Stacking-order dependence in thermoelectric transport of biased trilayer graphene

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We numerically study the thermoelectric and thermal transport in trilayer graphene with different stacking orders in the presence of interlayer bias under a strong perpendicular magnetic field. In the biased ABA-stacked case, we find that the thermoelectric conductivity displays different asymptotic behaviors with the varying of the temperature, similar to that of monolayer graphene. In the high-temperature regime, the transverse thermoelectric conductivity as saturates to a universal value 2.77k_B/e/h at the center of each Landau level, while it displays a linear temperature dependence at the low-temperature limit. The calculated transverse thermal conductivity as exhibits two plateaus away from the band center. The transition between the two plateaus is continuous, which is accompanied by a pronounced peak in the longitudinal thermal conductivity at. In the biased ABC-stacked case, it is found that both the thermoelectric conductivity and thermal conductivity have similar properties to the biased bilayer graphene, which is consistent with the behavior of a band insulator. The obtained results demonstrate the sensitivity of the thermoelectric conductivity to the band gap near the Dirac point. We also verify the validity of the Mott relation and the generalized Wiedemann-Franz law.

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I. INTRODUCTION

Recently, much attention has been paid to the thermoelectric transport properties of graphene both experimentally1–3 and theoretically.4–8 In experiments, the thermopower (the longitudinal thermoelectric response) and the Nernst signal (the transverse response) in the presence of a strong magnetic field are found to be large, reaching the order of the quantum limit k_B/e, where k_B and e are the Boltzmann constant and the electron charge, respectively.1–3 Besides monolayer graphene, bilayer graphene is also very interesting. Experiments have shown that bilayer graphene exhibits tunable bandgap9,10 in the presence of an applied bias voltage, exhibiting similar properties to semiconductors. Thermoelectric measurement11 shows that the room-temperature thermopower with a bias voltage can be enhanced by a factor of four compared to that of monolayer graphene. Theoretical calculations from the tight-binding models and band structure with a tunable band overlap between the hole bands, which are quite different from those of biased bilayer graphene.23,25,26,33 Owing to their distinctive band structures, ABA- and ABC-stacked trilayer graphene are expected to exhibit rich novel thermoelectric transport properties. However, theoretical understanding of the thermoelectric transport properties of trilayer graphene is limited compared to that of monolayer or bilayer graphene. In particular, the influence of different stacking orders on the thermoelectric transport properties has not been studied so far, which is highly desired.

In this paper, we carry out a numerical study of the thermoelectric transport properties in both ABA- and ABC-stacked trilayer graphene systems in the presence of electrostatic bias between the top and bottom graphene layers. We focus on studying the effects of disorder and thermal activation on the broadening of LLs and the corresponding thermoelectric transport properties. In the biased ABA-stacked case, the thermoelectric coefficients exhibit unique characteristics near the central LL due to the LL crossing of electron and hole bands, which are quite different from those of biased bilayer graphene. Both the longitudinal and the transverse band structure for ABA-stacked trilayer graphene contains both linear and hyperbolic bands, similar to the combined spectrum of monolayer graphene and bilayer graphene.27,28 While the ABC-stacked case presents approximately cubic bands.25 Moreover, the LL spectrum of the ABA-stacked case in a perpendicular magnetic field B can be viewed as a superposition of B-dependent monolayerlike LLs and B-dependent bilayerlike LLs.23,29,30 On the other hand, the LLs of the ABC-stacked case are given by E_n ∝ B^{3/2}√n(n-1)(n-2) with Berry’s phase 3π.53–25 Interestingly, when a bias voltage or a potential difference is applied to the top and bottom graphene layers, the ABA-stacked case exhibits a semimetallic band structure with a tunable band overlap between the conduction and valence bands.31,32 whereas the ABC-stacked case exhibits a semiconducting band structure with a tunable band gap, similarly to bilayer graphene.23,25,26,33
thermoelectric conductivities are universal functions of the effective bandwidth and temperature, and display different asymptotic behaviors in different temperature regimes. The Nernst signal displays an apparent plateau with $\gamma_0$, and changes sign near other LLs, while the thermopower behaves in an opposite manner. The peak values of the Nernst signal and thermopower are very large, compared with monolayer graphene due to the semimetallic band overlap near zero energy. The validity of the semiclassical Mott relation is found to remain valid at low temperatures. In the biased ABC-stacked case, we observe quite different behavior from the biased ABA-stacked case near the central LL. Around the Dirac point, the transverse thermoelectric conductivity displays an apparent plateau with $\gamma_0 = 3.16$ eV, $\gamma_1 = 0.39$ eV, and $\gamma_3 = 0.315$ eV.

