Cascade of isospin phase transitions in Bernal-stacked bilayer graphene at zero magnetic field

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Emergent phenomena arising from the collective behaviour of electrons is expected when Coulomb interactions dominate over the kinetic energy, and one way to create this situation is to reduce the electronic bandwidth. Bernal-stacked bilayer graphene intrinsically supports saddle points in the band structure that are predicted to host a variety of spontaneous symmetry-broken states26–31. Here we show that bilayer graphene displays a cascade of symmetry-broken states with spontaneous spin and valley isospin ordering at zero magnetic field. We independently tune the carrier density and electric displacement field to explore the phase space of isospin order. Itinerant ferromagnetic states emerge near the conduction and valence band edges with complete spin and valley polarization. At larger hole densities, twofold degenerate quantum states exist with complete spin and valley polarization at zero magnetic field. Both symmetry-broken states display enhanced layer polarization, suggesting a coupling to the layer degree of freedom. These states occur in the absence of a moiré superlattice and are intrinsic to natural graphene bilayers. Therefore, we demonstrate that bilayer graphene represents a related but distinct approach to produce collective behaviour from flat dispersion, complementary to engineered moiré structures.

Large electronic densities of states enhance the effects of interactions between electrons and give rise to collective electronic states such as Stoner ferromagnetism and Bardeen–Cooper–Schrieffer superconductivity. In quantum Hall ferromagnetism, a perpendicular magnetic field transforms a two-dimensional electronic system into a series of macroscopically degenerate flat bands that spontaneously polarize at half-filling. Moiré superlattices in magic-angle twisted bilayers and trilayers of graphene generate flat bands in zero field but without a moiré superlattice and are intrinsic to natural graphene bilayers. Therefore, we demonstrate that bilayer graphene represents a related but distinct approach to produce collective behaviour from flat dispersion, complementary to engineered moiré structures.

To explore this possibility, we probe the electronic compressibility of ultraclean bilayer graphene samples at low carrier densities and in the presence of perpendicular displacement fields. We fabricate dual-gated capacitor devices using graphite for the top- and bottom-gate electrodes and measure the penetration field capacitance $C_p$, related to the inverse compressibility ($\rho_n/\rho_v$)$^{-1}$, from the top-gate electrode to the bottom-gate electrode (Fig. 1b and Methods). Depending on the compressibility of bilayer graphene, $C_p$ varies between the geometric capacitance between the top and bottom gates (when the bilayer is incompressible or gapped) and zero (when the bilayer is infinitely compressible, as in a metal).

