Two graphene monolayers, placed on top of each other with a small rotational misalignment between their crystallographic axes, form a long-wavelength moiré superlattice. The electronic properties of such a superlattice depend on the relative twist angle, \( \theta \), between the graphene layers as well as their interlayer hybridization. A particularly interesting case is that of small-angle (\( \theta < 3^\circ \)) TBG (SA-TBG), where hybridization is strong, and which, for a certain range of angles, features intriguing interaction-driven phenomena including, but not limited to, superconductivity, correlated insulator states, and orbital ferromagnetism. The low-energy single-particle band structure of SA-TBG resembles that of monolayer graphene (MLG) but is characterized by a decreased Fermi velocity, \( v_F \), and a reduced Brillouin zone (BZ). Like the BZ of MLG, the reduced BZ is hexagonal and contains two minivalleys located at the \( k_m \) and \( k'_m \) high-symmetry points. The minivalleys are spaced apart by a relatively small (in comparison to MLG) distance \( \Delta k = (4\pi/a)\sin(\theta/2) \), where \( a \) is the lattice constant of MLG (Fig. 1a). In MLG, the intervalley separation is sufficiently large so as to suppress intervalley electron scattering, provided that atomically-sharp defects are absent. In this work, we show that the opposite is true for SA-TBG, where strong intervalley scattering significantly affects its transport properties.

Our devices are multi-terminal Hall bars composed of SA-TBG encapsulated between two relatively thin (<100 nm thick) slabs of hexagonal boron nitride (hBN). The Hall bars were produced by a combination of tear-and-stack and hot release methods, were endowed with quasi-one-dimensional contact and had typical width of about 2 \( \mu \)m as shown in Fig. 1. (See Supplemental Material\cite{supplementary}). Figure 1c shows a typical dependence of the longitudinal resistivity, \( \rho_{xx} \), versus the externally-induced total carrier density \( n \), measured in one of our devices at two representative temperatures, \( T \). At small \( n \), the \( \rho_{xx}(n) \) dependence resembles that of monolayer graphene (MLG); namely, it exhibits a sharp peak of about 2 k\( \Omega \) at the charge neutrality point (CNP) that rapidly drops to 20–50 \( \Omega \) with increasing \( |n| \). Upon further doping, \( \rho_{xx}(n) \) exhibits a steep rise at \( |n| \approx 6 \times 10^{12} \) cm\(^{-2} \), which corresponds to full filling of the first superlattice miniband in accord with previous studies on SA-TBG\cite{supplementary}. Three devices with \( \theta \) of 1.65°, 2.24° and 2.3°, respectively, were studied — all exhibiting similar transport characteristics (see Supplemental Material\cite{supplementary} for the angle determination procedure).

A notable feature of SA-TBG is that, by employing a single- or dual-gated device architecture, one can selectively populate the minivalleys by appropriately tuning the top and bottom gate voltages\cite{supplementary} (\( V_{bg} \) and \( V_{tg} \) respectively). This combination defines the relative displacement field between graphene layers, \( D = (C_{bg}V_{bg} - C_{tg}V_{tg})/2 \), and total carrier density, \( n = (C_{bg}V_{bg} + C_{tg}V_{tg})/e \). Here \( C_{tg,bg} \) are the top and bottom gate capacitances per unit area, and \( e \) is the electron charge. Figure 1d shows the calculated band structure of the 1.65° SA-TBG for the case of zero and finite \( D \), which clearly demonstrates the gate-induced imbalance in the population between the \( k_m \) and \( k'_m \) points in the latter case. Intuitively, as the minivalleys are predominantly formed from the energy bands of different graphene sheets, an applied electric field dopes the lay-
ers unequally resulting in such an imbalance. In the presence of a perpendicular magnetic field, the gate-induced imbalance also determines the relative offset of the Landau levels (LLs) hosted by each minivalley, a property that brings us a reliable method to explore the effects of interminivalley electron scattering in SA-TBG, as we now proceed to show.