We assume that each monolayer graphene has totally $L_\gamma$ zigzag chains with $L_z$ atomic sites on each zigzag chain. The size of the sample will be denoted as $N = L_x \times L_y \times L_z$, where $L_z = 3$ is the number of monolayer graphene planes along the $z$ direction. We have confirmed that the calculated results do not depend on the system sizes (as long as the system lengths are reasonably large). In the presence of an applied magnetic field perpendicular to the plane of biased trilayer graphene, the lattice model for ABA stacking in real space can be written in the tight-binding form,

$$H = -\gamma_0 \left( \sum_{\langle ij \rangle \sigma} e^{i \alpha_{ij}} c^\dagger_{i \sigma} c_{j \sigma} + \sum_{\langle ij \rangle \sigma} e^{i \alpha_{ij}} c^\dagger_{3 \sigma} c_{3 \sigma} \right) - \gamma_1 \left( \sum_{\langle ij \rangle \sigma} e^{i \alpha_{ij}} c^\dagger_{1 \sigma} B_1 c_{2 \sigma} A_1 + \sum_{\langle ij \rangle \sigma} e^{i \alpha_{ij}} c^\dagger_{2 \sigma} B_2 c_{3 \sigma} A_2 + \sum_{\langle ij \rangle \sigma} e^{i \alpha_{ij}} c^\dagger_{2 \sigma} B_2 c_{3 \sigma} A_3 \right) + \text{H.c.} + \sum_{\sigma} w_i c^\dagger_{i \sigma} c_{i \sigma} + c^\dagger_{2 \sigma} c_{2 \sigma} + c^\dagger_{3 \sigma} c_{3 \sigma} + \sum_{i \sigma} (e_1 c^\dagger_{i \sigma} c_{i \sigma} + e_2 c^\dagger_{3 \sigma} c_{3 \sigma}),$$

(1)

where $c^\dagger_{m \sigma}$ ($c_{m \sigma}$) are creation operators on sublattices $A_m$ and $B_m$ in the $m$th layer ($m = 1, \ldots, 3$), with $\sigma$ as a spin index. The sum $\sum_{\langle ij \rangle \sigma}$ denotes the intralayer nearest-neighbor hopping in three layers, $\sum_{\langle ij \rangle \sigma}$ stands for the interlayer hopping between the $B_1$ (B3) sublattice in the bottom (top) layer and the $A_2$ sublattice in the middle layer, and $\sum_{\langle ij \rangle \sigma}$ stands for the interlayer hopping between the $B_2$ sublattice in the middle layer and the $A_1$ (A3) sublattice in the bottom (top) layer, as described above. For the biased system, the top and bottom graphene layers gain different electrostatic potentials, and the corresponding energy difference is given by $\Delta_x = \epsilon_2 - \epsilon_1$ where $\epsilon_1 = -\frac{1}{2} \Delta_x$, and $\epsilon_2 = \frac{1}{2} \Delta_x$. For illustrative purpose,
a relatively large asymmetric gap $\Delta \gamma = 0.3\gamma_0$ is assumed. $w_i$ is a random disorder potential uniformly distributed in the interval $w_i \in [-W/2, W/2] \gamma_0$. The magnetic flux per hexagon $\phi = \sum a_j = \frac{2\pi}{\sqrt{3}}$ is proportional to the strength of the applied magnetic field $B$, where $M$ is assumed to be an integer and the lattice constant is taken to be unity.

