Collective Nature of the Reentrant Integer Quantum Hall States in the Second Landau Level

N. Deng,[1] A. Kumar,[1] M. J. Manfra,[1,2] L. N. Pfeiffer,[3] K. W. West,[3] and G. A. Csáthy[1,*]

[1]*Department of Physics, Purdue University, West Lafayette, Indiana 47907, USA*

[2]*Birck Nanotechnology Center, School of Materials Engineering and School of Electrical and Computer Engineering, Purdue University, West Lafayette, Indiana 47907, USA*

[3]*Department of Electrical Engineering, Princeton University, Princeton, New Jersey 08544, USA*

(Received 18 October 2011; published 23 February 2012)

We report an unexpected sharp peak in the temperature dependence of the magnetoresistance of the reentrant integer quantum Hall states in the second Landau level. This peak defines the onset temperature of these states. We find that in different spin branches the onset temperatures of the reentrant states scale with the Coulomb energy. This scaling provides direct evidence that Coulomb interactions play an important role in the formation of these reentrant states evincing their collective nature.

DOI: 10.1103/PhysRevLett.108.086803 PACS numbers: 73.43.−f, 73.20.Qt, 73.21.Fg

The density of states of a two-dimensional electron gas (2DEG) in a perpendicular magnetic field B consists of a set of discrete energy levels called the Landau levels. Since spin degeneracy is lifted, each orbital Landau level splits into two distinct energy levels. The second orbital Landau level (SLL) thus corresponds to Landau level filling factor ν in the $2 < \nu < 4$ range.

The SLL is astonishingly rich in novel ground states [1–3]. Recent experiments [3–9] suggest that in this region there are both fractional quantum Hall states (FQHS) of free composite fermions [10,11] as well as exotic FQHSs [12–15]. The study of the latter has enriched quantum many-body physics with numerous novel concepts such as paired composite fermion states with Pfaffian correlations, non-Abelian quasiparticles [12–21], and topologically protected quantum computing [22].

The eight reentrant integer quantum Hall states (RIQHSs) form another set of prominent ground states in the SLL [1]. The transport signatures of the RIQHSs are consistent with electron localization in the topmost energy level [1]. However, the nature of the localization is not yet well understood. Depending on the relative importance of the electron-electron interactions, the ground state can be either an Anderson insulator or a collectively pinned electron solid.

FQHSs owe their existence to the presence of the interelectronic Coulomb interactions [10,11]. Since FQHSs and RIQHSs alternate in the SLL, it was argued that Coulomb interactions must be important and, therefore, the RIQHSs in the SLL must be electron solids [1]. Subsequent density matrix renormalization group [23] and Hartree-Fock calculations [24] also favored the electron solid picture and predicted the solid phase similar to the Wigner crystal, but having one or more electrons in the nodes of the crystal [24]. Recently reported weak microwave resonances in one such RIQHS are suggestive of but are far from being conclusive on the formation of a collective insulator [25]. Our understanding of the RIQHSs in the SLL, therefore, is still in its infancy and the collective nature of these states has not yet been firmly established.

In this Letter we report a feature in the temperature dependent magnetoresistance which so far has only been observed in the RIQHSs in the SLL. We use this feature to define the onset temperature of the RIQHSs. The scaling of the onset temperatures with the Coulomb energy reveals that Coulomb interactions play a central role in the formation of RIQHSs and, therefore, these reentrant states are exotic electronic solids rather than Anderson insulators. We also report an unexpected trend of the onset temperatures within each spin branch. This trend is inconsistent with current theories and can be understood as a result of a broken electron-hole symmetry. Explaining such a broken symmetry of the RIQHSs is expected to impact our understanding of a similar asymmetry of the exotic FQHSs of the SLL, including the one at $\nu = 5/2$.

We performed magnetotransport measurements on a high quality GaAs/AlGaAs sample of density $n = 3.0 \times 10^{11}$ cm^{-2} and of mobility $\mu = 3.2 \times 10^7$ cm^2/Vs. Earlier we reported the observation of a new FQHS at $\nu = 2 + 6/13$ in this sample [3]. The sample is immersed into a He-3 cell equipped with a quartz tuning fork viscometer used for B-field independent thermometry [26].

