The magnetoplasmon spectrum of Landau level transitions in hexagonal boron nitride-encapsulated graphene is explored via infrared transmission magnetospectroscopy, as a function of the filling factor at fixed magnetic field. As the lowest Landau level occupancy is increased from half-filling, a non-monotonic progression of multiple cyclotron resonance peaks is observed, with a single peak evolving into four peaks and back to two, all with linewidths of order 0.5 meV. This provides a novel window on the interplay of electron interactions with broken spin and valley symmetries in the quantum Hall regime. Analysis of the peak energies shows an indirect enhancement of spin gaps below the Fermi energy, a Dirac mass at half-filling that is nearly 50% larger than when the lowest Landau level is completely full, and a small but clear particle-hole asymmetry. We suggest a key role is played by the boron nitride in enabling interaction-enhanced broken symmetries to be observed in graphene cyclotron resonance.
The cyclotron resonance (CR) of a translationally invariant system with a parabolic dispersion is insensitive to electron-electron interactions, a result known as Kohn’s theorem [1–3]. This has limited the utility of CR in the study of many-particle phenomena such as found in the fractional quantum Hall regime [4, 5]. The linear dispersion of graphene, on the other hand, is widely held to negate Kohn’s theorem [6–16]. In fact, a limited version of the theorem is predicted to survive for transitions between the evenly-spaced \( n = 0 \) and \( \pm 1 \) Landau levels (LL) [7–9, 17, 18]. Nonetheless, in the infrared magnetospectroscopy of graphite-gated, hexagonal boron nitride-encapsulated graphene, we find a remarkable progression of CR lineshapes for transitions between these lowest LLs, with a single peak appearing at half-filling of the lowest LL that splits into four peaks and then just two as this level becomes completely occupied. This provides an unprecedented spectroscopic view of the evolution of the magnetoplasmon spectrum in the regime of broken spin and valley symmetries. We propose that contact between graphene and the encapsulating layers of hexagonal boron nitride (hBN) breaks translation symmetry and serves to lift the remnants of Kohn’s theorem, enabling direct contributions to the CR energy by interaction effects. While this work specifically addresses physics in graphene, the approach is applicable in principle to any system with a linear dispersion and so may find utility in understanding the competing roles of interactions and symmetry breaking in Dirac, Weyl, or strongly correlated materials [19].

In a strong magnetic field, graphene develops a set of four-fold degenerate LLs (two each for electron spin and the \( K \) and \( K' \) valleys) with energies given by \( E_n = s_n \hbar \omega_c \sqrt{|n|} \), where \( \omega_c = \sqrt{2} v l_B \) is the cyclotron frequency, \( v \sim 10^6 \text{ m/s} \) is the band velocity, \( l_B = \sqrt{\hbar/eB} \) the magnetic length, \( s_n = \text{sign}(n) \), and \( n = 0, \pm 1, \pm 2... \) is the orbital index [20, 21]. If the sublattice symmetry of graphene is broken, as is common for hBN-encapsulated devices, a gap (or Dirac mass) \( M \) appears which shifts the \( |n| > 0 \) levels according to \( E_n = s_n \hbar \omega_c \sqrt{|n| + \mu^2} \), where \( \mu = M/\hbar \omega_c \), with the valley-polarized \( n = 0 \) level split by \( E_{0,K}(E_{0,K'}) = +(-)M \) [22]. Cyclotron resonance corresponds to the absorption of a photon which elevates an electron from one LL to another that differs by one unit of the orbital quantum number and leaves a hole behind; the selection rule for this transition is \( \Delta |n| = \pm 1 \). Since the conduction and valence bands share \( p \)-type symmetry, CR includes both inter- and intra-band transitions with energies nominally given by the LL separation

\[
\Delta E_{m,n} = \hbar \omega_c \left( s_m \sqrt{|m| + \mu^2} - s_n \sqrt{|n| + \mu^2} \right).
\]

A key consequence of electron interactions in graphene is to renormalize the band velocity \( v \to \tilde{v} \) [6, 23–25]; in a magnetic field \( \tilde{v} \) decreases weakly with increasing \( B \), or at fixed field it is roughly constant when the lowest LL is partially filled and declines sharply as higher LLs become occupied.
Many-particle contributions have been seen in CR and magneto-Raman experiments \cite{6,11-15}, and generally good agreement with theory has been found \cite{8,9,17,18,26}.

Strictly speaking, in a translationally-invariant system with a parabolic dispersion, the energy of the long-wavelength $q = 0$ magnetoplasmon mode of the CR excitation is pinned to the LL separation $\omega_c = eB/m^*$, as required by Kohn’s theorem \cite{3,27}. However, in practice band non-parabolicity \cite{28,30}, disorder \cite{31,32}, multiple effective masses \cite{33,34}, intense irradiation \cite{35}, or intentional breaking of translation invariance \cite{36,38} all allow for many-particle contributions to the CR energy. Yet these end-runs around Kohn’s theorem inevitably include effects beyond interactions \cite{29,39}, due to whichever mechanism is used to evade it. In contrast, the linear Dirac dispersion of graphene can be considered an extreme case of non-parabolicity where the carriers have no center of mass, and Kohn’s theorem ought no longer apply \cite{7,10}. In fact a limited version of the theory does survive, although just for transitions between the evenly-spaced $n = 0$ and $\pm 1$ LLs, for which the direct Coulomb and exchange interactions cancel and leave just the overall renormalization of the band velocity \cite{7,9,17,18}. For all other transitions, it should be possible to directly observe interaction contributions to the magnetoplasmon spectrum, promising that minimally-disordered graphene devices can become a laboratory for exploring the excited states of interacting many-particle systems. Nevertheless, here we find that transitions to or from the $n = 0$ LL show clear evidence of symmetry-breaking and many-particle contributions, strongly suggesting even this remnant of Kohn’s theorem is broken.