For ABC-stacked trilayer graphene in the presence of bias voltage, the Hamiltonian can be written as

$$H = -\gamma_0 \left( \sum_{(ij)\sigma} e^{i\alpha_{ij}} c^+_i \sigma c_j \sigma + \sum_{(ij)\sigma} e^{i\alpha_{ij}} c^+_i 2\sigma c_j 2\sigma \right) + \sum_{(ij)\sigma} e^{i\alpha_{ij}} c^+_i 3\sigma c_j 3\sigma - \gamma_1 \left( \sum_{(ij)\sigma} e^{i\alpha_{ij}} c^+_i 1\sigma B_1 c_{2\sigma} A_1 \right) + \sum_{(ij)\sigma} e^{i\alpha_{ij}} c^+_i 2\sigma B_2 c_{3\sigma} A_3 \right) + \sum_{(ij)\sigma} e^{i\alpha_{ij}} c^+_i 3\sigma B_3 c_{1\sigma} A_1 \right) + \text{H.c.} + \sum_{i\sigma} w_i (c^+_i 1\sigma c_{1\sigma} + c^+_i 2\sigma c_{3\sigma} + c^+_i 3\sigma c_{1\sigma}) \right) + \sum_{i\sigma} (\epsilon_1 c^+_i 1\sigma c_{1\sigma} + \epsilon_2 c^+_i 3\sigma c_{3\sigma}) \right).$$

The sum $\sum_{(ij)\sigma}$ denotes the intralayer nearest-neighbor hopping in three layers, $\sum_{(ij)\sigma}$ stands for the interlayer hopping between the $B_1$ ($B_2$) sublattice in the bottom (middle) layer

and the $A_2$ ($A_3$) sublattice in the middle (top) layer, and $\sum_{i\sigma}$ stands for the interlayer hopping between the $B_2$ ($B_3$) sublattice in the middle (top) layer and the $A_1$ ($A_2$) sublattice in the bottom (middle) layer, as described above.

In the linear response regime, the charge current in response to an electric field and a temperature gradient can be written as $\mathbf{J} = \mathbf{\sigma} \mathbf{E} + \alpha (-\nabla T)$, where $\mathbf{\sigma}$ and $\alpha$ are the electrical and thermoelectric conductivity tensors, respectively. The transport coefficient $\sigma_{\alpha\beta}$ can be calculated by Kubo formula $\sigma_{\alpha\beta}$ and $\sigma_{xx}$ can be obtained based on the calculation of the Thouless number. We exactly diagonalize the tight-binding Hamiltonian in the presence of disorder. Then the transport coefficients can be calculated using the obtained energy spectra and wave functions. In practice, we first calculate the $T = 0$ conductivities $\sigma_{\beta\beta}(E_F)$, and then use the relation to obtain the finite-temperature electrical and thermoelectric conductivity tensors. Here, $f(x) = 1/[e^{(x-E_F)/k_B T} + 1]$ is the Fermi distribution function. At low temperatures, the second equation can be approximated as

$$\alpha_{\beta\beta}(E_F, T) = -\frac{\pi^2 k_B^2 T}{3e} \frac{d\sigma_{\beta\beta}(E_F, T)}{dE} \bigg|_{E=E_F},$$

FIG. 2. (Color online) Thermoelectric conductivities at finite temperatures of biased ABA-stacked trilayer graphene. (a), (b) $\alpha_{xx}(E_F, T)$ and $\alpha_{xy}(E_F, T)$ as functions of the Fermi energy $E_F$ at different temperature $T$. (c) shows the temperature dependence of $\alpha_{xx}(E_F, T)$ for trilayer graphene. (d) Comparison of the results from numerical calculations and from the generalized Mott relation at two characteristic temperatures, $k_B T/W_L = 0.1$ and $k_B T/W_L = 1.5$. Here, the width of the central LL $W_L/\gamma_0 = 0.0069$. The asymmetric gap $\Delta = 0.3\gamma_0$. The system size is taken to be $N = 96 \times 24 \times 3$, the magnetic flux $\phi = 2\pi/48$, and the disorder strength $w = 0.1$. 