Using the top- and bottom-gate voltages (Fig. 1b), we vary the nominal carrier density $n_n \equiv (C_n V_n + C_v V_v)/\epsilon_0$, whereas the out-of-plane displacement field $D_\parallel \equiv (C_n V_n - C_v V_v)/2\epsilon_0$ is fixed, where $C_n$, $C_v$ are the geometric capacitances of the top and bottom gates with respect to the bilayer, $\epsilon_0$ is the dielectric constant, $e$ is the electron charge, $V_n = V_n - V_0$ and $V_v = -V_0$. At $D_\parallel = 0$, the measured $C_p$ shows a jump in charge compressibility at charge neutrality, with holes exhibiting higher compressibility compared with electrons (Fig. 1c). In the presence of a large, positive displacement field, the small peak at charge neutrality develops into a large, incompressible peak (Fig. 1d), as expected for the tunable bandgap in the band structure of bilayer graphene. A series of peaks appear on either side of charge neutrality (Fig. 1f) that do not match the features in the non-interacting electronic structure (Fig. 1c–e). The non-interacting bands of bilayer graphene develop a single divergence in the density of states on each side of the gap (Fig. 1e and Extended Data Fig. 2) arising from a Van Hove singularity and Lifshitz transition (LT; Fig. 1c) at the saddle point within each band4,50. However, two peaked regions appear on the hole side of the gap (Fig. 1f, labelled as A and B).
Fig. 1 | Flat dispersion in Bernal-stacked bilayer graphene. a, Bernal-stacked (AB) bilayer graphene structure. b, Schematic of \( C_\parallel \) measurement. Effective gate voltages are defined as \( V_e = V_{g} - V_\parallel \) and \( V_\parallel = -V_g \), where \( V_g \) is the d.c. voltage applied to the bilayer graphene. c, Measured \( C_\parallel \) versus carrier density \( n_\parallel \) at zero electric displacement and \( D_\perp = 0 \) (left). Corresponding calculated, non-interacting band structure along the K→Γ direction showing two of the four Dirac nodes (right). Although only one valley is shown, the non-interacting bands are 4\times degenerate due to the combined spin and valley degrees of freedom. The inset shows the Fermi contour showing the effect of trigonal warping around K and relative position of the Dirac nodes (points). The dashed line indicates the \( k \) slice along which the bands are shown. d, Measured \( C_\parallel \) at \( D_\perp = 0.75 \) \( \text{V nm}^{-1} \) showing an incompressible peak in the bandgap of the bilayer (left). Calculated, non-interacting band structure with emerging flat dispersion driven by the \( D_\parallel \) field (right). e, Calculated density of states (DOS) versus carrier density for two values of interlayer potential asymmetry, \( \Delta_\parallel \), showing a single Van-Hove singularity per conduction/valence band. The gap between the bands collapses to a single point (\( n = 0 \)) when plotted versus density, rather than energy. The inset shows the band surfaces with isoenergy contours for different carrier densities (only the valence band is shown for \( \Delta_\parallel = 0 \)). Details of the trace from \( D_\perp = 0.75 \) \( \text{V nm}^{-1} \) (as shown in d) showing two prominent peaks on the hole side of the incompressible peak, labelled as A and B, with the corresponding trace from \( D_\perp = -0.75 \) \( \text{V nm}^{-1} \) showing shoulder features on the electron side of charge neutrality, with label C indicating a broad shoulder region where additional features emerge in finite magnetic field. g, Map of \( C_\parallel \) measured as a function of \( n_\parallel \) and \( D_\perp \) at 40 mK. Equivalent map for electrons with \( D_\perp < 0 \) (which is clearer than \( D_\perp > 0 \) for electrons; Supplementary Fig. 3 shows the full map) (right). Positions of the peaks in f depend linearly on the magnitude of \( D_\perp \).

separated by highly compressible minima. By measuring these peaks as a function of \( D_\perp \) (Fig. 1g), we see that their location in electron density varies linearly with the displacement field over a wide range and tends to zero as \( D_\perp \to 0 \).

To understand the origin of the \( D_\perp \)-tunable peaks in compressibility, we explore quantum oscillations in the presence of a small perpendicular magnetic field \( B_\perp \) with respect to the graphene bilayer. In Fig. 2b, we show a map of \( C_\parallel \) measured at \( B_\perp = 2 \) T. A series of cyclotron gaps are evident in the \( C_\parallel \) map as narrow vertical features at integer filling factors, \( \nu \equiv n/\hbar eB_\perp \), with Planck constant \( \hbar \), appearing due to Landau-level formation in the bilayer. Near \( D_\perp = 0 \), the density separation between the cyclotron gaps is \( \Delta n_\parallel \approx 9.7 \times 10^{10} \text{cm}^{-2} \) or \( \Delta \nu = 2 \) from \( \nu = -4 \) to \( \nu = +4 \), and \( \Delta n_\parallel \approx 1.9 \times 10^{11} \text{cm}^{-2} \) or \( \Delta \nu = 4 \) for \( |\nu| \geq 4 \) (Fig. 2a). The set of states from \( -4 < \nu < 4 \) emerges from the zero-energy Landau level (ZLL), characterized by degenerate spin, valley and Landau-level orbital quantum numbers. A density separation of \( \Delta \nu = 2 \) or \( \Delta \nu = 1 \) is expected in the ZLL due to Coulomb-driven exchange splitting, particularly in the presence of a layer-polarizing displacement field. Outside the ZLL and at small magnetic fields, exchange generally lifts the \( 4\times \) degeneracy of each Landau level gradually as the displacement field is increased from zero. However, at finite displacement field and within the boundaries of regions A and B observed at \( B = 0 \) (between the \( C_\parallel \) minima at \( B = 0 \), the Landau levels exhibit sharp transitions to \( \Delta \nu = 1 \) within region B and \( \Delta \nu = 2 \) within region A (Fig. 2b). Starting at a finite \( D_\perp \) value that depends on density, cyclotron gaps emerge for \( \nu = -3, -4 \) and \( -5 \) within region B and \( \nu = -6, -8 \) and \( -10 \) within region A. Interestingly, a region with \( \Delta \nu = 1 \) also emerges for electrons (Fig. 2b, labelled as C), although the parent state at zero field is not as evident as in regions A and B.