The sample used in this study is an 820 $\mu$m$^2$ sheet of monolayer graphene sandwiched between $\approx 40$-nm-thick flakes of hexagonal boron nitride, assembled using a dry-stacking technique \cite{40} and placed on a $\approx 4$-nm-thick flake of single-crystal graphite which lies on a lightly-doped, oxidized Si wafer. Electrical contacts to the edge of the graphene were made using 3/60-nm-thick films of Cr/Au, defined by standard electron beam lithography techniques. Aluminum foil with a $\approx 90$-$\mu$m diameter aperture is placed immediately on top of the sample (visible inset to Fig. 1(c)) to restrict the infrared light to the region immediately surrounding the sample. All spectroscopic data in this work were acquired at a temperature of 300 mK and a fixed magnetic field of 8.0 T, using a Bruker V80v broadband Fourier-transform infrared spectrometer with instrumental resolution set to 0.5 meV (except for exploratory traces at other resolutions, samples of which are included in the Supplemental Material). Unpolarized blackbody light from the spectrometer was coupled through a KBr window into a Blue Fors cryogen free dilution refrigerator with a 14 T solenoid, focused to and defocused from the sample using custom parabolic optics, and funneled via a compound parabolic collector (Winston cone) to a composite Si bolometer held at 4 K for detection. Traces
were acquired at specific LL filling factors \( \nu = 2\pi n_s l_B^2 \) (where \( n_s \) is the carrier sheet density), and normalized by traces acquired at a higher filling of \( \nu = +18 \). Absorption features common to both traces that do not depend on the gate voltage—including from the substrate or elsewhere in the beam path—divide to unity, while absorptions in the target trace appear as dips in the normalized transmission. Each trace at a given filling factor was averaged for approximately four hours. An example of the normalizing procedure and discussion of incidental graphite features is included in the Supplemental Material.

In graphene, several interband CR transitions \( T_i \) can be observed simultaneously at fixed filling factor, comprising nominally degenerate pairs of inter-LL excitations \( n = -i \to i - 1 \) and \( 1 - i \to i \) with energies given by Eq. 1. Figure 1(a) shows a color map of transitions \( T_1 \) through \( T_5 \) acquired as a function of \( \nu \), where the square-root dependence of the energies on the LL indices is immediately apparent. A schematic of the allowed transitions at \( \nu = 0 \) is drawn in Fig. 1(b), and a linecut at \( \nu = 0 \) in Fig. 1(c) highlights the sharp \( T_1 \) transition. The remarkably narrow resonances are a consequence of recent improvements in sample fabrication techniques, and are key to our observations. In Fig. 1(d) we show \( T_1 \) at \( \nu = 0 \) in three separate devices from the present and two prior works [13, 41], in which the half-width at half-max, \( \Gamma \), has clearly decreased. In fact, the lower two traces in Fig. 1(d) provide a comparison of two common gating methods: the middle trace has a distant, doped Si/SiO\(_2\) substrate on which the encapsulated monolayer rests [13, 42], while the lower trace which has a local graphite gate [43]. By chance these two devices have similar mobilities of 200,000 cm\(^2\)/Vs, but show differing ranges of the quantum scattering time \( \tau_q [44] \), likely due to screening of charged impurities in the SiO\(_2\) by the graphite. The CR lifetimes \( \tau_{CR} = \hbar/\Gamma \) in Fig. 1(e) are similarly improved, and in fact the \( \sim 600 \) fs quoted is a lower limit as even narrower lines with \( \tau_{CR} \approx 2.5 \) ps (\( \Gamma = 0.26 \) meV) are seen in this device at higher instrumental resolution—an order of magnitude greater than \( \tau_q [45] \). Perhaps not coincidentally, this is close to the transport time derived from the mobility [45]. As previously noted in the CR of AlGaN/GaN heterostructures [16], \( \tau_{CR} \) can be several times larger than \( \tau_q \) (which can be reduced by e.g. density variations across the sample) suggesting that \( \tau_{CR} \) is a better measure of the total carrier scattering time.

In Figure 2 we focus on the \( T_1 \) transition, which displays a marked non-monotonic evolution from a single peak at \( \nu = 0 \) into four peaks around \( \nu = +1 \), and then becomes two peaks for \( \nu \geq +2 \) which both fade away when \( \nu \to 6 \) and the \( n = +1 \) LL is completely filled; a sudden sharp rise in the lower energy peak above \( \nu = 5 \) signals the extinction of the resonance. Linecuts in Fig. 2(b) show details at half-integer \( \nu \). The features at \( \nu = 1 \) and 2 persist over a wide range of \( \nu \).
This is a real effect and not due, for instance, to small variations $\delta n_s$ in the carrier density across the sample: from the width of the Dirac peak in the zero-field resistance vs density, we estimate $\delta n_s \approx 2 \times 10^{10}$ cm$^{-2}$, or $\delta \nu \approx 0.1$ at 8 T, much smaller than the range of fillings over which the $\nu = 1$ and 2 features are seen [45]. At half-integer fillings below $\nu = 2$, only single broad resonances appear that nevertheless maintain the full spectral weight, suggesting all transitions are present if undifferentiated [45]. This could indicate the presence of dark magnetoexciton modes serving as additional scattering channels, but which do not contribute to optical transitions. Indeed there are up to 16 distinct transitions between the 0 and $\pm 1$ LLs [7], though only the four that conserve spin and valley are optically active.

