using arguments similar to those of Eisenstein *et al.*[14] in their study of the $\nu=\frac{5}{2}$ polarization, destruction of $\frac{4}{3}$ structure is consistent with $\uparrow\downarrow$ as changing the Zeeman energy forces the system into a different spin state. The change in e^* from $e/3$ to $e/5$ and the reduced value of $\Delta_{4/3}$ for the reemergent state is interpreted as a transition from $\uparrow\downarrow$ to $\updownarrow$. We speculate that $\frac{4}{3}\updownarrow=(\frac{14}{15}\uparrow,\frac{2}{5}\downarrow)$ as Δ and σ^c_{xx} for $\frac{4}{3}\updownarrow$ and $\frac{7}{5}$ data are very similar and this assignment closely resembles $\frac{7}{5}\uparrow\uparrow=(\frac{15}{15}\uparrow,\frac{2}{5}\downarrow)$ [although it is possible that the $\frac{7}{5}$ state is only partially polarized $(\frac{4}{5}\uparrow,\frac{3}{5}\downarrow)$].

(ii) The daughter $\frac{7}{5}$ state with $e^*=e/5$ is largely unaffected on increasing B *until* the parent $\frac{4}{3}$ transition in which e^* of the $\frac{4}{3}$ state changes to $e/5$. The hierarchical basis of the $\frac{7}{5}$ state is then absent and it can no longer exist, as observed.

(iii) $\uparrow\downarrow$ assignments for $\frac{8}{5}$ and $\frac{10}{7}$ are more tentative and are based on the disappearance of $\frac{8}{5}$ structure as B is increased and the complete absence of $\frac{10}{7}$ structure at higher fields than the observed $\frac{11}{7}$ state.

(iv) For the emergent $\frac{11}{7}$ structure, $\Delta_{11/7}\simeq 30$ mK and $\sigma^c_{xx}=(1.07\pm0.1)(e/5)^2/h$ (not shown in Fig. 4 for clarity, see Ref. 13). We tentatively conclude that the polarization is therefore only partial and speculate that $\frac{11}{7}\updownarrow=(\frac{34}{35}\uparrow,\frac{3}{5}\downarrow)$ would similarly be consistent with σ^c_{xx}. The $\frac{8}{5}\uparrow\downarrow$ and $\frac{11}{7}\updownarrow$ interplay (Figs. 1 and 2) is analogous to the $\frac{4}{3}\uparrow\downarrow$ and $\frac{7}{5}\uparrow\uparrow$ situation, which leads to the disappearance of $\frac{11}{7}$ structure at high B.

(v) At high B, only the parent $\frac{5}{3}$ and the reemergent $\frac{4}{3}$ structure remains, since the basis of the $\frac{4}{3}$ hierarchy is swept away by the change in e^* at the $\frac{4}{3}$ transition and states beyond the $\frac{8}{5}\uparrow\downarrow$ state in the $\frac{5}{3}$ hierarchy will be similarly affected.

Fractions $\nu<1$ occur at high $B_\perp$ in our samples which precludes ground-state spin reversal. However, projected $\nu<1$ assignments that could be realized in low-density samples can be extrapolated from electron-hole symmetry. With the inclusion of spin but *not* Landau-level mixing, this is between states ν and $2-\nu$ and from Eq. (1) anticipated $\nu<1$ configurations at low $B_\perp$ are $\frac{1}{3}\uparrow\uparrow$, $\frac{2}{5}\uparrow\downarrow$; $\frac{2}{3}\uparrow\downarrow$, $\frac{3}{5}\uparrow\uparrow$ (or $\updownarrow$). The $\nu<1$ and $1<\nu<2$ comparison is not so simple, however. In our angular study, $\nu=\frac{4}{3}$ is moved to absolute fields B comparable to $B_\perp$ of $\nu<1$ states and maximum polarization of the $\frac{4}{3}$ state $[(\frac{3}{3}\uparrow,\frac{1}{3}\uparrow)]$ is not achieved. At $\nu=\frac{4}{3}$, $B_\perp$ is small (5 T) and the cyclotron energy is unchanged in the angular study. Landau-level-mixing effects might therefore be important and will break electron-hole symmetry with $\nu=\frac{2}{3}$. The effect of $B_\parallel$ further complicates the $\frac{2}{3}$ and $\frac{4}{3}$ comparison and more work is required to resolve these problems.

Finally, the configurational sequence in Eq. (1) can be understood from a hierarchical argument[4] and to simplify matters we refer to the projected $\nu<1$ low-B assignments. Starting with a $\frac{1}{3}\uparrow\uparrow$ parent state, for n electrons the total spin is $n/2$. In Haldane's notation[9] the daughter state $\frac{2}{5}=1/(m\pm\alpha/p')=1/(3-\frac{1}{2})$. The number of quasiparticles that form the $\frac{2}{5}$ state is $n/p'=n/2$. If all quasiparticle excitations are spin reversed, which is anticipated at low B,[6–8,10,13] the $\frac{2}{5}$ state will then be unpolarized. A $\frac{2}{3}\uparrow\downarrow$ state (at low B) can similarly be obtained starting from $\nu=1\uparrow\uparrow$ and noting that $\frac{2}{3}=1/(1+\frac{1}{2})$. This scheme can be extended down the series of Eq. (1).

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