Figure 1 compares the magnetoresistance of one of our devices measured at \( T = 16 \) K, for zero and finite \( D/\varepsilon_0 = 0.35 \) V/nm at the same total \( n \) (where \( \varepsilon_0 \) is the vacuum permittivity). The carrier density was verified via the Hall effect measurements presented in Fig. 1a, which shows that the Hall resistance, \( R_{xy} \), and, therefore, \( n \), is identical for both \( D \) values. At zero \( D \), \( \rho_{xx} \) grows with increasing \( B \) but remains featureless: Shubnikov-de-Haas oscillations (SdHO) in this device disappear at 15 K for this \( B \) range (see below). In striking contrast, a clear oscillatory pattern develops in \( \rho_{xx}(B) \) data when a finite \( D/\varepsilon_0 = 0.35 \) V/nm is applied across the graphene layers. The oscillations are even more visible in the derivative of the resistivity with respect to \( B \), \( d\rho_{xx}/dB \), (inset in Fig. 1f) because of the eliminated magnetoresistance background.

Figure 2 details our observations further by comparing the low-field magnetoresistance of another SA-TBG device (2.24°) at two characteristic \( T \) and \( D \neq 0 \). At \( T = 4.2 \) K, \( \rho_{xx} \) exhibits the \( 1/B \)-periodic pattern ascribed to SdHO. Because the applied displacement field creates a small difference in the size of the Fermi surfaces associated with the \( k_m \) and \( k'_m \) minivalleys (see Fig. 1i), two oscillations of slightly different frequency emerge. The sum of these produces a familiar beating pattern. At \( T > 20 \) K, a different oscillation series, characterized by a much lower frequency, dominates the \( \rho_{xx}(B) \) behavior. In Fig. 2b we plot the amplitude, \( \Delta \rho_{xx} \), of these oscillations as a function of the inverse magnetic field \( 1/B \) and demonstrate their \( 1/B \)-periodicity. This periodicity is further verified by plotting the oscillations’ extrema indices, \( N \), against the values of the inverse magnetic field, \( 1/B_N \), at which they appear: all peaks (dips) fall onto straight lines, the slope of which defines the oscillation frequency as \( B_0 = \frac{1}{2\pi} \left( \frac{d(1/B_N)}{dN} \right)^{-1} \) (inset of Fig. 2b).

Having revealed the \( 1/B \)—character of the high-\( T \) resistance oscillations, it is instructive to explore these magnetoresonance oscillations in SA-TBG by FFT analysis. As expected from the beating pattern, the FFT spectrum at \( T = 4.2 \) K is dominated by two closely-spaced peaks, labeled as \( B_1 \) and \( B_2 \) (Fig. 2c), containing information on the carrier density in each minivalley, \( n_{1,2} \), via \( B_{1,2} = n_{1,2}\hbar/e \), where \( \hbar \) is Planck’s constant and \( g = 4 \) is the minivalley degeneracy. At \( T = 20 \) K, the FFT spectrum consists of a single peak at \( B_0 \) (labeled accordingly) that matches the periodicity determined from the \( 1/B_N(N) \) fit (inset of Fig. 2c). Interestingly, the \( B_0 \) peak is also visible at \( T = 4.2 \) K and 11 K in the FFT.
spectra, but, because of the complicated beating pattern in \( \rho_{xx}(B) \), these oscillations were obscured in previous magnetotransport studies on SA-TBG, whereas in large \( \theta > 3^{\circ} \) non-encapsulated devices, they were presumably absent\(^{21}\). Critically, we find that the obtained \( B_0 \) is identical to the difference \( B_2 - B_1 \) indicating that the period of the high-\( T \) magnetooscillations is controlled by the carrier density imbalance, \( \Delta n = n_2 - n_1 \), between the minivalleys (see below).

Figure 3 shows the \( d\rho_{xx}/dB_0 \) for \( \theta = 2.24^{\circ} \) mapped onto a \( (B,T) \) plane. Such representation allows for a convenient illustration of the evolution of magnetooscillation patterns in SA-TBG as a function of \( T \); fast SdHO, clearly visible at liquid helium \( T \), vanish at \( \sim 15 \) K whereas the slow high-\( T \) oscillations persist even above 50 K. We also studied the effect of in-plane dc current, \( I_{dc} \), on magnetoresistance and found that, in contrast to SdHO, which are readily damped by the application of only \( I_{dc} \approx 10 \) \( \mu \)A because of Joule heating, the amplitude of the high-\( T \) magnetooscillations is resilient to \( I_{dc} \), up to at least 75 \( \mu \)A. Interestingly, we also observed that upon increasing \( I_{dc} \), the phase of these oscillations flips several times, additionally distinguishing them from SdHO (see below and Supplemental Material\(^{13}\)).

