The zero-energy state in graphene in a high magnetic field

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The fate of the charge-neutral Dirac point in graphene in a high magnetic field \( H \) has been investigated at low temperatures \( (T \sim 0.3 \, \text{K}) \). In samples with small gate-voltage offset \( V_0 \), the resistance \( R_0 \) at the Dirac point diverges steeply with \( H \), signalling a crossover to an insulating state in high field. The approach to the insulating state is highly unusual. Despite the steep divergence in \( R_0 \), the profile of \( R_0 \) vs. \( T \) in fixed \( H \) saturates to a \( T \)-independent value below 2 K, consistent with gapless charge-carrying excitations.

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The discovery of the quantum Hall effect (QHE) in monolayer graphene crystals provides a new system for investigating relativistic Dirac-like excitations in solids [1, 2, 3, 4, 5, 6]. In a magnetic field \( H \), the system forms Landau Levels (indexed by \( n \)) that are 4-fold degenerate. The Hall conductivity \( \sigma_{xy} \) is accurately quantized as the chemical potential \( \mu \) is changed from the hole part to electron part of the Dirac spectrum. Considerable attention has focussed on the \( n = 0 \) Landau Level (LL), especially on the nature of the electronic state at the charge-neutral point \( (\mu = 0) \) in an intense magnetic field \( H \). Several groups [7, 8, 9, 10, 11, 12] have predicted that a large field stabilizes the quantum Hall ferromagnetic (QHF) state, in which the pseudospins describing the valley degree of freedom become ferromagnetically ordered. In a second group of theories [13, 14], interaction causes an effective degree of freedom to become ferromagnetically ordered. In a magnetic field \( H \), the Hall conductivity \( \sigma_{xy} \) rises steeply from 4 kΩ to 190 kΩ. The curve of the conductivity \( \sigma_{xy} \) plotted vs. \( H \) (Panel b) displays plateaus given by [12]

\[
\sigma_{xy} = \frac{\nu e^2}{h} \left( n + \frac{1}{2} \right),
\]

where \( n \) indexes the 4-fold degenerate LL and \( \nu \) indexes individual sublevels. In K7, the ‘zero’ plateau \( \sigma_{xy} \approx 0 \) at \( V_g = 0 \) is already visible at \( H = 8 \, \text{T} \).

Narrowing our focus to the \( n = 0 \) LL, we examine \( R_{xx} \) in the \( n = 0 \) LL as a function of \( T \), with \( H \) fixed at 14 T (Fig. 2a). We see that, from 40 to 0.3 K, \( R_0 \) rises steeply from 4 kΩ to 190 kΩ. The curve of the conductivity \( \sigma_{xx} \) plotted vs. \( V_g \) reveals a 2-peak structure that implies splitting of the 4-fold degeneracy by a gap \( \Delta \) (Fig. 2b). At 100 K, the two peaks are already resolved. With decreasing \( T \), the minimum at \( V_g = 0 \) initially deepens rapidly, but saturates below \( \sim 2 \, \text{K} \). The Hall conductivity \( \sigma_{xy} \) at 0.3 K (thin curve) displays a well-defined plateau on which \( \sigma_{xy} \approx 0 \). Because the next plateau is at \( \sigma_{xy} = 2(\gamma^2/h) \), we infer that each of the peaks in \( R_{xx} \) is comprised of 2 unresolved sublevels. The opening of the gap causes \( \sigma_{xx} \) (at \( V_g = 0 \)) to fall rapidly with decreasing \( T \), until saturation occurs below 2 K.