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which is the semiclassical Mott relation. The thermopower and Nernst signal can be calculated subsequently from

\[
S_{xx} = \frac{E_x}{V_x T} = \rho_{xx} \alpha_{xx} - \rho_{yx} \alpha_{yx}, \\
S_{xy} = -\frac{E_y}{V_x T} = \rho_{xx} \alpha_{yx} + \rho_{yx} \alpha_{xx}.
\]

(5)

The thermal conductivity, measuring the magnitude of the thermal currents in response to an applied temperature gradient, which usually includes electron and phonon contributions. In our numerical calculations, phonon-related thermal conductivity is omitted. The electronic thermal conductivities \(\kappa_{ji}\) at finite temperature assume the forms

\[
\kappa_{ji}(E_F, T) = \frac{1}{e^2 T} \int d\epsilon \sigma_{ji}(\epsilon)(\epsilon - E_F)^2 \left( -\frac{\partial f(\epsilon)}{\partial \epsilon} \right) - T \alpha_{ji}(E_F, T) \sigma^{-1}_{ji}(E_F, T) \alpha_{ji}(E_F, T).
\]

(6)

For diffusive electronic transport in metals, it is established that the Wiedemann-Franz law is satisfied between the electrical conductivity \(\sigma\) and the thermal conductivity \(\kappa\) of electrons

\[
\frac{\kappa}{\sigma T} = L,
\]

(7)

where \(L\) is the Lorentz number and takes a constant value: \(L = \frac{\pi^2}{3} k_B^2 e^2\). The validity of this relation will be examined for the present trilayer graphene.

III. THERMOELECTRIC TRANSPORT IN BIASED ABA-STACKED TRILAYER GRAPHENE SYSTEMS

We start from numerically diagonalizing the tight-binding Hamiltonian with the presence of disorder scattering. In the biased ABA-stacked case, the density of states shows a central peak around the Dirac point, and the width of the peak increases with the increase of the disorder strength. While in the biased ABC-stacked case, the density of states exhibits a pronounced valley around the Dirac point. Due to these clear differences in the density of states for both cases, the thermoelectric transport properties exhibit different behavior.

We first show the calculated thermoelectric conductivities at finite temperatures for biased ABA-stacked trilayer graphene. As shown in Figs. 2(a) and 2(b), the transverse thermoelectric conductivity \(\alpha_{xy}\) displays a series of peaks, while the longitudinal thermoelectric conductivity \(\alpha_{xx}\) oscillates and changes sign at the center of each LL. At low temperatures, the peak of \(\alpha_{xy}\) at the central LL is higher and narrower than others, which indicates that the impurity scattering has less effect on the central LL. These results are qualitatively

![Fig. 3.](image-url)