Taken together, the high-\( T \) character, peculiar frequency, and fragile phase identify these oscillations as an SA-TBG analogue of magneto-intersubband oscillations (MISO) discovered in wide quantum wells (QW) and studied in related systems.\(^{23,31}\) In QW, the oscillations emerge when a two-dimensional electron system (2DES) occupies two or more energy bands capable of electron exchange.\(^{23,24,31}\) In particular, when the LLs from different subbands become aligned within the thermal window around the Fermi level, elastic interband scattering gives rise to excess resistivity. In the opposite case, when the subbands are misaligned, interband scattering is suppressed. As a result, the resistance experiences \( 1/B \)-periodic oscillations with a period proportional to the difference in filling factors between the two subbands. In the assumption that the intraband scattering time, \( \tau \), does not depend on the subband index, the oscillations’ functional form in the limit of small \( I_{dc} \) reads\(^{20,27,31}\)

\[
\Delta \rho = \frac{2\tau}{\tau_{\text{inter}}} \rho_0 \delta_1 \delta_2 \cos(2\pi \Delta \nu/g).
\]  

Here \( \rho_0 \) is the Drude resistivity, \( \delta_{1,2} = \exp(-\pi/\omega_c \tau_{q_{1,2}}) \) are the Dingle factors of the two subbands labeled by indices 1 and 2 and expressed in terms of the cyclotron frequency, \( \omega_c \), and the quantum scattering times \( \tau_{q_1} \) and \( \tau_{q_2} \). \( g \) is the subband degeneracy, \( \Delta \nu = \nu_2 - \nu_1 \) is the difference in filling factors, \( \nu_{1,2} = n_{1,2}h/Be \), of the subbands, and \( \tau_{\text{inter}} \) is the interband scattering time. Note, while initially derived for 2DES with parabolic spectrum, Eq. (1) becomes generally applicable when expressed in terms of the filling factors. Indeed, \( \Delta \nu/g \) universally determines the condition where the LLs in both subbands are aligned. In addition, we mention that the conditions of our experiments actually correspond to high filling factors and low \( T \) where the effects of non-parabolicity and associated non-equidistant LL spectrum are negligible.
To validate the interpretation of the observed high-$T$ magnetooscillations in SA-TBG in the context of MISO physics, we plot the experimentally determined $B_0$ (from Fig. 2b) as a function of $\Delta n$ in the inset of Fig. 2. This difference in carrier density, $\Delta n$, was obtained by a simple electrostatics argument that accounts for the partial screening of the applied field by the graphene layer (Supplemental Material)\textsuperscript{13}. Additionally, for some $D$, we also verified the aforementioned $\Delta n$ values by FFT analysis at liquid helium $T$ as well, where the beating of SdHO can be used to determine $\Delta n$. For all our devices, the obtained $B_0(\Delta n)$ dependence was found to be linear over a wide range of $\Delta n$ and accurately followed the functional form $B_0 = h\Delta n/4e$, where $4$ represents the degeneracy of each minivalley. This substantiates the interpretation of the observed oscillations in terms of MISO, where minivalleys now take on the role of the subbands.

Another important characteristic of MISO is its fragile phase with respect to dc bias\textsuperscript{33,35}. In the presence of magnetic field, an electric current, $I_{dc}$, of high density generates a substantial Hall field perpendicular to the current flow that initiates additional impurity-assisted tunneling of electrons between the tilted LLs\textsuperscript{37,38}. The probability of such tunneling events oscillates with magnetic field and is maximized when the Hall voltage drop across the cyclotron diameter matches an integer multiple of the cyclotron energy\textsuperscript{33}. This leads to the modification of the resonant condition for MISO which manifests itself in multiple phase reversals upon ramping $I_{dc}$. This interesting behaviour was also found in our SA-TBG devices, which exhibited the aforementioned phase flips with respect to $I_{dc}$ (Supplemental Material and Fig. S2\textsuperscript{13}) further supporting the origin of the observed oscillations.