The behavior of \( R_0 \) described here is qualitatively different from that in, for e.g., Ref. [15]. To understand the difference, we have examined several samples (Table I). As mentioned, the offset gate \( V_0 \) is a crucial parameter. Figure 2 compares the curves of \( R_{xx} \) (\( n = 0 \) LL) in the

| Sample | \( V_0 \) (V) | \( R_0(14) \) (kΩ) | \( \mu_e(1/T) \) |
|--------|-------------|----------------|-------------|
| K5     | 3           | 80             | 0.3         |
| K7     | 1           | 190            | 1.3         |
| K8     | 12          | 15             | 0.6         |
| K18    | 20          | 7.5            | 0.9         |
| K22    | -0.6        | >280           | 2.5         |
| K29    | 22.5        | 7              | 0.2         |

TABLE I: Sample parameters. \( V_0 \) is the gate voltage needed to bring \( \mu \) to 0. \( R_0(14) \) is \( R_0 \) measured at \( H = 14 \, \text{T} \) and \( T = 0.3 \, \text{K} \). \( \mu_e \) is the electron mobility at \( H = 0 \).
FIG. 1: (color online) The resistance $R_{xx}$ (Panel a) and Hall conductivity $\sigma_{xy}$ (b) in Sample K7 versus (shifted) gate voltage $V'_g = V_g - V_0$ at 0.3 K with $H$ fixed at 8, 11 and 14 T. Peaks of $R_{xx}$ at finite $V'_g$ correspond to the filling of the $n=1$ and $n=2$ LLs. At $V'_g = 0$, the peak in $R_{xx}$ grows to 190 kΩ at 14 T. The inset shows in false color a graphene crystal (dark red) with Au leads deposited (yellow regions). The bar indicates 5 μm. Panel b shows the quantization of $\sigma_{xy}$ at the values $(4e^2/h)(n + \frac{1}{2})$. At 0.3 K, $\sigma_{xy} = 0$ in a 2-Volt interval around $V'_g = 0$.

samples K5, K7, K8 and K29, all measured at 14 T at $T = 0.3$ K. For each sample, we have plotted $R_{xx}$ vs. the unshifted gate voltage $V_g$, so its peak automatically locates $V_0$. It is clear that K7 ($V_0 = 1$ V) has the highest peak, followed by K5 ($V_0 = 3$ V), whereas K8 ($V_0 = 12$ V) and K29 (22.5 V) have peaks that are severely suppressed. Further insight into this pattern of suppression is given below when we examine the $H$ dependence of $R_0$ at low $T$.

Sample self-heating may also obscure the divergence. We find that, below 1 K, self-heating becomes serious when the dissipation exceeds ~2 pW. The measurements of $R_{xx}$ vs. $V_g$ were repeated at 3 currents ($I = 0.6$, 2 and 15 nA) at $T = 0.3$ K. The results at $I = 0.6$ and 2 nA are virtually identical. However, the curve at 15 nA is 30% smaller near $V'_g = 0$, consistent with heating. Hence, we have kept $I$ at 2 nA to eliminate self-heating as a problem. Heating at the contacts is negligible because of the small contact resistances (~1 kΩ) relative to $R_0$.

Hereafter, we focus on $R_0$, or equivalently, the Dirac-point conductivity $\sigma_{xx}^0 = L/wR_0$ ($L$ and $w$ are the length and width). Curves of the conductivity versus $\log_{10} T$ are shown in Fig. 3b at selected fields. In low fields ($H < 9$ T), the $T$ dependence of $\sigma_{xx}^0$ is quite mild. As $H$ is increased to 14 T, the opening of the gap $\Delta$ (between the $n=0$ sublevels) causes the conductance to decrease sharply below 40 K. However, instead of falling to 0, $\sigma_{xx}^0$ saturates below 2 K to a $T$-independent residual value $\sigma_{res}$, as anticipated in the discussion of Fig. 2b. The existence of this residual $\sigma_{res}$, which is highly sensitive to $H$, is one of our key findings.

The field dependence of $\sigma_{res}$ is best viewed as a divergent $R_0$. Figure 3b shows the rising profile of $R_0$ vs. $H$.
in sample K7 at selected temperatures. The divergent form at the lowest $T$ (0.3 K) strongly suggests that the system is rapidly approaching a field-induced crossover (or transition) to an insulating state.