The CR energy is given by the energy difference of participating levels, with the restriction that the spin and valley orientations are preserved. In the absence of interactions or symmetry-breaking mechanisms, the energies of the four component transitions of $T_1$ are degenerate and equal to the LL separation. Any gap, e.g. due to Zeeman splitting, that is the same in both LLs is invisible to CR. Thus the multi-peak structure in Fig. 2 particularly at $\nu = 1$, requires that both spin and valley symmetries are broken to a differing extent in the $n = 0$ and $\pm 1$ LLs. In the present device architecture, an obvious valley-symmetry-breaking mechanism results from contact of graphene with the encompassing hBN [47, 48], generating a mass $M$ that splits the $n = 0$ level and shifts the higher levels. Additionally, the spin splittings of the $n = 0$ and $n = \pm 1$ levels must also differ in energy, which is most readily achieved by invoking Coulomb interactions. But as noted above, in graphene the remnant of Kohn’s theorem allows only for a filling-independent renormalization of the velocity, $v \rightarrow \tilde{v}$, for $T_1$. Thus Eq. 1 should apply and even with finite $\mu$, the $T_1$ energy at $\nu = 0$ should be precisely the same as the higher energy peak when $T_1$ splits at $\nu = 2$ [17]. Yet in Fig. 2 it is clearly higher, by 1.3 meV. Thus even this surviving part of Kohn’s theorem must be broken. This cannot be due to non-parabolicity. Disorder is a potential culprit, although it is quite weak as seen by how narrow the CR linewidths are compared to the observed shifts and splittings. On the other hand, the interaction of graphene with the hBN substrate gives rise to a long wavelength moiré pattern that breaks translation symmetry of the graphene lattice [47, 49]. Thus we propose the hBN is responsible for lifting the final restrictions of Kohn’s theorem, enabling all interaction effects to contribute to the CR energies on an equal footing. This is a fortunate outcome, as interactions are relatively strongest for the $T_1$ transition.

To study the impact of Coulomb contributions, in Figure 2(c) we introduce schematics representing the simplest model of transitions between the $n = 0$ and $\pm 1$ LLs that aligns with the observed CR. These are drawn for $\nu = 0, +1$ and +2 with each of the four spin- and valley-resolved
levels shown explicitly. The level shifts and gap sizes are exaggerated, and represent single particle states enhanced by Coulomb interactions. At \( \nu = +2 \), all orbitals are filled or empty so interaction effects should be minimal. Assuming that all spin gaps are equal but allowing for the possibility of a weak valley splitting in \( n \neq 0 \) LLs, this schematic gives two doubly-degenerate resonances with energies split by

\[
\Delta E_{0,1}^{k,l} - \Delta E_{0,1}^{i,j} = \left[ E_1 + \frac{1}{2} (\Delta v_0 - \Delta v_1) \right] - \left[ E_1 + \frac{1}{2} (\Delta v_0 + \Delta v_1) \right] = \Delta v_{0,1}^{\nu=2}.
\]

Here \( \Delta E_{0,1}^{k,l} \) is the energy of the \( k \) and \( l \) lines in Fig. 2(c) and \( \Delta v_{0,1}^{\nu=2} \) is the difference of the \( n = 0 \) and 1 valley gaps. If the gaps varied in any other way, e.g. if the Zeeman splitting is enhanced in one level over another, this picture would return additional resonances not present in the data. To be fair, the \( \nu = 2 \) linecut in Fig. 2(b) shows what may be small extra peaks, and the higher of the two resonances is broadened with a suggestive shoulder on its high energy side; but for now we focus on just the brightest peaks persisting in the color map, and leave exploration of finer features to future higher quality devices. Fitting the two peaks at \( \nu = 2 \) with Lorentzians, we find \( \Delta v_{0,1}^{\nu=2} = 5.0(1) \) meV. Since any valley splitting of the \( n = 1 \) valley is likely to be small, this should be a good measure of the valley gap in the zeroth LL and yields a Dirac mass of \( M = 2.5 \) meV.

The single peak at \( \nu = 0 \) requires that all four transitions are degenerate, and by the schematic are equal to the LL separation plus half the difference of the valley gaps in the zeroth and \( \pm 1 \) LLs. In particular the \( n = 0 \) and \( \pm 1 \) Zeeman gaps must be equal, or else further CR lines would be seen. Actually, the \( \nu = 0 \) resonance is the broadest of the \( T_1 \) set, suggesting there may be unresolved lines, perhaps from a differential enhancement of the Zeeman energies. As noted, the energy of \( T_1(\nu = 0) \) is greater than the upper peak at \( \nu = +2 \), so the valley gap (or rather, the gap difference \( \Delta v_{0,1} \)) must also be larger at \( \nu = 0 \). We estimate its size by taking \( E_1 \) to be the average of the \( \nu = 2 \) peak energies, and writing \( \Delta v_{0,1}^{\nu=0} = 2(\Delta E_{meas} - E_1) = 7.3(5) \) meV, suggesting the Dirac mass at \( \nu = 0 \) is significantly enhanced over the \( \nu = 2 \) value.

For the four peaks near \( \nu = +1 \), each transition must sample a different combination of valley and spin gaps in order to avoid degeneracies. The gaps as sketched in Fig. 2 are not a unique selection, but show the simplest way to obtain four distinct energies. Support for this picture is gained by considering how the transitions labeled \( a \) through \( d \) at \( \nu = 0 \) evolve into transitions \( e \) through \( h \) at \( \nu = +1 \), to become \( i \) through \( l \) at \( \nu = +2 \). From \( \nu = 0 \) to \( +1 \), the Fermi energy moves up and fills half of the \( K' \) valley, extinguishing transition \( b \) and initiating \( f \), which is now the lowest energy peak. Meanwhile transition \( a \) becomes \( e \), the highest peak at \( \nu = 1 \), upon the enhancement of the spin gap \( \Delta z_{0,} \). This leaves transitions \( c \rightarrow g \) and \( d \rightarrow h \). The latter is the second highest
peak in Fig. 2(a), while g becomes the third highest. With concomitant gap enhancements, this naturally explains the evolution of the peak structure. Note that g and h, which were degenerate at \( \nu = 0 \) (when carrying labels c and d), must now differ at \( \nu = 2 \) due to the enhancement of a spin gap well below the Fermi energy.