In Fig. 1(a) we show the dependence of the Hall resistance R_{xy} in the SLL on B (bottom scale) and on the Landau level filling factor $\nu = hn/eB$ (top scale). Here h is Planck's constant and e the electron charge. For $2 < \nu < 3$ only the lower of the two spin-split energy levels of the SLL is occupied, hence the term lower spin branch. $3 < \nu < 4$ corresponds to the occupation of the upper spin branch. Shown in Fig. 1(a) there are several regions of ν for which R_{xy} has plateaus quantized to an integer, i.e., $R_{xy} = h/ie^2$ with $i = 2$, 3, and 4. The three wide plateaus at $\nu < 2.17$, $2.83 < \nu < 3.17$, and $\nu > 3.83$ are quantized to $h/2e^2$, $h/3e^2$, and $h/4e^2$, respectively, and contain the filling factors $\nu = 2$, 3, and 4 marked by arrows. These are the well known integer quantum Hall plateaus. In contrast, other eight plateaus, shaded in Fig. 1(a), are quantized to an integer but form at a filling factor range which does not contain an integer ν. For example, the region labeled $R2a$ exhibits $R_{xy} = h/2e^2$ and it stretches between $2.26 < \nu < 2.32$, a region which does not contain $\nu = 2$. These eight states are the RIQHSs [1] and we label the ones located between $2 < \nu < 3$ with $R2a$, $R2b$, $R2c$, and $R2d$ and the ones between $3 < \nu < 4$ with $R3a$, $R3b$, $R3c$, and $R3d$. RIQHSs have historically been predicted [27] and observed [28,29] in high Landau levels (i.e., $\nu > 4$), but in contrast to the SLL, in high Landau levels there are only four RIQHSs in each Landau level.

Because of the delicate nature of the RIQHS in the SLL [1–6,30–35] there is only scarce information available on their temperature dependence [30,31,34]. In Figs. 1(b) and 1(c) we show the detailed temperature evolution of the longitudinal resistance R_{xx} and R_{xy} of $R2b$, respectively. The $R_{xx}(B)|_{T=6.9\text{ mK}}$ curve has a wide zero flanked by two sharp spikes. As the temperature is raised, the spikes in R_{xx} persist but they move closer to each other and the width of the zero decreases. At 32.6 mK $R_{xx}(B)$ does still exhibit the two spikes but instead of a zero it has a nonzero local minimum. The location in B field of this minimum is T independent and it defines the center $\nu_c = 2.438$ of the $R2b$ state. At 35.7 mK the two spikes of $R_{xx}(B)$ have moved closer to each other and between them there is still a local minimum, albeit with a large resistance. A small increase in T of only 2 mK leads to a qualitative change. Indeed, in contrast to curves at lower T, $R_{xx}(B)|_{T=37.7\text{ mK}}$ exhibits a single peak only. As the temperature is further raised, this single peak rapidly decreases until it merges into a low resistance background. Simultaneously with the described changes of R_{xx}, R_{xy} evolves from the quantized value $h/2e^2$ to its classical value $B/ne = h/\nu_c e^2$.

The behavior seen in Fig. 1 can be better understood by measuring T dependence at a fixed ν. In Fig. 2 we show R_{xy} versus T near the center ν_c of the various RQIHSs. It is found that R_{xy} assumes the classical Hall resistance at high T and it is quantized to $h/2e^2$ or $h/3e^2$ at low T. Since 80% of the change in R_{xy} between these two values occurs over only 5 mK, this change is very abrupt and it clearly separates the RIQHS at low T from the classical gas at high T. We interpret the inflection point in R_{xy} versus T as being the onset temperature T_c of the RIQHS. For reliable measurements in the vicinity of T_c the temperature is swept slower than 10 mK/h.

A transition from the classical Hall value to a quantized R_{xy} with decreasing T is observed not only for the RIQHSs in the SLL but also in the vicinity of any developed integer or fractional quantum Hall state and it is due to localization in the presence of a B field. As seen in Fig. 2, the $R_{xx}(T)|_{\nu=\text{fixed}}$ curves for the RIQHSs are nonzero at high T, they vanish at low T, and they exhibit a sharp peak at the onset temperature T_c defined above. In contrast, $R_{xx}(T)|_{\nu=\text{fixed}}$ of a quantum Hall state changes monotonically, without the presence of a peak. We found no reports in the literature of a similar peak in any other ground state of the 2DEG. The sharp peak in $R_{xx}(T)|_{\nu=\text{fixed}}$ is, therefore, a signature of localization so far only observed in the RIQHSs of the SLL and the peak temperature can be used as an alternative definition for the onset temperature T_c.