In light of the importance of $V_0$, it is instructive to see how the profile of $R_0$ vs. $H$ varies between samples. Figure 3: (color online) The $T$ dependence of $\sigma_{xx}^0$ in K7 (= $L/\mu R_0$) and the $H$ dependence of $R_0$ at low temperature. Panel a shows curves of $\sigma_{xx}^0$ vs. $\log_{10} T$ with $H$ fixed at 6-14 T. For $H > 8$ T, the gap $\Delta$ causes $\sigma_{xx}^0$ to decrease markedly until saturation at the residual value $\sigma_{xx}$ occurs below 2 K. Panel (b) displays the steep increase in $R_0$ vs. $H$ in K7 at selected $T$. At 0.3 K, $R_0$ appears to diverge at a field near 18 T (see Fig. 4b). Panel (c) compares the $R_0(H)$ profiles in Samples K7, K18 and K22. In Sample K18 ($V_0 = 20$ V), the divergence in $R_0$ becomes apparent only above 14 T, whereas in K22 ($V_0 = -0.6$ V) $R_0$ starts to diverge at fields lower than in K7. In K7, we have plotted $R_0$ values measured by sweeping $V_0$ at fixed $H$ (solid symbols) with $R_0$ measured by sweeping $H$ with $V'_0$ fixed at 0 (solid curve), to show consistency.

FIG. 3: (color online) The $T$ dependence of $\sigma_{xx}^0$ in K7 (= $L/\mu R_0$) and the $H$ dependence of $R_0$ at low temperature. Panel a shows curves of $\sigma_{xx}^0$ vs. $\log_{10} T$ with $H$ fixed at 6-14 T. For $H > 8$ T, the gap $\Delta$ causes $\sigma_{xx}^0$ to decrease markedly until saturation at the residual value $\sigma_{xx}$ occurs below 2 K. Panel (b) displays the steep increase in $R_0$ vs. $H$ in K7 at selected $T$. At 0.3 K, $R_0$ appears to diverge at a field near 18 T (see Fig. 4b). Panel (c) compares the $R_0(H)$ profiles in Samples K7, K18 and K22. In Sample K18 ($V_0 = 20$ V), the divergence in $R_0$ becomes apparent only above 14 T, whereas in K22 ($V_0 = -0.6$ V) $R_0$ starts to diverge at fields lower than in K7. In K7, we have plotted $R_0$ values measured by sweeping $V_0$ at fixed $H$ (solid symbols) with $R_0$ measured by sweeping $H$ with $V'_0$ fixed at 0 (solid curve), to show consistency.

FIG. 4: (color online) (Panel a) The contour plot of $R_0(T, H)$ (K7) in the $T$-$H$ plane (vertical bar shows values of $R_0$). The contour lines emphasize the unusual approach to the insulating state. At low $T$, $R_0$ is unchanged on a horizontal path ($H$ held fixed), but it rises rapidly on a vertical path (increasing $H$ at fixed $T$). Panel (b) displays log $R_0$ vs. $1/\sqrt{T - \hbar}$ in K7 at $T = 0.3$ K, with $\hbar = H/H_0$, where $H_0 = 18$ T. The linear segment at large $R_0$ shows that the divergence is consistent with $R_0(h) \sim \exp[2b/\sqrt{T - \hbar}]$ with $b \sim 0.7$. In Panel (c), the plot of $(d\ln R_0/dH)^{2/3}$ vs. $H$ shows a high-field linear segment that extrapolates to zero at $\sim H_0$ (18 T).

The variation of $R_0(T, H)$ in K7 is conveniently represented in a contour plot in the $T$-$H$ plane (Fig. 4).
Below \( \sim 2 \) K, the contour lines are horizontal, which implies that \( R_0 \) is unchanged if the sample is cooled in fixed \( H \). This provides evidence that \( \sigma_{res} \) involves gapless excitations. However, if \( T \) is fixed, \( R_0 \) rises steeply with \( H \), implying proximity to the insulating state (deep-red region). When a system approaches the insulating state, its resistivity generally diverges as \( T \to 0 \), as a result of either strong localization (variable-range hopping) or the opening of a mobility gap (weak localization is not relevant here because of the intense \( H \)). In both cases, decreasing \( T \) reduces the conductance because the itinerant states are severely depopulated. Hence, the pattern in Fig. 4a is most unusual. The gaplessness of ant states are severely depopulated. Hence, the pattern in