The detailed evolution of this peak structure is likely to contain information relevant to the formation of stable magnetoplasmon modes. For instance in Fig. 2(a), the upper two peaks at \( \nu = +1 \) roughly extrapolate to the energy of the single \( \nu = 0 \) peak. With increasing \( \nu \) in the neighborhood of \( \nu = +1 \), these two begin to separate, and then disappear before the lower two peaks. As \( \nu \to 2 \), those lower peaks coalesce and then rapidly rise up to become the higher side of the \( \nu = 2 \) splitting, while the lower side of the \( \nu = 2 \) splitting appears without any precursor (note these lower peaks must have come from the highest and lowest \( \nu = 1 \) transitions, e and f). The two-peak splitting then remains until nearly the end as \( \nu \to 6 \), \( n = 0 \) is filled, and \( T_1 \) is extinguished. (We briefly note the presence of faint features near \( \nu = 1/2 \) and \( 3/2 \), possibly related to fractional fillings or collective excitations localized near disorder [50]).

Within this simple schematic model, we can derive the differences between spin and valley gaps in neighboring LLs. Each transition energy is taken to be the difference of the initial and final LL energies, plus or minus half of each relevant valley and Zeeman gap. For instance, the energy of transition e is \( \Delta E_e = E_{-1} + 1/2(\Delta v_{0,-1} + \Delta z_{0,-1,-}) \), where for transitions from \( n = -1 \to 0 \) the difference in valley gaps is \( \Delta v_{0,-1} = \Delta v_0 - \Delta v_{-1} \), and for spin gaps is \( \Delta z_{0,-1,-} = \Delta z_{0,-} - \Delta z_{-1,-} \). Here + (−) labels the \( K \) (\( K' \)) valley. While there are four peaks in the CR, there are six free parameters in the model if all relevant valley and spin gap differences are counted along with \( E_1 \), so some simplifying assumptions are needed to extract any quantitative values. The model is sensitive to the precise value of \( E_1 \), so we use the three differences between transition energies, \( \Delta E_{a,b} = \Delta E_a - \Delta E_b, \Delta E_{b,c}, \) and \( \Delta E_{c,d} \). The full procedure is described in the Supplemental Material and yields three valley and spin gap differences \( \Delta v = 5.0 \text{ meV}, \Delta z_+ = 2.1 \text{ meV}, \) and \( \Delta z_- = 4.3 \text{ meV} \). These are, respectively, the energy differences of the \( n = 0 \) and \( |n| = 1 \) valley gaps, and of the Zeeman gaps in the \( K(+) \) and \( K'(-) \) valleys of each LL. While the valley gap is nearly identical to that found at \( \nu = +2 \), the spin gaps are significantly larger than the bare Zeeman energy at this field, \( E_Z = 0.93 \text{ meV} \), indicating a clear role for electron interactions. This \( \Delta z_+ \) splitting is notable for occurring well below the Fermi energy when the system is in the \( \nu = 1 \) spin polarized state, reminiscent of indirect exchange splittings seen in GaAs quantum wells [51]. Meanwhile the size of \( \Delta z_- \) at the Fermi energy compares well to a transport gap at \( \nu = -1 \) of \( \sim 5 \) meV, at 9 T, found in Ref. [52]. Casting the energies as effective \( g \)-factors, the spin gap at the
Fermi level has $g_{z/-}^* = 4.7$, and the buried spin gap (in the K valley) has $g_{z/+}^* = 2.2$. We note a recent optical study of multilayer epitaxial graphene found multiple peaks in the CR attributed to broken spin and valley symmetries from which $g_{z}^* = 4.8$ was extracted, in close agreement [15].

In Fig. 3(a) we zoom out to show the full filling factor range of $T_1$. In either the single- or many-particle picture, a $\nu \rightarrow -\nu$ symmetry is expected, but on close inspection small yet non-negligible differences appear. For instance, the 5.0 meV magnitude of the $\nu = \pm 2$ splittings is virtually identical, but the hole-side peaks are a full 1.0 meV lower in energy. Moreover on closer examination of the $\nu = -1$ and +1 data in Fig. 3(b) and (c), the two lower energy transitions are both separated by 1.7 meV and exhibit a slow ramp up with increasing $|\nu|$, but the hole-side pair is lower in energy by 1.1 meV. Meanwhile, the two higher-energy peaks on the hole side nearly overlap, whereas they are still individually resolved around $\nu = +1$; and this higher-energy pair are closer together (or overlapping) on the left side of the graph but begin to separate toward the right as $\nu$ decreases (increases) in the $-1$ (+1) data, breaking a $\nu \rightarrow -\nu$ symmetry. Finally, the highest energy peak on the hole side is only 0.4 meV lower than the electron side. Overall, the observed symmetry breaking is a $\sim 0.8\%$ effect, too small to have been noticed in early broadband spectroscopic studies [9] but matching an asymmetry apparent in the data of Ref. [13] [41]. For completeness we note swept-field CR studies in graphene-on-oxide and multilayer epitaxial graphene report a far larger asymmetry [15] [53]. Such asymmetry is not anticipated in many-particle theories to date, although it can arise in tight-binding models of monolayer graphene due to next-nearest-neighbor hopping [54]. For instance, a field-dependent asymmetry between the $-n \rightarrow n - 1$ and $1 - n \rightarrow n$ transitions given by $E_{\text{asym}} = 3\sqrt{2}\hbar\omega_c t'a/\tau t_B \approx 0.56$ meV at 8 T was derived in Ref. [55], albeit in the high-$n$ limit (where $t$ ($t'$) is the nearest (next-nearest) neighbor hopping, and $a$ the C-C atom distance in graphene). This value is within a factor of two of the asymmetry seen here.