Figure 3 represents the stability diagram of the RIQHSs in the ν^*-T plane. Here $\nu^* = \nu - 2(3)$ is the partial filling

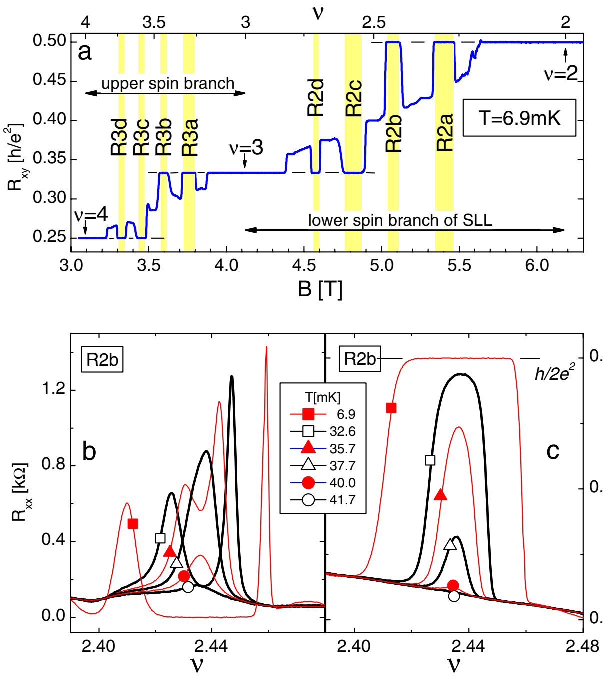

FIG. 1 (color online). The Hall resistance of the eight RIQHSs in the SLL at 6.9 mK [panel (a)] and the temperature evolution of the RIQHS labeled $R2b$ (bottom panels).

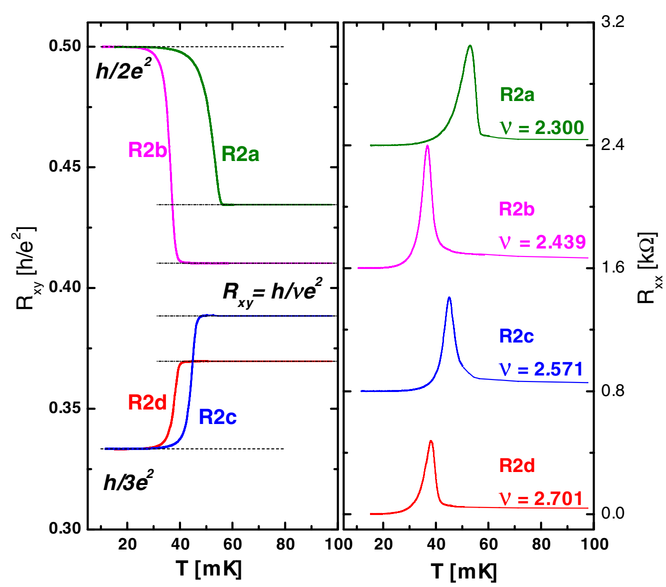

FIG. 2 (color online). The evolution of the magnetoresistance of RIQHSs of the lower spin branch with temperature near the center ν_c of each RIQHS. For clarity, $R_{xx}(T)$ curves are shifted vertically by 0.8 kΩ.

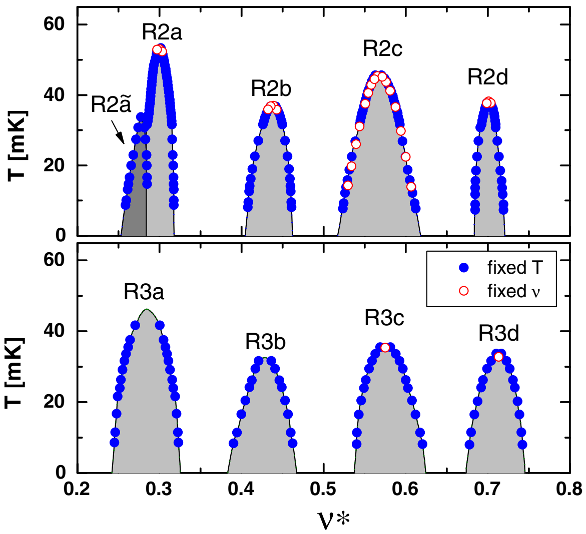

FIG. 3 (color online). The phase boundaries of the eight RIQHS in the SLL in the ν^*-T plane. The RIQHSs are stable within the shaded areas. Below 33 mK the $R2a$ state has a split-off state labeled $R2\tilde{a}$.