Moreover, unlike SdHO, which also emerge as a result of the Landau quantization, the intersubband oscillations are not sensitive to the smearing of the Fermi distribution, and therefore are damped only through the broadening of LLs, parameterized via the Dingle factors in Eq. 1.\textsuperscript{13} Our data reveals this expected behaviour too: namely, the FFT amplitude of the interminivalley oscillations features a slow $\text{exp}(−\gamma T^2)$ decay ($\gamma = 11.5 \times 10^{-4}$ K$^2$), as compared to the relatively fast SdHO thermal damping governed by the conventional Lifshitz-Kosevich (LK) law (dashed purple line in Fig. 3). This behaviour is also consistent with the robustness of the interminivalley oscillations to heating induced by large $I_{dc}$ (See Supplemental Material and Fig. S2 for details\textsuperscript{13}).

The observed high-$T$ magnetooscillations provide a convenient tool to estimate the relative ratio between inter- and intraminivalley scattering rates in SA-TBG by fitting them with Eq. 1. From the exponential damping of the oscillations’ amplitude with decreasing $B$, one can extract the quantum scattering time, while $\rho_0$ can be obtained from the zero-$B$ data leaving $\tau/\tau_{\text{inter}}$ as the only fitting parameter. We have performed such an analysis for our smallest angle device and, from the data shown in Fig. 4, found that at $T = 16$ K, $\tau$ and $\tau_{\text{inter}}$ are comparable, indicating the significance of interminivalley scattering processes at small $\theta$ (see Supplemental Material and Fig. S3 for details\textsuperscript{13}). We also note a drop in the oscillations’ amplitude with increasing $\theta$. This indicates the suppression of interminivalley scattering at larger twist angles likely because a larger momentum gain, $\sim \Delta k$, is required to initiate such transitions. However, more accurate comparison on samples with identical quality is
enhancement to the reduced
the observed exp(−γT^2) behaviour of the FFT amplit
tude (Fig. 3f) suggests LL broadening induced by e-e scattering[24] and thus, its rate, 1/τ_{ee}, can be con
evienently estimated via an analysis of the oscillations’
thermal damping[23]. Assuming that this thermal damp
ning is solely encoded in δ₁δ₂ through the temperature
dependence of the quantum scattering times (see Eq. [1]),
and that these are identical in both minivalleys (a
reasonable assumption when Δn ≪ n),[27] one obtains the
T−dependent amplitude of the interminivalley oscilla
tions: δ₁δ₂ = e^{-2π/ω−τ_e(T)}. Since τ_{ee}^{-1}(T) = τ_{0}^{-1} + τ_{ee}^{-1}(T), where τ_{0} is the T−independent elastic quantum
scattering time, one can extract τ_{ee}^{-1} from the FFT mag
nitude of the interminivalley oscillations. Figure 3 shows
the results of such an analysis and plots the τ_{ee}^{-1}(T) de
pendence. For θ = 2.24°, we find that the obtained esti
mates exceed the e-e scattering rate in MLG (blue dashed
curve in Fig. 3c) at identical n.[32,41,49] We attribute this
enhancement to the reduced v_e = 0.75v_0 in the SA-TBG
of this θ as compared to that of MLG, v_0 = 10^6 m/s. In
deed, by renormalizing the 1/τ_{ee}(T) dependence for MLG
from Fig. [3]-by the ratio of the Fermi velocities in these
systems and accounting for the two fold increase in the
degeneracy of SA-TBG as compared to MLG, we obtain the
scattering rate for SA-TBG close to that found experimen
tally.

Our experiments also raise important questions about
scattering processes in twisted moiré systems. The
observed high-T oscillations indicate the presence of
some scattering mechanism(s) enabling electrons to gain
enough momentum (∼ Δk) to escape from their mini
valley and scatter to another one (k_{m} ↔ k'_{m}). This, in
turn, may imply the presence of scatterers with a spa
tial scale of the order of 1/Δk ∼ λ_m, where λ_m is the
superlattice period. A possible candidate is twist an
gle disorder, regularly observed in devices fabricated by
the methods used here.[23] An alternative scenario involves
acoustic phonon-assisted processes.[44,50] However, in our
data, the interminivalley oscillations are present even at
liquid helium T (Fig. 3f), at which the allowed phase
space for phonon momenta is not sufficient to ensure the
momentum mismatch Δk. At T = 4.2 K, phonons with
momenta q < k_BT/ℏs ≈ 4 × 10^7 m⁻¹ are populated
(where s ≈ 20 km/s is the characteristic speed of sound
in graphene), i.e., those having momenta over an order of
magnitude smaller than Δk at the studied twist an
gles. The only phonon branch that at such low T can
be populated up to the required momenta, is the breath
ing mode,[27] however, little is known on its impact on
SA-TBG resistivity.[45]