FIG. 3. (Color online) Thermoelectric conductivities at finite temperatures of biased ABC-stacked trilayer graphene. (a), (b) \(\alpha_{xy}(E_F, T)\) and \(\alpha_{xx}(E_F, T)\) as functions of Fermi energy at different temperatures. (c) Comparison of the results from numerical calculations and from the generalized Mott relation at two characteristic temperatures, \(k_BT/\Delta_g = 0.001\) and \(k_BT/\Delta_g = 0.05\). The asymmetric gap \(\Delta_g = 0.3\gamma_0\). The other parameters are chosen to be the same as in Fig. 2.
similar to those found in monolayer graphene, but some differences exist due to the overlap of electron and hole bands. At low temperatures, more oscillations are observed in the higher LLs than monolayer graphene, in consistent with the further lifting of the LL degeneracy in the biased ABA-stacked case. As shown in Fig. 2(b), around zero energy, the peak value of $\alpha_{xx}$ shows different trends with increasing temperature. It first increases with $T$ at low-temperature region, and then decreases with $T$ at high temperatures. This is due to the competition between $\frac{\sigma_{0}^2}{2\gamma} T^{3} e^{-\frac{1}{3}}$ and $\frac{d\sigma_{ji}(\varepsilon,T)}{d\varepsilon}$ of Eq. (4). The peak value of $\alpha_{xx}$ could either increase or decrease depending on the relative magnitudes of these two terms. At high temperatures, $\sigma_{ji}(\varepsilon,T)$ becomes smooth, and consequently $\alpha_{xx}$ begins to decrease. In Fig. 2(c), we find that $\alpha_{yy}$ shows different behavior depending on the relative strength of the temperature $k_{B} T$ and the width of the central LL $W_{L}$ ($W_{L}$ is determined by the full-width at the half-maximum of the $\sigma_{xx}$ peak). When $k_{B} T \ll W_{L}$ and $E_{F} \ll W_{L}$, $\alpha_{yy}$ shows linear temperature dependence, indicating that there is a small energy range where extended states dominate, and the transport falls into the semiclassical Drude-Zener regime. When $E_{F}$ is shifted away from the Dirac point, the low-energy electron excitation is gapped due to Anderson localization. When $k_{B} T$ becomes comparable to or greater than $W_{L}$, $\alpha_{yy}$ for all LLs saturates to a constant value $2.77 k_{B} e/h$. This matches exactly the universal value (in $2 k_{B} e/h$ predicted for the conventional integer quantum Hall effect (IQHE) systems in the case where thermal activation dominates, with an additional degeneracy factor 4. The saturated value of $\alpha_{yy}$ in the biased ABA-stacked case is in accordance with the fourfold degeneracy at zero energy. In the presence of bias voltage, the valley degeneracy of the LLs usually is lifted by the interlayer potential asymmetry, so that the twelvefold energy levels (four and eight levels from the monolayerlike and the bilayerlike sub-bands, respectively) split into six different levels with twofold spin degeneracy. However, near the Dirac point, the interlayer potential asymmetry causes hybridization of the linear and parabolic chiral bands, which leads to the fourfold degeneracy for zero energy Landau levels. To examine the validity of the semiclassical Mott relation, we compare the above results with those calculated from Eq. (4), as shown in Fig. 2(d). The Mott relation is a low-temperature approximation and predicts that the thermoelectric conductivities have linear temperature dependence. This is in agreement with our low-temperature results, which proves that the semiclassical Mott relation is asymptotically valid in the Landau-quantized systems, as suggested in Ref. 38.

IV. THERMOELECTRIC TRANSPORT IN BIASED ABC-STACKED TRILAYER GRAPHENE SYSTEMS

For biased ABC-stacked trilayer graphene, we show the calculated $\alpha_{xx}$ and $\alpha_{xy}$ at finite temperatures in Fig. 3. As seen from Fig. 3(a), $\alpha_{xy}$ displays a pronounced valley at low temperature, in striking contrast to the ABA-stacked case with a peak at $E_{F} = 0$. These results are qualitatively similar to those found in biased bilayer graphene. This behavior can be understood as due to the split of the valley degeneracy in the central LL by an opposite voltage bias added to the top

![FIG. 4. (Color online) The thermopower $S_{xx}$ and the Nernst signal $S_{xy}$ as functions of the Fermi energy in (a), (b) biased ABA-stacked trilayer graphene, and (c), (d) biased ABC-stacked trilayer graphene at different temperatures. The parameters in these two systems are chosen to be the same as in Figs. 2 and 3, respectively.](115414-5)
layer and the bottom layer. This is consistent with the opening of a sizable gap between the valence and conduction bands in biased ABC-stacked trilayer graphene.\cite{note2} \(\alpha_{xx}\) oscillates and changes sign around the center of each split LL. In Fig. 3(c), we also compare the above results with those calculated from the semiclassical Mott relation using Eq. (4). The Mott relation is found to remain valid at low temperatures.