factor of the lower (upper) spin level. As described earlier, at a given ν the RIQHSs develop below the temperature of the peak in the $R_{xx}(T)|_{\nu^*=\text{fixed}}$ curve. Such peaks are shown in Fig. 2 for $\nu^* \approx \nu_c^*$, but similar peaks are also present for nearby filling factors (not shown). Open symbols in Fig. 3 are the temperatures of the peak as plotted against ν^*. Similarly, the RIQHSs develop between the spikes of the $R_{xx}(\nu)|_{T=\text{fixed}}$ curves, such as the ones shown in Fig. 1(b). The filling factors ν^* of the spikes for each RIQHS measured at a given temperature are marked with closed symbols in Fig. 3. The excellent overlap of the two data sets in Fig. 3 shows that the two definitions used above self-consistently define the stability boundary of each RIQHS. The shaded areas within each boundary of Fig. 3 represent the RIQHSs. FQHSs can develop only outside these shaded areas. The locations ν^*_{high} and ν^*_{low} of the spikes of the $R_{xx}(\nu)|_{T=\text{fixed}}$ curve measured at the lowest $T = 6.9$ mK of our experiment are listed in Table I.

We note that the $R2a$ state is different from the rest of the RIQHSs as it splits into two RIQHSs with a decreasing temperature. Such a split is signaled by an R_{xy} deviating from $h/2e^2$ as well as a nonzero R_{xx} in the vicinity of $\nu = 2 + 2/7$ and it has already been reported in Ref. [2]. The split-off RIQHS is marked as $R2\tilde{a}$ and with a darker shade in Fig. 3. We note that our data are similar to that in Refs. [1] in that the Ria, $i = 2, 3$ is the most stable state. Other studies find the $R2c$ state to be the most stable of RIQHSs [4–6,25,30–35].

Each stability boundary shown in Fig. 3 can be fitted close to their maxima with a parabolic form $T_c(\nu^*) = T_c(\nu_c^*) - \beta(\nu^* - \nu_c^*)^2$. The obtained parameters are listed in Table I. T_c obtained from the fit is within 1 mK from the peak temperature obtained from Fig. 2. The centers ν_c^* of the RIQHSs in the upper spin branch are in excellent agreement with the earlier reported values [1]. Those of the upper spin branch, however, have not yet been documented and they differ significantly from those of the lower spin branch. Indeed, $\nu^*_{c,R2\alpha} \neq \nu^*_{c,R3\alpha}$ for $\alpha = a, b, c,$ or d, the difference being the largest for the states a and d. Such a difference is not expected from the theory [23,24] and we think it is due to the interaction of the electrons in the topmost Landau level with those in the filled lower levels. Furthermore, we establish that the centers ν_c^* of RIQHSs in both spin branches obey particle-hole symmetry, as assumed by the theory [23,24]. In short $\nu^*_{c,Ria} = 1 - \nu^*_{c,Rid}$ and $\nu^*_{c,Rib} = 1 - \nu^*_{c,Ric}$ for $i = 2$, 3, relations which hold within our measurement error for the filling factor of ± 0.003.

In contrast to the centers of RIQHSs, other parameters of RIQHSs from Table I do not obey particle-hole symmetry. These parameters are the maximum onset temperatures $T_c(\nu_c^*)$, the fit parameter β describing the curvature of the stability diagrams near $T_c(\nu_c^*)$, and the widths $\Delta\nu = \nu^*_{\text{high}} - \nu^*_{\text{low}}$ of the stability regions of the RIQHSs at $T = 6.9$ mK. Indeed, particle-hole symmetry within a spin branch would imply a scaling of T_c with the Coulomb energy E_C and, therefore, a monotonically decreasing $T_c(\nu_c^*)$ with an increasing ν_c^*. Here $E_C = e^2/\epsilon l_B$ and $l_B = \sqrt{\hbar/eB}$ is the magnetic length. Data from Table. I, however, clearly show that contrary to this expectation $T_c(\nu_c^* = 0.568) > T_c(\nu_c^* = 0.438)$ [31]. We thus find that the particle-hole symmetry within one spin branch assumed in current theories [23,24] is violated. The nonmonotonic dependence of T_c on ν_c^* is, furthermore, at odds with the sequence of the one- and two-electron bubbles suggested [23,24]. These findings are puzzling and they show that there is still much left to be understood about the RIQHSs. Possible causes include Landau level mixing, disorder, or finite thickness effects. The origin of the broken symmetry described above is most likely related to and, therefore, it will influence the understanding of a

TABLE I. Parameters extracted from the ν^*-T diagram. T_c and β are in units of mK.