In Figure 4(a) and (b) the transition energies at $\nu = 0$ and 2 are plotted as a function of transition number $T_i$, parameterized as an effective velocity $v_{\text{eff}}(T_i) = \Delta E_{\text{meas}}(T_i)/\Delta E_{i,-1}(\nu = 10^6; \mu = 0)$. In both cases $v_{\text{eff}}$ rises from $T_1$ to $T_2$ and thereafter gradually decreases [13]. We analyze the data at $\nu = 2$ using two approaches, and predict $v_{\text{eff}}$ for the $\nu = 0$ data with the resulting fit parameters. First we fit the data with the gapped but non-interacting model using Eq. 1 and find the velocity $\tilde{v} = 1.116 \times 10^6$ m/s and the Dirac mass $M = 2.51$ meV. By design the splitting at $T_1$ is well-matched, but otherwise this approach fails badly for both the $\nu = 2$ and $\nu = 0$ data. We find much better results with the theory of Ref. [17] that accounts for many-particle contributions to CR in a single-mode approximation [56]. The fit parameters are
the velocity and mass, and include an overall Coulomb interaction $V_C = \sqrt{\frac{\pi}{2}} \frac{e^2}{(4\pi\epsilon l_B)} = 50$ meV set by choosing $\epsilon = 4$ for hBN [57]. This theory provides a good account of the variation of $v_{eff}$ and the $T_2$ splittings using $\tilde{v} = 1.105 \times 10^6$ m/s and $M = 2.76$ meV. However, no single set of parameters can simultaneously match the $T_1$ data at both $\nu = 0$ and $\nu = +2$. This is presumably due to the limited version of Kohn’s theorem active in the theory of Ref. [17], but not the data.

We have seen the Dirac masses extracted from CR differ depending on the filling factor: to sum up, the valley gap (extracted in discussion of the LL schematic of Fig. 2 not from the curve fits) is 7.3, 5.0, and 5.0 meV for $\nu = 0, 1,$ and 2, respectively. We note the hBN-induced gap in graphene is expected to be enhanced by interactions at half-filling [49, 58, 59], as we observe. Interestingly, these values are smaller than the 15 meV gap extracted from measurements of the thermally activated transport at zero field and half-filling, shown in Fig. 4(c) and (d). But it is not yet clear if CR of the $T_1$ transition sees the full strength of interaction contributions; this discrepancy in gap sizes suggests that the remnant Kohn’s theorem may not be not entirely eradicated by the presence of the boron nitride. On the other hand, the linear fit in Fig. 4(d) covers rather less than an order of magnitude change in temperature, so the gap size extracted from transport may not be accurate.

Finally, in Figure 5 we show the evolution of the second transition $T_2$, consisting of excitations from $n = -2 \rightarrow +1$ and $-1 \rightarrow +2$. For $|\nu| < 2$, the orbital levels in $T_2$ are always completely filled or empty. Inspection of the color map and linecuts shows that a splitting is just resolved at $\nu = 0$, with peaks of approximately equal strength. This evolves into a bright and sharp peak at $\nu = +1$ accompanied by a much weaker resonance on the high energy side, while at $\nu = +2$ the splitting persists but most of the spectral weight has shifted to the higher energy peak, which is also slightly broadened. Similar to $T_1$, at half-integer fillings only a single broad resonance is seen, and the integrated intensity remains constant over this range of $\nu$ (see Supplemental Material). The peaks are split by 2.8 meV at $\nu = 0$ and 4.7 meV at $\nu = +2$.

The $T_2$ transitions are nominally degenerate, and can be split by interactions alone [8, 17]. In particular at $\nu = 0$ and 2, two peaks of equal weight separated by 3.1 meV are predicted due to an asymmetry at finite doping between the $-2 \rightarrow 1$ and $-1 \rightarrow 2$ transitions, even if $\mu = 0$. Though close to the splittings seen here, the spectral weight distribution at $\nu = 2$ is not captured. The agreement can be improved by accounting for the Dirac mass. By fitting the data with the theory of Ref. [17] (see Fig. 4), we find the two peaks of equal weight at $\nu = 0$ can be reasonably accounted for. At $\nu = 2$, finite $\mu$ leads to a prediction of a doubly-degenerate higher-energy peak with two separate peaks at lower energies, which may account for the higher spectral weight of the upper
sitions \(-i \to i - 1\) and \(-(i - 1) \to i\) is generically broken for finite doping, but the difference is
the same for \(\nu \to (-\nu)\), preserving an overall particle-hole symmetry)

When applied to graphene, and potentially other linearly dispersing or strongly-correlated sys-
tems as well \[19\], cyclotron resonance becomes a high-resolution tool for exploring phenomena
related to the interplay of many-particle interactions and broken symmetries. Although Kohn’s
theorem, which limits the ability of CR to see interaction phenomena, is lifted by a linear dis-
ersion, a weak version remains for the lowest three, evenly-spaced LLs. This can be removed as
well by breaking translation symmetry of the graphene lattice, which here enables us to observe a
non-monotonic progression of interaction-enhanced excitations between LLs with broken spin and
valley symmetries. With continually improving device fabrication techniques, it should soon be
possible to directly observe excited states in the fractional quantum Hall regime.

**Acknowledgements** We acknowledge informative discussions with K. Shizuya, H. Fertig, and
Leo Li, and are grateful for support from the Institute of Materials Science and Engineering at
Washington University in St. Louis.

[1] G. Dresselhaus, A. Kip, & C. Kittel, “Observation of cyclotron resonance in germanium crystals,”
Physical Review 92, 827 (1953).

[2] D. J. Hilton, T. Arikawa, & J. Kono, “Cyclotron resonance,” in “Characterization of Materials”, John
Wiley and Sons (2012). Ed. Kaufmann, Elton N.

[3] W. Kohn, “Cyclotron resonance and de Haas-van Alphen oscillations of an interacting electron gas,”
Physical Review 123, 1242 (1961).

[4] D. C. Tsui, H. L. Stormer, & A. C. Gossard, “Two-dimensional magnetotransport in the extreme
quantum limit,” Physical Review Letters 48, 1559 (1982).

[5] I. V. Kukushkin, J. H. Smet, K. Von Klitzing, & W. Wegscheider, “Cyclotron resonance of composite
fermions,” Nature 415, 409 (2002).

[6] Z. Jiang, E. A. Henriksen, L.-C. Tung, Y. J. Wang, M. E. Schwartz, M. Y. Han, P. Kim, & H. L.
Stormer, “Infrared Spectroscopy of Landau Levels of Graphene,” Physical Review Letters 98, 197403
(2007).