The observation of sharp transitions between Landau-level spacing \( \Delta \nu \) at fixed \( B_\perp \) suggests a sequence of phase transitions between the phases with broken spin–valley (isospin) symmetry. That is, the changes in Landau-level degeneracy are driven by \( D_\perp \)-field-induced spontaneous symmetry breaking, even at relatively low magnetic fields, in contrast to the usual exchange-driven symmetry breaking generated by large magnetic fields. At \( B = 0 \), these transitions manifest as sharp minima in \( C_\parallel \) whereas the symmetry-broken regions occur within regions of reduced compressibility (Fig. 1f, labelled peaks).

The assignment between \( \Delta \nu \) and Fermi-surface degeneracy within each region at fixed field can be further supported by investigating the magnetic-field dependence of the symmetry-broken states as \( B_\perp \to 0 \). In Fig. 3b, we show the field dependence of the Landau levels taken along a density cut at fixed \( D_\perp \) (Fig. 3a, along the dashed line). The displacement field is fixed at \( D_\perp = 0.67 \) \( \text{V nm}^{-1} \) to sample a portion of each symmetry-broken region as well as the \( 4\times \)-degenerate regions away from the symmetry-broken states. Quantum oscillations with differing frequencies are observed in the resulting field dependence with clear boundaries near the edges of each symmetry-broken region (Fig. 3b). Taking the Fourier transform along the field direction in \( 1/B_\perp \) for a range of densities in each region (Fig. 3b, top) reveals the dominant oscillation frequency in each sampling window (Fig. 3c).

This oscillation frequency depends linearly on the total carrier density \( n_\parallel \). To compare the frequencies over a range of densities,
we normalize the oscillation frequency \( f_0 \) by the carrier density in units of magnetic flux, namely, \( n_0 \Phi_0 \), where \( \Phi_0 = \hbar / e \), equivalent to the \( k \)-space area enclosed by an orbit around the Fermi surface. Using this normalization, Fourier transforms can be averaged over a window of densities to optimize the signal-to-noise ratio (the boxed regions in Fig. 3b are chosen to maximize the window with consistent oscillations) and plotted together (Fig. 3c). At large hole densities (Fig. 3b, region D), the resulting Fourier transform displays a dominant peak at \( f_0/n_0\Phi_0 = 1/4 \) (with additional harmonics arising from weak, gradual symmetry breaking from the finite magnitude of \( D_x \)), indicating that the carriers are predominantly divided evenly between four coexisting Fermi surfaces, as expected for a 4x-degenerate band. Within the intermediate density range at which \( \Delta \nu = 2 \) Landau levels are shown at 2T (Fig. 2b,c, region A), a Fourier analysis reveals a peak at \( f_0/n_0\Phi_0 = 1/2 \), or a reduction to two Fermi surfaces. The peak at 1/2 is consistent with the separation of Landau levels \( \Delta \nu = 2 \) at fixed field, further suggesting that the system has spontaneously developed an exchange splitting that favours filling two of the four isospin flavours. Notably, an additional low-frequency peak emerges in curve A that indicates some coexistence between the Fermi surfaces with differing areas (discussed later). In regions B and C (for holes and electrons, respectively), the peak Fourier amplitude occurs at \( f_0/n_0\Phi_0 = 1 \). Thus, only a single Fermi surface (with one isospin flavour) is occupied in these density ranges, each of which is, therefore, a ferromagnetic spin- and valley-polarized state, akin to a quantum Hall ferromagnet, but driven by the \( D_x \) field.

To determine the precise nature of the symmetry-broken states, it is necessary to experimentally probe the flavour degrees of freedom. Using the fact that in-plane magnetic fields preferentially couple to spin, we can follow the evolution of phase boundaries between the symmetry-broken states in \( B_z \) to reveal changes in spin ordering (Fig. 4a). Taking a cut along the density direction at fixed \( D_z = 0.69 \ V \ \text{nm}^{-1} \), we track the positions of capacitance minima, taken to be the boundaries of each symmetry-broken phase as we

**Fig. 2 | Degeneracy of symmetry-broken states in \( B_z \).** a, Trace of \( C_p \) measured at \( B_z = 2 \ T \) and \( D_z = 0 \) with 2x-degenerate incompressible peaks (separation between subsequent filling factors, \( \Delta \nu \)) in the ZLL and 4x-degenerate peaks elsewhere. b, Map of \( C_p \) measured at \( B_z = 2 \ T \) versus \( n_0 \) and \( D_z \), with Landau-level filling factors (\( \nu_{LL} \)) shown above. c, Details of line traces taken along the matching dotted lines in b, showing sharply different \( \Delta \nu \) from \( D_z = 0 \).