To conclude, we have observed high-T magnetoooscil
lations in SA-TBG when a finite displacement field is
applied across the graphene layers. Although similarly

periodic in 1/B, these oscillations show a clearly distinct
temperature and dc current dependence from SdHO and
are controlled by the difference in the minivalleys’ filling
factors. Drawing a parallel with MISO, we have shown
that the observed oscillations originate from interminival
ley scattering, allowed by the reduced size of the mini
Brillouin zone in SA-TBG. By analyzing the amplitude
of these high-T oscillations, we estimated the relative ra
tio between interminivalley and intraminivalley scatter
ing times, τ/τ_{inter}, which we found to be of similar order
in the θ = 1.65° device. Finally, from the temperature
dependence of the oscillations, we obtained information
on τ_{ee}^{-1}, which we found to exceed that of MLG due to
a reduced v_e in this system. Our study points to the
presence of a scattering mechanism(s) of unknown na
ture with large momentum transfer and highlights the
importance of interminivalley momentum relaxation in
the resistivity of twisted moiré system,[45] which has to be
accounted for in future studies.

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SUPPLEMENTAL MATERIAL

Supplementary Section 1. Device fabrication

Our devices consisted of hBN-encapsulated twisted bilayer graphene, which we fabricated using a combination of cut-and-stack and hot release methods. Monolayer graphene, few-layer graphite, and 30-80 nm-thick hBN crystals were mechanically exfoliated on a Si/SiO\textsubscript{2} substrate, and sizable, uniform flakes were selected using optical contrast. Then, using a homemade transfer system with \( \mu \text{m} \)-accuracy and a polycarbonate (PC) membrane stretched over a small (8 mm\( \times \)8 mm\( \times \)4 mm) polydimethylsiloxane (PDMS) polymer block on a glass slide, we assembled hBN and graphite stacks on a Si/SiO\textsubscript{2} wafer. To minimize strain on the hBN, we picked up at 50-70 \( ^\circ \)C, when the membrane was minimally sticky enough to allow for a clean pickup. The graphite was picked up at room temperature, and then the entire stack was “ironed” and then released on a clean Si/SiO\textsubscript{2} wafer at high temperatures (160 – 170 \( ^\circ \)C). After removing the polymer membrane, we annealed the hBN and graphite stack at 350 \( ^\circ \)C for 3 hours while flowing argon and hydrogen in order to ensure the removal of any residues. We then assembled the hBN and twisted bilayer graphene stack using a “cut-and-stack” method described previously. After picking up the top hBN and twisted graphene, we “ironed” the entire stack at room temperature. The three-layer stack was then released onto the previously fabricated and cleaned bottom hBN and graphite gate at roughly 160 \( ^\circ \)C. After this point, we avoided heating the stack to reduce the possibility of twist angle relaxation. The resulting heterostructure is shown in Fig. S1a. The final stack was inspected using dark-field microscopy (Fig. S1b) and atomic force microscopy (AFM), and bubble- and blister-free areas were selected to use for Hall bars (Fig. S1c).

To fabricate the devices, we covered the heterostructures by a protective polymethyl-methacrylate (PMMA) resist and used electron beam lithography (EBL) to define contact regions (Fig. S1d). We then performed a mild O\textsubscript{2} plasma cleaning before using reactive ion etching (RIE) with a plasma generated from CHF\textsubscript{3} and O\textsubscript{2} gases to selectively etch away the hBN in the parts of the heterostructure unprotected by the lithographic mask. 3 nm chromium and 50-70 nm gold was then evaporated into the contact regions via thermal evaporation at high vacuum (Fig. S1e). We repeat the same EBL and thermal evaporation procedures to define a metallic top gate (3 nm chromium and 30-40 nm gold). Finally, we repeat the same EBL and RIE procedures to define the final Hall bar geometry, using, in this case, a plasma generated by Ar, O\textsubscript{2} and CHF\textsubscript{3} gases.