[7] A. Iyengar, J. Wang, H. A. Fertig, & L. Brey, “Excitations from filled Landau levels in graphene,”
Physical Review B 75, 125430 (2007).

[8] Y. A. Bychkov & G. Martinez, “Magnetoplasmon excitations in graphene for filling factors \(\nu \leq 6\),”
Physical Review B 77, 125417 (2008).
[9] K. Shizuya, “Many-body corrections to cyclotron resonance in monolayer and bilayer graphene,” Physical Review B 81, 075407 (2010).
[10] R. Roldán, J.-N. Fuchs, & M. O. Goerbig, “Spin-flip excitations, spin waves, and magnetoexcitons in graphene Landau levels at integer filling factors,” Physical Review B 82, 205418 (2010).
[11] Z.-G. Chen, Z. Shi, W. Yang, X. Lu, Y. Lai, H. Yan, F. Wang, G. Zhang, & Z. Li, “Observation of an intrinsic bandgap and Landau level renormalization in graphene/boron-nitride heterostructures,” Nature Communications 5, 4461 (2014).
[12] C. Faugeras, S. Berciaud, P. Leszczynski, Y. Henni, K. Nogajewski, M. Orlita, T. Taniguchi, K. Watanabe, C. Forsythe, P. Kim, R. Jalil, A. K. Geim, D. M. Basko, & M. Potemski, “Landau Level Spectroscopy of Electron-Electron Interactions in Graphene,” Physical Review Letters 114, 126804 (2015).
[13] B. J. Russell, B. Zhou, T. Taniguchi, K. Watanabe, & E. A. Henriksen, “Many-Particle Effects in the Cyclotron Resonance of Encapsulated Monolayer Graphene,” Physical Review Letters 120, 047401 (2018).
[14] J. Sonntag, S. Reichardt, L. Wirtz, B. Beschoten, M. I. Katsnelson, F. Libisch, & C. Stampfer, “Impact of Many-Body Effects on Landau Levels in Graphene,” Physical Review Letters 120, 187701 (2018).
[15] Y. Jiang, Z. Lu, J. Gigliotti, A. Rustagi, L. Chen, C. Berger, W. de Heer, C. J. Stanton, D. Smirnov, & Z. Jiang, “Valley and Zeeman Splittings in Multilayer Epitaxial Graphene Revealed by Circular Polarization Resolved Magneto-Infrared Spectroscopy,” Nano Letters 19, 7043 (2019).
[16] I. O. Nedoliuk, S. Hu, A. K. Geim, & A. B. Kuzmenko, “Colossal infrared and terahertz magneto-optical activity in a two-dimensional Dirac material,” Nature Nanotechnology 14, 756 (2019).
[17] K. Shizuya, “Many-body effects on Landau-level spectra and cyclotron resonance in graphene,” Physical Review B 98, 115419 (2018).
[18] A. A. Sokolik & Y. E. Lozovik, “Many-body filling factor dependent renormalization of Fermi velocity in graphene in strong magnetic field,” Physical Review B 99, 085423 (2019).
[19] P. Rao & I. Sodemann, “Cyclotron resonance inside the Mott gap: A fingerprint of emergent neutral fermions,” Physical Review B 100, 155150 (2019).
[20] J. W. McClure, “Diamagnetism of graphite,” Physical Review 104, 666 (1956).
[21] V. P. Gusynin, S. G. Sharapov, & J. P. Carbotte, “Anomalous Absorption Line in the Magneto-Optical Response of Graphene,” Physical Review Letters 98, 157402 (2007).
[22] V. P. Gusynin, S. G. Sharapov, & J. P. Carbotte, “Unusual Microwave Response of Dirac Quasiparticles in Graphene,” Physical Review Letters 98, 256802 (2006).
[23] J. Gonzalez, F. Guinea, & M. Vozmediano, “Non-Fermi liquid behavior of electrons in the half-filled honeycomb lattice (A renormalization group approach),” Nuclear Physics B 424, 595 (1994).
[24] S. Das Sarma, E. H. Hwang, & W.-K. Tse, “Many-body interaction effects in doped and undoped graphene: Fermi liquid versus non-Fermi liquid,” Physical Review B 75, 121406 (2007).
[25] D. C. Elias, R. V. Gorbachev, A. S. Mayorov, S. V. Morozov, A. A. Zhukov, P. Blake, L. A. Ponomarenko, I. V. Grigorieva, K. S. Novoselov, F. Guinea, & A. K. Geim, “Dirac cones reshaped by
interaction effects in suspended graphene,” Nature Physics 7, 701 (2011).

[26] A. A. Sokolik, A. D. Zabolotskiy, & Y. E. Lozovik, “Many-body effects of Coulomb interaction on Landau levels in graphene,” Physical Review B 95, 125402 (2017).

[27] C. Kallin & B. I. Halperin, “Excitations from a filled Landau level in the two-dimensional electron gas,” Physical Review B 30, 5655 (1984).

[28] E. Batke, H. L. Stormer, A. C. Gossard, & J. English, “Filling-factor-dependent cyclotron mass in space-charge layers on GaAs,” Physical Review B 37, 3093 (1988).

[29] A. H. MacDonald & C. Kallin, “Cyclotron resonance in two dimensions: Electron-electron interactions and band nonparabolicity,” Physical Review B 40, 5795 (1989).

[30] M. Manger, E. Batke, R. Hey, K. Friedland, K. Köhler, & P. Ganser, “Filling-factor-dependent electron correlations observed in cyclotron resonance,” Physical Review B 63, 121203 (2001).

[31] Z. Schlesinger, S. J. Allen, Jr, J. C. M. Hwang, P. Platzman, & N. Tzoar, “Cyclotron resonance in two dimensions,” Physical Review B 30, 435 (1984).