**Fig. 3 | Fermi-surface signatures in quantum oscillations of the symmetry-broken states.** a, Density (filling factor, \( \nu_{LL} \)) versus \( D_z \), map of \( C_p \) measured at \( B_z = 2.3 \ T \). b, Measured \( C_p \) versus \( n_0 \) and \( B_z \) at fixed \( D_z = 0.67 \ V \ \text{nm}^{-1} \) (along the dashed line in a). c, Fast Fourier transform of magneto-oscillations from b taken along 1/\( B_z \) and averaged over the range of densities indicated by the boxed regions in b. The width of each averaging window is selected to optimize the signal-to-noise ratio. The resulting frequencies are normalized by the total carrier density \( n_0 \) and magnetic flux quantum \( \Phi_0 \) for cross-comparison. The insets show the schematic bands for each curve with the implied degeneracy of the occupied flavours indicated (with valleys superimposed).
sweep a pure $B_\|$ field (Fig. 4b). The resulting A–B phase boundary decreases in density as the magnitude of $B_\|$ increases (Fig. 4c). If the primary effect of $B_\|$ is to introduce a Zeeman splitting of the spin states in the system, we can interpret the reduction in density of the A–B transition point as resulting from the difference in Zeeman energy between the two phases with $B_\|$ (ref. 13). From fixed-field measurements (Fig. 2) and quantum oscillations (Fig. 3), we can attribute region B to a onefold (1×)-degenerate state with full spin and valley polarization. The change in A–B boundary suggests that the spin order differs between the two phases, suggesting that region A is symmetry broken, but without complete spin polarization below $B_\| \leq 4T$. The D–A phase boundary also shifts slightly in the presence of a large in-plane field since the fully symmetric, 4×-degenerate region D is spin unpolarized. Together with the fact that the width of region A along the density axis increases as $B_\|$ is applied, this suggests region A may possess some degree of partial spin polarization that is stabilized by finite $B_\|$. Considering that the twofold features in region A observed in fixed $B_\|$ are weaker in comparison to region B, it is likely that region A is only partially polarized. This interpretation is consistent with the observation of an additional low-frequency peak in curve A below 1/4 (Fig. 3c), indicating that the occupation is shared between at least two Fermi surfaces with differing $k$-space geometry.

Finite $B_\|$ also causes the boundaries between symmetry-broken states to become sharper (Fig. 4d) compared with $B_\|=0$, with decreasing $C_s$ minima trending towards negative compressibility as $B_\|$ increases (Fig. 4b). Negative compressibility is an indication of electronic ordering in the system, typically emerging in proximity to phase transitions into ordered states16. This trend supports the interpretation of the sharp $C_s$ minima as occurring at or near the phase boundaries between the symmetry-broken phases. Notably, the boundaries of the 1×-degenerate electron state (Fig. 4d, region C) are more clearly resolved at $B_\|=8$ T compared with $B_\|=0$, lending further support to the assignment of the symmetry-broken state observed in Figs. 2 and 3.

To further probe the character of the symmetry-broken states, we make use of layer sensitivity provided by the capacitance measurement. By applying out-of-phase signals to the top and bottom gates of the sample as we measure the response on the bilayer, we nullify the contribution from electronic compressibility and instead probe the tendency of the system to polarize across the layers of the bilayer18. This measurement—the differential capacitance $C_{\text{diff}}$—is positive for
one sign of the displacement field and negative for the other (Fig. 4e). The measured $C_{diff}$ is related to layer polarization $p = n_1 - n_2$ (for layer densities $n_1$ and $n_2$) by integration, namely, $p \propto \int C_{diff} \, dn$, and therefore, the sign indicates the direction of polarization; however, the magnitude of $C_{diff}$ indicates the tendency of the system to fill one layer or the other as carriers are added (Extended Data Fig. 3). The overall slope of $C_{diff}$ (Fig. 4e–g) largely reflects the simple electrostatic effect of $D_{int}$; however, the small oscillations observed on top of this background are instead related to the symmetry-broken states. The regions indicated by filled symbols in Fig. 4f align with the symmetry-broken regions A and B measured in $C_p$. By comparing the curves measured at 1.8 K (solid) to the ones measured at 40.0 K (dashed) where the signature of the symmetry-broken states is quenched, we find that $C_{diff}$ is enhanced at the phase boundaries at low temperatures (Extended Data Fig. 4). Taking the difference of the low- and high-temperature data ($\Delta C_{diff} = C_{diff}(B) - C_{diff}(A)$) and integrating along the density direction, we obtain the enhancement in layer polarization due to the symmetry-broken phases as $\Delta \rho = \Delta C_{diff} \, dn$. (Fig. 4g). This excess polarization $\Delta \rho$ shows that the degree of layer polarization is enhanced within the symmetry-broken phases, suggesting that charge transfer occurs between the layers due to isospin polarization.