We further calculate the thermopower \(S_{xx}\) and the Nernst signal \(S_{xy}\). In Figs. 4(a) and 4(b), we show the calculated \(S_{xx}\) and \(S_{xy}\) in biased ABA-stacked trilayer graphene. As we can see, \(S_{xy}\) (\(S_{xx}\)) has a peak (peaks) at the central LL (the other LLs), and changes sign near the other LLs (the central LL). At zero energy, both \(\rho_{xy}\) and \(\alpha_{xx}\) vanish, leading to a vanishing \(S_{xy}\). Around zero energy, because \(\rho_{xx}\alpha_{xx}\) and \(\rho_{xy}\alpha_{xy}\) have opposite signs, depending on their relative magnitudes, \(S_{xx}\) could either increase or decrease when \(E_F\) is increased passing the Dirac point. In the biased ABA-stacked case, \(S_{xy}\) is dominated by \(\rho_{xx}\alpha_{xy}\), consequently, \(S_{xx}\) increases to positive value as \(E_F\) passing zero. At low temperatures, the peak value of \(S_{xx}\) near zero energy is \(\pm 0.81k_B/e\) (\(\pm 69.8\) \(\mu V/K\)) at \(k_B T = 0.4W_L\), which is in agreement with the measured value (\(\pm 70\) \(\mu V/K\)).\cite{note43} With the increase of temperature, the peak height increases to \(\pm 3.75k_B/e\) (\(\pm 323.14\) \(\mu V/K\)) at \(k_B T = 1.5W_L\). On the other hand, \(S_{xy}\) has a strong peak structure around zero energy, which is dominated by \(\rho_{xy}\alpha_{xy}\). The peak height is \(7.82k_B/e\) (673.85 \(\mu V/K\)) at \(k_B T = 1.5W_L\). The large magnitude of \(S_{xy}\) and \(S_{xx}\) near zero energy can be attributed to the semimetal-type dispersion of biased ABA-stacked trilayer graphene, and the fact that the system is in the vicinity of a quantum Hall liquid to insulator transition, where the imbalance between the particle and hole types of carriers should be significant. The thermoelectric effects are very sensitive to such an imbalance in Dirac materials in comparison with conventional metals.

In Figs. 4(c) and 4(d), we show the calculated \(S_{xx}\) and \(S_{xy}\) in the biased ABC-stacked case. As we can see, \(S_{xy}\) (\(S_{xx}\)) has a peak (peaks) around zero energy (the other LLs), and changes sign near the other LLs (zero energy). These results are qualitatively similar to those found in the biased ABA-stacked case. In our calculation, we find that \(S_{xx}\) is always dominated by \(\rho_{xx}\alpha_{xx}\), consequently, \(S_{xx}\) decreases to negative values as \(E_F\) passing zero. This is different from the biased ABA-stacked case. At low temperatures, the peak value of \(S_{xx}\) near zero energy is \(\pm 0.86k_B/e\) (\(\pm 74.11\) \(\mu V/K\)) at \(k_B T = 0.01\Delta_g\). With the increase of temperature, the peak height increases to \(\pm 2.23k_B/e\) (\(\pm 192.16\) \(\mu V/K\)) at \(k_B T = 0.1\Delta_g\). On the other hand, \(S_{xy}\) has a peak structure around zero energy, which is dominated by \(\rho_{xy}\alpha_{xy}\). The peak height is 4.69 \(k_B/e\) (404.14 \(\mu V/K\)) at \(k_B T = 0.1\Delta_g\).