[32] E. A. Henriksen, S. Syed, Y. J. Wang, H. L. Stormer, L. N. Pfeiffer, & K. W. West, “Disorder-mediated splitting of the cyclotron resonance in two-dimensional electron systems,” Physical Review B 73, 241309 (2006).

[33] G. M. Summers, R. J. Warburton, J. G. Michels, R. J. Nicholas, J. J. Harris, & C. T. Foxon, “New phases of the 2D electron system in the ultra-quantum limit observed by cyclotron resonances,” Physical Review Letters 70, 2150 (1993).

[34] N. Cooper & J. T. Chalker, “Theory of spin-split cyclotron resonance in the extreme quantum limit,” Physical Review Letters 72, 2057 (1994).

[35] T. Maag, A. Bayer, S. Baierl, M. Hohenleutner, T. Korn, C. Schüller, D. Schuh, D. Bougeard, C. Lange, R. Huber, M. Mootz, J. E. Sipe, S. W. Koch, & M. Kira, “Coherent cyclotron motion beyond Kohn’s theorem,” Nature Physics 12, 119 (2015).

[36] E. Batke, D. Heitmann, J. Kotthaus, & K. Ploog, “Nonlocality in the two-dimensional plasmon dispersion,” Physical Review Letters 54, 2367 (1985).

[37] Y. Zhao, D. C. Tsui, M. B. Santos, M. Shayegan, R. A. Ghanbari, D. A. Antoniadis, & H. I. Smith, “Grating-induced cyclotron-resonance anomaly in GaAs/AlₓGa₁₋ₓAs heterostructures,” Physical Review B (1995).

[38] I. V. Kukushkin, J. H. Smet, D. Schuh, W. Wegscheider, & K. Von Klitzing, “Dispersion of the Composite-Fermion Cyclotron-Resonance Mode,” Physical Review Letters 98, 066403 (2007).

[39] C. Kallin & B. I. Halperin, “Many-body effects on the cyclotron resonance in a two-dimensional electron gas,” Physical Review B 31, 3635 (1985).

[40] L. Wang, I. Meric, P. Y. Huang, Q. Gao, Y. Gao, H. Tran, T. Taniguchi, K. Watanabe, L. M. Campos, D. A. Muller, J. Guo, P. Kim, J. Hone, K. L. Shepard, & C. R. Dean, “One-Dimensional Electrical Contact to a Two-Dimensional Material,” Science 342, 614 (2013).

[41] E. A. Henriksen, P. Cadden-Zimansky, Z. Jiang, Z. Q. Li, L.-C. Tung, M. E. Schwartz, M. Takita, Y. J.
Wang, P. Kim, & H. L. Stormer, “Interaction-Induced Shift of the Cyclotron Resonance of Graphene Using Infrared Spectroscopy,” Physical Review Letters 104, 067404 (2010).

[42] C. R. Dean, A. F. Young, I. Meric, C. Lee, L. Wang, S. Sorgenfrei, K. Watanabe, T. Taniguchi, P. Kim, K. L. Shepard, & J. Hone, “Boron nitride substrates for high-quality graphene electronics,” Nature Nanotechnology 5, 722 (2010).

[43] A. A. Zibrov, C. Kometter, H. Zhou, E. M. Spanton, T. Taniguchi, K. Watanabe, M. P. Zaletel, & A. F. Young, “Tunable interacting composite fermion phases in a half-filled bilayer-graphene Landau level,” Nature 549, 360 (2017).

[44] P. T. Coleridge, R. Stoner, & R. Fletcher, “Low-field transport coefficients in GaAs/Ga$_{1-x}$Al$_x$As heterostructures,” Physical Review B 39, 1120 (1989).

[45] See Supplemental Material.

[46] S. Syed, M. J. Manfra, Y. J. Wang, R. J. Molnar, & H. L. Stormer, “Electron scattering in AlGaN/GaN structures,” Applied Physics Letters 84, 1507 (2004).

[47] M. Yankowitz, J. Xue, D. Cormode, J. D. Sanchez-Yamagishi, K. Watanabe, T. Taniguchi, P. Jarillo-Herrero, P. Jacquod, & B. J. LeRoy, “Emergence of superlattice Dirac points in graphene on hexagonal boron nitride,” Nature Physics 8, 382 (2012).

[48] B. Hunt, J. D. Sanchez-Yamagishi, A. F. Young, M. Yankowitz, B. J. Leroy, K. Watanabe, T. Taniguchi, P. Moon, M. Koshino, P. Jarillo-Herrero, & R. C. Ashoori, “Massive Dirac Fermions and Hofstadter Butterfly in a van der Waals Heterostructure,” Science 340, 1427 (2013).

[49] J. Jung, A. M. DaSilva, A. H. MacDonald, & S. Adam, “Origin of band gaps in graphene on hexagonal boron nitride,” Nature Communications 6, 6308 (2015).

[50] A. Fischer, A. Dzyubenko, & R. Römer, “Localized collective excitations in doped graphene in strong magnetic fields,” Physical Review B 80, 165410 (2009).

[51] O. Dial, R. C. Ashoori, L. N. Pfeiffer, & K. W. West, “High-resolution spectroscopy of two-dimensional electron systems,” Nature 448, 176 (2007).

[52] A. F. Young, C. R. Dean, L. Wang, H. Ren, P. Cadden-Zimansky, K. Watanabe, T. Taniguchi, J. Hone, K. L. Shepard, & P. Kim, “Spin and valley quantum Hall ferromagnetism in graphene,” Nature Physics 8, 550 (2012).

[53] R. S. Deacon, K.-C. Chuang, R. J. Nicholas, K. S. Novoselov, & A. K. Geim, “Cyclotron resonance study of the electron and hole velocity in graphene monolayers,” Physical Review B 76, 081406 (2007).

[54] N. M. R. Peres, F. Guinea, & A. H. Castro Neto, “Electronic properties of disordered two-dimensional carbon,” Physical Review B 73, 125411 (2006).