Exchange splitting driven by displacement-field-induced transitions is an exciting direction for further study in bilayer graphene. There is an interesting parallel between the role of the $D_{int}$ field in driving the bilayer into isospin-polarized states and quantum Hall ferromagnetism.

Although the $B_0$ field induces ferromagnetism by forming extremely flat Landau levels at odd integer fillings in any sufficiently clean two-dimensional electronic system, in bilayer graphene, the $D_{int}$ field provides an alternative tuning parameter for generating flat dispersion. Crucially, $D_{int}$ can drive the system into isospin-polarized states at $B = 0$. This mechanism is reminiscent of, but distinct from, the observation of correlated states in moiré systems where a periodic superlattice induces flat bands and large densities of states. Here no superlattice is required and, more importantly, the band dispersion is tunable by an external parameter in situ. This presents a distinct advantage over many moiré systems where the bandwidth is fixed by the angle of rotation or mismatch between the constituent lattice periodicities. For this reason, and due to the ease with which bilayer graphene can be prepared and experimentally measured relative to ABC trilayers and other graphene multilayers, bilayer graphene is a strong candidate for further investigation in the context of interaction-driven physics.

Online content

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Methods

Device fabrication. Each of the devices in this study were fabricated using a single boron nitride (BN) crystal cut into two pieces to ensure a symmetric gating environment for the bilayer graphene. Each BN flake was etched into two using reactive ion etching. We first picked up a graphite flake using one of the two BN pieces and put it onto a substrate with prepatterned markers. The stack was annealed at 350 °C for at least 3h with H2/Ar gas. Then, the other half of the BN flake was used to pick up the bilayer graphene and transferred onto the BN/graphite stack. Another graphite flake was then put on top in a separate transfer process. The whole device structure was assembled using standard dry-transfer techniques. Graphite gates were deliberately employed to minimize charge disorder and ensure high device quality. Afterwards, electrical contact was defined by electron-beam lithography and made through a top contact method using Cr/Pd/Au.

For device I, the top and bottom BN flakes were aligned at 0°, whereas in device II, the two BN flakes were rotated by 90° with respect to one another. In all the measurements, no superlattice peaks were observed. Therefore, it is unlikely that the bilayer graphene is aligned to any of the BN flakes. We believe that the relative BN angles are not important for our reported observations as we also performed the same measurement in a purposely misaligned sample and the same phenomenon was consistently observed. All the data in the main text were measured in device II, unless otherwise specified.

Capacitance measurements. All the capacitance data in the study were measured by balancing a cryogenic capacitance bridge circuit at a fixed point in the parameter space and measuring the off-balance voltage via a low-temperature amplifier circuit (Extended Data Fig. 5). Sinusoidal a.c. voltage signals were simultaneously applied to the device and a reference capacitor whose value was separately measured. The phase and amplitude of the signal applied to the reference capacitor were varied to nullify the signal at the balance point of the bridge circuit. As the capacitance of the device changes, an off-balance voltage develops at the balance point whose magnitude is proportional to the device capacitance.

The sample and reference capacitances are far smaller than the parasitic capacitances of the cryostat cabling and are difficult to directly measure due to signal shunting. A custom low-temperature amplifier is, therefore, fabricated and placed as close to the sample as possible in the cryostat (typically a few millimetres away). This amplifier transforms the impedance of the bridge circuit and provides a small amount of voltage gain. In the \( C_{\text{ref}} \) measurements, the amplifier and bridge balance point were connected to the back gate of each device, whereas in \( C_{\text{ref}} \) measurements, they were connected to the bilayer graphene. Excitation signals in this study were \( V_r = 5 \text{mV} \), applied at a frequency of 51.747 kHz either to the top gate (as in the \( C_{\text{ref}} \) measurements) or simultaneously to the top and bottom gates (as in the \( C_{\text{ref}} \) measurements). All the measurements were performed between 40 and 100 mK, unless otherwise specified.