FIG. S1. Device fabrication a, Optical photograph of a twisted bilayer graphene encapsulated between two slabs of hBN. b, Dark field image of the obtained heterostructure with blister-free area that was further examined by AFM (red). c, AFM topography of regions used for Hall bars. Ripples are an artifact of the AFM scanning and are not physical. White bar is 5 \( \mu \text{m} \). Area used for Hall bar indicated by dashed black rectangle. d, Protective PMMA mask used to define contact areas via RIE. e, Gold contacts are evaporated on to graphene after RIE. f, Optical photograph of the finalized device.
Supplementary Section 2. Angle determination

Because of the relatively large $\theta$ for two of our devices, we were limited by the robustness of the gate dielectric and could not reach full-filling of the moiré bands, which is the standard approach of determining twist angle. To circumvent this problem, we performed further magnetoresistance measurements and, for one of the devices, resolved well-defined Brown-Zak oscillations, whose fundamental frequency provided an accurate tool for determining the superlattice period and therefore the corresponding twist angle $2.24^\circ$. For the second device, where Brown-Zak oscillations were not detected, we performed a detailed measurements of the effective mass of charge carriers, $m$, using $T$-induced damping of the SdHO and compared it with that of the $2.24^\circ$ SA-TBG. This comparison pointed to 10–20% lighter charge carriers in the $2.24^\circ$ SA-TBG at the same $n$ which provided a rough estimate for the twist angle $2.3^\circ$ using a theoretical framework. Lastly, from the superlattice-induced insulating states at full-filling ($n_s$) readily observed in the $\rho_{xx}(n)$ dependencies of our smallest-angle SA-TBG (Fig. 1c) and the relation $n_s = 4 \times \frac{8 \sin^2(\theta/2)}{\sqrt{3}a^2}$, we determine $\theta \approx 1.65^\circ$. 
Supplementary Section 3. Connection between the applied gate voltages and density imbalance

As discussed in the main text, to validate the interpretation of the observed high-$T$ magnetoooscillations in SA-TBG in the context of MISO physics, one has to know the difference in carrier density $\Delta n$ between the two minivalleys for a given displacement field $D$ applied between the graphene layers. This difference can obtained by a simple electrostatics consideration that accounts for the electrostatic screening effect:

\[
\frac{e d}{\epsilon_0 \epsilon} \left( \epsilon_0 D - \frac{1 + \epsilon}{4} (n_1 - n_2) e \right) = \frac{\hbar v_0}{2 \sqrt{\pi}} \left( s_1 \sqrt{|n_1|} - s_2 \sqrt{|n_2|} \right), \quad (S2)
\]

\[
n = n_1 + n_2, \quad (S3)
\]

where $d$ is the distance separating the graphene layers, $e$ is the electron charge, $\epsilon$ is the dielectric constant for twisted bilayer, $v_0$ is the Dirac velocity and $\epsilon_0$ is the vacuum permittivity; the band indices are given by $s_i = n_i/|n_i|$. To find $\Delta n$ for each value of $n$ and $D$ used in the experiment, the two equations are solved simultaneously using $\epsilon = 2.7$ known from previous experiments.
Supplementary Section 4. Interminivalley oscillations under strong dc bias

To further explore the properties of the observed high-\(T\) interminivalley oscillations, we have performed the measurements of the differential resistance, \(r = \frac{dV}{dI}\), as a function of DC current \(I_{\text{dc}}\) and magnetic field, \(B\), in our \(\theta = 2.3^\circ\) device. Figure S2 shows the \(r(B)\) dependence along with its derivative \(\frac{dr}{dB}\) (red) measured at \(I_{\text{dc}} = 0 \ \mu\text{A}\) (black) and \(T = 20 \ \text{K}\). The data resembles that shown in Fig. 2a of the main text. The differentiation improves the visibility of the oscillations by removing the smooth non-oscillating \(I_{\text{dc}}\)-dependent background. With increasing \(I_{\text{dc}}\), the amplitude of the oscillations remains practically unaffected whereas the phase changes for certain values of \(I_{\text{dc}}\) as apparent from the checkboard pattern of the \(\frac{dr}{dB}(I_{\text{dc}}, B)\) map shown in Fig. S2a. This pattern persists up to the highest currents of \(I_{\text{dc}} \approx 75 \ \mu\text{A}\) applied in our experiments with no apparent decay of the oscillations’ amplitude. For comparison, SdHO are washed out with the application of only \(I_{\text{dc}} \sim 10 \ \mu\text{A}\) if similar measurements are performed at \(T = 4.2 \ \text{K}\). These observations clearly indicate that the observed high-\(T\) magnetoooscillations are entirely different than SdHO.