[55] P. Plochocka, C. Faugeras, M. Orlita, M. L. Sadowski, G. Martinez, M. Potemski, M. O. Goerbig, J.-N. Fuchs, C. Berger, & W. A. de Heer, “High-Energy Limit of Massless Dirac Fermions in Multilayer Graphene using Magneto-Optical Transmission Spectroscopy,” Physical Review Letters 100, 087401 (2008).

[56] A. H. MacDonald, H. C. A. Oji, & S. M. Girvin, “Magnetoplasmon Excitations from Partially Filled
Landau Levels in Two Dimensions,” Physical Review Letters 55, 2208 (1985).

[57] Y. Hattori, T. Taniguchi, K. Watanabe, & K. Nagashio, “Layer-by-Layer Dielectric Breakdown of Hexagonal Boron Nitride,” ACS Nano 9, 916 (2015).

[58] J. C. W. Song, A. V. Shytov, & L. S. Levitov, “Electron Interactions and Gap Opening in Graphene Superlattices,” Physical Review Letters 111, 266801 (2013).

[59] H. Kim, N. Leconte, B. L. Chittari, K. Watanabe, T. Taniguchi, A. H. MacDonald, J. Jung, & S. Jung, “Accurate Gap Determination in Monolayer and Bilayer Graphene/h-BN Moiré Superlattices,” Nano Letters 18, 7732 (2018).
FIG. 1. Cyclotron resonance transitions in graphite-gated monolayer graphene. (a) Color map of the normalized transmission spectra $\Delta T/T$ in the mid-infrared as a function of the LL filling factor, $\nu$, measured in the device shown inset to (c). Several sharp CR transitions are visible, labeled $T_1$ through $T_5$. The higher noise in the region of $T_1$ is due to overall lower transmission in the 60-150 meV energy range, compared to the other transitions (and below 60 meV the signal goes to zero) [45]. (b) Schematic showing the allowed Landau level transitions at $\nu = 0$, consisting of nominally degenerate pairs. (c) Representative linecut of the color map, at $\nu = 0$. (d) Evolution of CR lineshape at $\nu = 0$ with sample quality: the top trace is from a graphene-on-SiO$_2$ device with mobility of 17,000 cm$^2$/Vs [41], the middle is from an hBN-encapsulated device on SiO$_2$ [13], and the bottom is the present graphite-gated, hBN-encapsulated device; the latter two have the same mobility, $\mu \approx 200,000$ cm$^2$/Vs. (e) CR lifetime $\tau_{CR} = \hbar/\Gamma$ at $\nu = 0$ ($\Gamma$ the half-width at half-max) vs a spread of quantum scattering times $\tau_q$, derived from Shubnikov-de Haas oscillations acquired for a range of carrier densities at 3 K (colors correspond to traces (d)). The dashed line marks $\tau_{CR} = \tau_q$. 
FIG. 2. Evolution of transition $T_1$ vs filling factor. (a) High-resolution map of $T_1$ vs filling factor from $\nu = -0.07$ to $\nu = +2.5$. Traces were acquired every $\delta \nu = 0.026$, with additional traces at $\nu = 3, 3.5, ... 6$. Starting with a single bright peak at $\nu = 0$, four peaks appear near $\nu = +1$ which become two peaks at $\nu = 2$ and higher. By $\nu = 6$ the $T_1$ transition is extinguished as the participating levels are filled. (b) Detail of transitions by linecuts at half-integer fillings. The linewidths at $\nu = +1$ are the narrowest observed with $\tau_{CR}$ reaching 2.5 ps, or a resonance quality factor $Q = 220$. The widths are not intrinsic, as they continue to narrow at higher instrumental resolution (see Supplemental Material). In between integer values of $\nu$, only a single broad resonance is resolved. (c) Schematic of transitions involving the $n = -1, 0,$ and $+1$ LLs. Solid (dashed) lines indicate the $K (K')$ valleys, with valley gaps $\Delta v_i$ and spin splittings $\Delta z_i$ explicitly included. The Fermi energy $E_F$ is shown as a dotted line. Spin and valley-preserving CR transitions are shown in red, and labeled $a-d$, $e-h$, and $i-l$ for those at $\nu = 0, 1$ and 2, respectively. Gaps indicated in this schematic represent single particle levels enhanced by electron-electron interactions as discussed in the text.
FIG. 3. Particle-hole asymmetry. (a) Closeup of $T_1$ transition from Fig. 1(a), showing a small but clear asymmetry in the spectrum for positive vs. negative filling factors. (b) and (c) show higher resolution (finer abscissa spacing) measurements to compare the four-fold splittings at $\nu = -1$ and $+1$, respectively.
FIG. 4. **Comparison with theory, and activated transport.** The (a) \( \nu = +2 \) and (b) \( \nu = 0 \) transition energies are shown parameterized as effective band velocities, \( v_{\text{eff}} \), as a function of transition number \( T_i \) (open circles). The dashed lines are predictions of the transition energies by two approaches: the tan lines show Eq. 1 that gives the single-particle LL separation with a finite Dirac mass, while the blue lines are from the many-particle theory of Ref. [17]. Here both models are fit to the data in (a), and the fit parameters are then used to generate the curves in (b). (c),(d) Temperature-dependent resistance for charge-neutral graphene at zero magnetic field, and Arrhenius activation plot yielding a gap of 15.0 meV, respectively.
FIG. 5. **Evolution of the second CR transition** $T_2$ **vs filling factor.** (a) High-resolution map of $T_2$ for the same filling factor range as Fig. 1. The horizontal band at 298 meV is due to harmonics of 60 Hz. Two peaks are just resolved at $\nu = 0$, while the spectrum is dominated by a single peak at $\nu = +1$, and two peaks appear again at $\nu = +2$ albeit with more intensity in the higher peak. The resonances are broader than those at $T_1$, and exhibit a high energy tail which may be a result of multiple reflections in the substrate [41]. (c) Schematic of $T_2$ transitions, with the $K$ ($K'$) valley shows as a solid (dashed) line. The Zeeman gap is suppressed.