The \( C_{\text{ref}} \) measurements were performed by tuning the amplitude of two out-of-phase signals simultaneously applied to the top and bottom gates to nullify the contribution of the electronic compressibility at the balance point. Due to careful device design and the use of symmetric BN layers, the optimal ratio of top and bottom excitation amplitudes was exactly 1. Parallel capacitance from unintentional geometric contributions was separately nullified using a signal applied to the reference capacitor in the same fashion as in the \( C_{\text{ref}} \) measurements.

Data availability

Source data are provided with this paper. All other data that support the plots within this paper and other findings of this study are available from the corresponding author upon reasonable request.

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Author contributions

S.C.d.l.B. and S.A. performed the measurements and calculations, analysed the data and wrote the manuscript with input from all the authors. Z.Z. fabricated the samples. K.W. and T.T. grew the hexagonal boron nitride crystals. S.C.d.l.B, Q.M. and Z.Z. conceived of the measurements, and Q.M., P.J.-H. and R.A. supervised the project.

Competing interests

The authors declare no competing interests.

Additional information

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Extended Data Fig. 1 | Evolution of bilayer and rhombohedral trilayer graphene bands with potential asymmetry. a Evolution of bilayer and rhombohedral (ABC) trilayer graphene bands along $k_x$ as interlayer asymmetry, $\Delta_g$, is increased from zero. ‘Flatness’ of the dispersion in both cases is apparent as the number of states near the valence band edge occupies a large extent in $k_x$. c Qualitative comparison of the $k$-space extent of the valence band edge, $|k_{VBE}|$, for bilayer and ABC trilayer graphene as a function of $\Delta_g$ extracted from the matching points in a–b. Bands were calculated using the same tight-binding model and parameters as in Supplementary Fig. 2, following Ref. 41.
Extended Data Fig. 2 | Single particle density of states for bilayer graphene. a Calculated single particle density of states (DOS) versus carrier density for fixed values of interlayer potential asymmetry, $\Delta g$, showing a single Van Hove singularity per conduction/valence band. The gap between bands collapses to a single point ($n=0$) when plotted versus density, rather than energy. Inset: Band surfaces with isoenergy contours for different carrier densities. b Inverted density of states, $\text{DOS}^{-1}$, showing a sharp minimum (dark feature) at the position of the large Van Hove singularity (VHs) trending toward larger hole density as potential asymmetry increases (in analogy to $D\perp$-field in experiment). A smaller VHs is also seen for electrons, but broadens with increasing $\Delta g$. The DOS was obtained by numerical integration of tight binding bands using the model and parameters of Ref. 41.
Extended Data Fig. 3 | Layer polarization in the zero-energy Landau level. **a** Map of $C_{\text{eff}}$ measured in Device II at $B = 4$ T and 1.7 K. Sloped lines throughout the map arise from cyclotron gaps in the graphite gate electrodes. **b** Density cuts taken at fixed displacement fields indicated by the dotted lines in **a**.
Extended Data Fig. 4 | Displacement field dependence of $C_{\text{diff}}$ at high and low temperatures. Maps of $C_{\text{diff}}$ measured in Device II at a 1.8 K and b 40 K. In a, regions of enhanced layer polarization at finite displacement field correspond to phase transitions between symmetry-broken states. c Difference map showing the data in b subtracted from those in a. $B = 0$ in both measurements.
Extended Data Fig. 5 | Capacitance bridge amplifier circuit. Circuit schematic showing the dual-amplifier configuration used to measure $C_p$ and $C_{\text{diff}}$ in the same device. Combinations of AC and DC voltages are applied to terminals $V_{\text{tg}}$, $V_{\text{bl}}$, $V_{\text{bg}}$ in order to gate the bilayer graphene, excite charge for the capacitance measurement, and to simultaneously power the relevant amplifier. Amplifier 2 is used for $C_p$ measurements with an AC excitation on the top gate, while Amplifier 1 is used for $C_{\text{diff}}$ measurements, with AC excitations applied to the top and bottom gates simultaneously.