Qualitatively, in the presence of magnetic field, an electric current of high density generates a substantial Hall field perpendicular to the current flow. In sufficiently clean systems, this field initiates scattering-assisted transitions of electrons between the tilted Landau levels\(^{37–39}\). The probability of these transitions is maximized when the Hall voltage drop across the cyclotron diameter matches an integer multiple of the cyclotron energy\(^{31}\) giving rise to peculiar magnetoooscillations sensitive to the value of \(I_{\text{dc}}\). The observed phase flip and the corresponding checkboard pattern shown in Fig. S2a demonstrates the modification of the intersubband scattering in SA-TBG by Hall voltage-induced transitions\(^{33–36}\).

**FIG. S2.** Interminivalley oscillations in the presence of strong dc field. **a**, The derivative of the differential resistance with respect to magnetic field, \(\frac{dr}{dB}\), mapped against \(B\) and \(I_{\text{dc}}\). **b**, The interminivalley oscillations observed in \(r(B)\) (black) and its derivative \(\frac{dr}{dB}\) at \(I_{\text{dc}} = 0 \ \mu\text{A}\). \(\theta = 2.3^\circ, T = 20 \ \text{K}, n = 1.5 \times 10^{12} \ \text{cm}^{-2}\).
Supplementary Section 5: Estimating $\tau/\tau_{\text{inter}}$ in SA-TBG

The observed interminivalley oscillations reported in the main text provide a tool for characterizing momentum relaxation processes in SA-TBG. Namely, they enable estimating the ratio of the interminivalley scattering to the total scattering rate. Figure S3 plots the amplitude of the interminivalley oscillations, $\rho_{xx}$, as a function of inverse magnetic field, $1/B$, obtained by subtracting a smooth, non-oscillating background from the data reported in Fig. 1f of the main text. The oscillations feature a damped sinusoidal dependence on $1/B$ described by Eq. (1) of the main text. In the assumption of equal Dingle factors in both minivalleys $\delta_1 = \delta_2$ (a fair assumption as long as $\delta n \ll n$), this equation becomes:

$$\Delta \rho = \frac{2\tau}{\tau_{\text{inter}}} \rho_0 e^{-\frac{2\pi m/e B\tau_q}{B}} \cos \left(2\pi B_0/B\right),$$

where $\rho_0$ is the resistivity at $B = 0$, $m$ is the effective mass of the charge carriers, $\tau_q$ is the quantum scattering time, and $\tau/\tau_{\text{inter}}$ is the ratio between the intraminivalley and interminivalley scattering times. Using $m$, obtained from the analysis of the SdHO thermal damping, $\tau_q$ from the Dingle plot shown in the inset of Fig. S3 and $\rho_0$ from the zero-$B$ measurements, we plot the calculated $\Delta \rho(1/B)$ dependence which shows good agreement with experiment for $\tau/\tau_{\text{inter}} = 3/4$. For comparison, we also illustrate the case when interminivalley scattering is weaker, $\tau/\tau_{\text{inter}} = 1/3$, which shows a reduced oscillations amplitude. These findings indicate the importance of interminivalley scattering in SA-TBG.

![Figure S3](image_url)

**FIG. S3.** Estimating $\tau/\tau_{\text{inter}}$. Solid: Oscillation amplitude, $\Delta \rho_{xx}$, as a function of $1/B$, obtained by subtracting a smooth non-oscillating background from the data in Fig. 1f (red curve). Red and blue dashed lines: intersubband resistivity oscillations calculated from Eq. (1) for $\tau/\tau_{\text{inter}} = 3/4$ and $\tau/\tau_{\text{inter}} = 1/3$ and experimentally determined $\rho_0 = 122 \, \Omega$, $m = 0.044 m_e$, and $\tau_q = 0.6 \, \text{ps}$. Inset: Dingle plot for the interminivalley oscillations in $\theta = 1.65^\circ$ SA-TBG along with the best fit to an exponential decay function (black dashed line) yielding $\tau_q = 0.6 \, \text{ps}$. 