	$R2a$	$R2b$	$R2c$	$R2d$	$R3a$	$R3b$	$R3c$	$R3d$
ν_c^*	0.300	0.438	0.568	0.701	0.284	0.429	0.576	0.712
$T_c(\nu_c^*)$	53.0	37.1	45.8	38.0	46.3	32.3	36.1	33.8
$\beta \times 10^{-4}$	10	3.9	2.4	8.5	2.1	2.0	1.6	2.3
ν^*_{high}	0.317	0.461	0.613	0.719	0.324	0.463	0.621	0.742
ν^*_{low}	0.258	0.407	0.523	0.684	0.245	0.388	0.540	0.677

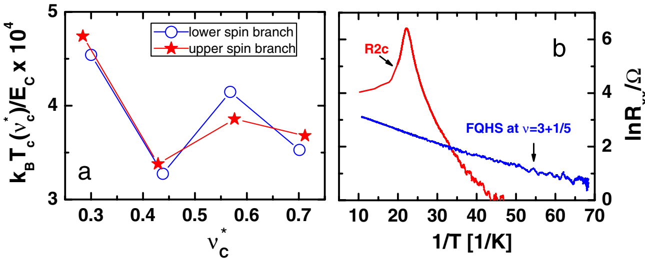

FIG. 4 (color online). The variation with the filling factor ν_c^* of the reduced onset temperatures $T_c(\nu_c^*)/E_C$ at the center of the RIQHSs in the SLL [panel (a)]. Lines are guides to the eye. Arrhenius plot for $R2c$ and for the $\nu = 3 + 1/5$ FQHS [panel (b)].

similar symmetry breaking of the Pfaffian and anti-Pfaffian construction for the $\nu = 5/2$ FQHS [36–43].

The onset temperatures $T_c(\nu_c^*)$ in the higher spin branch are consistently smaller than those in the lower spin branch. We notice, however, a startlingly similar dependence within each spin branch. A particularly revealing plot is that of the reduced onset temperatures $T_c(\nu_c^*)/E_C$ against the filling factor ν_c^*. As shown in Fig. 4(a), there is a surprisingly good collapse of $T_c(\nu_c^*)/E_C$ for the different spin branches. This collapse shows that Coulomb interactions play a central role in the formation of the RIQHSs in the SLL and provides direct evidence that these states reflect collective behavior of the electrons rather than single particle localization. The lack of collapse of $T_c(\nu_c^*)/\hbar\omega_C$ (not shown) means that $T_c(\nu_c^*)$ does not scale with the cyclotron energy $\hbar\omega_C$.

In a recent study an activated dependence of $R_{xx}(T)$ is found for the $R2c$ state [34]. In our sample we find a significant deviation from such a dependence and, as a consequence, the definition of an activation energy is no longer possible. Figure 4(b) shows such a plot for the $R2c$ state, together with the activated resistance of a suitably chosen FQHS measured in order to rule out thermometry artifacts. Our data suggest that nonactivated behavior might be an inherent property of the RIQHSs. The peak in the $R_{xx}(T)|_{\nu=\text{fixed}}$ curves could be due to interpenetrating RIQHS, a collective low T insulator and the high T classical electron fluid. In such an interpretation the nonactivated behavior seen in Fig. 4(b) is a consequence of the coexistence of these two phases.

The work at Purdue was supported by the DOE Grant No. DE-SC0006671. L. N. P. and K. W. W. were supported by the Princeton NSF-MRSEC and the Moore Foundation. We benefited from discussions with G. Giuliani and S. Das Sarma.