\[
\begin{align*}
\text{These considerations lead us to quantify the divergence} & \text{ of either strong localization (variable-range hopping) or} \\
& \text{the opening of a mobility gap (weak localization is not} \\
& \text{relevant here because of the intense \( H \)). In both cases,} \\
& \text{decreasing \( T \) reduces the conductance because the itinerant} \\
& \text{states are severely depopulated. Hence, the pattern} \\
& \text{in Fig. 4a is most unusual. The gaplessness of \( \sigma_{res} \) sug-} \\
& \text{gests that, below 2 K, these excitations are protected} \\
& \text{from the effects of changing \( T \). Paradoxically, they are} \\
& \text{not protected from an increasing \( H \), which reduces the} \\
& \text{current carried at an exponential rate.} \\
\end{align*}
\]

In the theory in Refs. [15, 18], the current at the Dirac point is carried by a pair of edge states. The change in \( R_{xx} \) vs. \( H \) is interpreted as an increased scattering rate as the edge states are pushed closer to the edge.

In samples with small \( V_0 \), however, the steep increase in \( R_0 \) to \( \sim 200 \) kΩ (Fig. 5b) suggests a different regime in which Coulomb exchange may be the dominant energy scale. From Figs. 3b and 4a, we infer that \( R_0 \) appears to be diverging towards an insulating state in high fields. These considerations lead us to quantify the divergence in \( R_0 \).

We find that \( R_0 \) fits very well to the form \( R_0 \sim \xi(h)^2 \), where the correlation length \( \xi \) has the Kosterlitz-Thouless (KT) form

\[
\xi_{KT} \sim \exp(b/\sqrt{1-h}), \quad (h = H/H_0),
\]

with \( H \) replacing \( T \). Plotting \( \ln R_0 \) vs. \( \sqrt{1-h} \), we find that the high-field portion becomes linear (Fig. 4b) when \( H_0 \) is adjusted to be 17-18 T. From the slope, we find that the parameter \( b \) has the value \( \sim 0.7 \), in agreement with simulations of the KT transition. For self consistency, we may also let the data determine \( H_0 \) by linear extrapolation. By Eq. 2, we have \( d \ln R_0/dH \sim (H_0 - H)^{-3/2} \). Hence a plot of \( (d \ln R_0/dH)^{-2/3} \) vs. \( H \) should cross the \( H \)-axis at \( H_0 \). Indeed, this quantity, plotted in Fig. 4c, becomes linear at large \( H \) and extrapolates to zero at \( \sim 18 \) T in agreement with (b).

Although consistency with Eq. 2 alone does not prove that a KT transition occurs at \( H_0 \), the fit does reveal the striking exponential character of the divergence in \( R_0 \). We adopt as a working hypothesis that this reflects the approach to a KT transition. The ordered state at large \( H \) is destroyed by the spontaneous appearance of defects which increase exponentially in density at \( H_0 \). The steep fall of \( R_0 \) below \( H_0 \) may reflect the current carried by these defects.

In both the QHF state [7, 8, 10, 11] and the exciton-gap state [13, 14], the 4-fold degeneracy in the \( n = 0 \) LL is completely lifted in large \( H \) to produce an insulator. In the QHF, the starting symmetry is \( SU(4) \) if we treat the spin and pseudospin degrees on equal footing [8]. The reduction of the \( SU(4) \) to lower symmetries by the effects of Zeeman energy, disorder or lattice discretization is discussed in Refs. [8, 10, 11, 12]. According to Ref. [11] random disorder may force the pseudospins into the plane. If the ordered state indeed has \( U(1) \) symmetry, the XY order is susceptible to a KT transition [7, 11]. It would be very interesting to relate the fit of \( R_0 \) to Eq. 2 with charged, topological excitations envisioned in the KT transition. Measurements are in progress at higher \( H \) to provide further evidence for a transition to the insulating state and to clarify its nature.

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