*gcsathy@purdue.edu

[1] J. P. Eisenstein et al., Phys. Rev. Lett. **88**, 076801 (2002).
[2] J. S. Xia et al., Phys. Rev. Lett. **93**, 176809 (2004).
[3] A. Kumar et al., Phys. Rev. Lett. **105**, 246808 (2010).
[4] I. P. Radu et al., Science **320**, 899 (2008).
[5] M. Dolev et al., Phys. Rev. Lett. **107**, 036805 (2011).
[6] N. Samkharadze et al., Phys. Rev. B **84**, R121305 (2011).
[7] M. Dolev et al., Nature (London) **452**, 829 (2008).
[8] R. L. Willett, L. N. Pfeiffer, and K. W. West, Phys. Rev. B **82**, 205301 (2010).
[9] V. Venkatachalam et al., Nature (London) **469**, 185 (2011).
[10] R. B. Laughlin, Phys. Rev. Lett. **50**, 1395 (1983).
[11] J. K. Jain, Phys. Rev. B **40**, 8079 (1989).
[12] G. Moore and N. Read, Nucl. Phys. **B360**, 362 (1991).
[13] M. Greiter, X. G. Wen, and F. Wilczek, Phys. Rev. Lett. **66**, 3205 (1991).
[14] C. Nayak and F. Wilczek, Nucl. Phys. **B479**, 529 (1996).
[15] P. Bonderson, V. Gurarie, and C. Nayak, Phys. Rev. B **83**, 075303 (2011).
[16] K. Park, V. Melik-Alaverdian, N. E. Bonesteel, and J. K. Jain, Phys. Rev. B **58**, R10167 (1998).
[17] N. Read and E. Rezayi, Phys. Rev. B **59**, 8084 (1999).
[18] V. Scarola, K. Park, and J. K. Jain, Nature (London) **406**, 863 (2000).
[19] E. H. Rezayi and F. D. M. Haldane, Phys. Rev. Lett. **84**, 4685 (2000).
[20] N. Read and D. Green, Phys. Rev. B **61**, 10 267 (2000).
[21] H. Lu, S. Das Sarma, and K. Park, Phys. Rev. B **82**, R201303 (2010).
[22] S. Das Sarma, M. Freedman, and C. Nayak, Phys. Rev. Lett. **94**, 166802 (2005).
[23] N. Shibata and D. Yoshioka, J. Phys. Soc. Jpn. **72**, 664 (2003).
[24] M. O. Goerbig, P. Lederer, and C. M. Smith, Phys. Rev. B **68**, 241302 (2003); **69**, 115327 (2004).
[25] R. M. Lewis et al., Phys. Rev. B **71**, R081301 (2005).
[26] N. Samkharadze et al., Rev. Sci. Instrum. **82**, 053902 (2011).
[27] A. A. Koulakov, M. M. Fogler, and B. I. Shklovskii, Phys. Rev. Lett. **76**, 499 (1996).
[28] M. P. Lilly et al., Phys. Rev. Lett. **82**, 394 (1999).
[29] R. R. Du et al., Solid State Commun. **109**, 389 (1999).
[30] G. A. Csáthy et al., Phys. Rev. Lett. **94**, 146801 (2005).
[31] W. Pan et al., Phys. Rev. B **77**, 075307 (2008).
[32] C. R. Dean et al., Phys. Rev. Lett. **100**, 146803 (2008).
[33] V. Umansky et al., J. Cryst. Growth **311**, 1658 (2009).
[34] J. Nuebler et al., Phys. Rev. B **81**, 035316 (2010).
[35] J. Xia et al., Phys. Rev. Lett. **105**, 176807 (2010).
[36] M. Levin, B. I. Halperin, and B. Rosenow, Phys. Rev. Lett. **99**, 236806 (2007).
[37] S. S. Lee et al., Phys. Rev. Lett. **99**, 236807 (2007).
[38] X. Wan et al., Phys. Rev. B **77**, 165316 (2008).
[39] H. Wang, D. N. Sheng, and F. D. M. Haldane, Phys. Rev. B **80**, 241311 (2009).
[40] W. Bishara and C. Nayak, Phys. Rev. B **80**, 121302 (2009).
[41] M. R. Peterson, K. Park, and S. Das Sarma, Phys. Rev. Lett. **101**, 156803 (2008).
[42] A. Wójs, C. Toke, and J. K. Jain, Phys. Rev. Lett. **105**, 096802 (2010).
[43] E. H. Rezayi and S. H. Simon, Phys. Rev. Lett. **106**, 116801 (2011).