V. THERMAL CONDUCTIVITY IN BIAS- AND ABC-STACKED TRILAYER GRAPHENE SYSTEMS

We now focus on thermal conductivities. In Fig. 5, we show results of the transverse thermal conductivity \(\kappa_{xy}\) and the longitudinal thermal conductivity \(\kappa_{xx}\) for biased ABA-stacked trilayer graphene at different temperatures. As seen from Figs. 5(a) and 5(b), \(\kappa_{xy}\) exhibits two flat plateaus away from the central LL. The values of the plateaus in \(\kappa_{xy}\) are \(\pm 0.0045 \gamma_0 k_B/h\) \(\pm 0.048\) \(\text{nW/(K-m)}\) at \(k_B T = 0.1W_L\). With
FIG. 6. (Color online) (a), (b) Thermal conductivities $\kappa_{xy}(E_F, T)$ and $\kappa_{xx}(E_F, T)$ as functions of the Fermi energy in biased ABC-stacked trilayer graphene at different temperatures. (c), (d) Comparison between the thermal conductivity as functions of the Fermi energy obtained from numerical calculations and from the Wiedemann-Franz law at two characteristic temperatures. The parameters chosen here are the same as in Fig. 3.

the increase of temperature, the values of the plateaus increase to $\pm 0.023 \gamma_0 k_B / h$ $[\pm 0.24 \text{nW/(K} \cdot \text{m})]$ at $k_B T = 0.5W_L$. At low temperatures, the transition between these two plateaus is smooth and monotonic, while at higher temperatures, $\kappa_{xy}$ exhibits an oscillatory feature at $k_B T = 0.5W_L$ between two plateaus. On the other hand, $\kappa_{xx}$ displays a peak near the center LL, and its peak value increases quickly with $T$. The peak height is $0.028 \gamma_0 k_B / h$ $[0.3 \text{nW/(K} \cdot \text{m})]$ at $k_B T = 0.5W_L$. To test the validity of the Wiedemann-Franz law, we compare the above results with those calculated from Eq. (7), as shown in Figs. 5(c) and 5(d). The Wiedemann-Franz law predicts that the ratio of the thermal conductivity $\kappa$ to the electrical conductivity $\sigma$ of a metal is proportional to the temperature. This is in agreement with our low-temperature results, but apparent deviation is seen at higher temperatures.

In Fig. 6, we show the calculated thermal conductivities $\kappa_{xx}$ and $\kappa_{yy}$ for the biased ABC-stacked case. As seen from Figs. 6(a) and 6(b), around zero energy, a flat region with $\kappa_{yy} = 0$ is found at low temperatures, which is accompanied by a valley in $\kappa_{xx}$. These features are clearly in contrast to those of the ABA-stacked case due to the presence of an energy gap between the valence and conduction bands. When temperature increases to $k_B T = 0.06\Delta_g$, the plateau with $\kappa_{xx} = 0$ disappears, while $\kappa_{xx}$ displays a large peak. The peak height $\kappa_{xx}$ is $0.14\gamma_0 k_B / h$ $[1.49 \text{nW/(K} \cdot \text{m})]$ at $k_B T = 0.06\Delta_g$. In Figs. 6(c) and 6(d), we also compare the above results with those calculated from the Wiedemann-Franz law using Eq. (7). We find that the Wiedemann-Franz law remains valid at low temperatures.
ABA-stacked case. Around the Dirac point, the transverse thermoelectric conductivity $\alpha_{xy}$ exhibits a pronounced valley with $\alpha_{xy} = 0$ at low temperatures, in striking contrast to the ABA-stacked case with a peak. The validity of the semiclassical Mott relation between the thermoelectric and electrical transport coefficients is verified to be satisfied only at very low temperatures. Furthermore, the transverse thermal conductivity $\kappa_{xy}$ has a pronounced plateau with $\kappa_{xy} = 0$, which is accompanied by a valley in $\kappa_{xx}$. These are consistent with the opening of sizable gap between the valence and conduction bands in the biased ABC-stacked case.

We mention that in our numerical calculations, the flux $2\pi/M$ in each hexagon gives a magnetic field of the strength $B \sim 1.3 \times 10^5$ Tesla. Thus the magnetic field $B$ we used is about 2700 Tesla. This magnetic field is much stronger than the ones that can be achieved in the experimental situation, as limited by current computational capability. In our calculation, the system size is taken to be $N = 96 \times 24 \times 3$, and $M$ is taken to be $L_x$ or $L_y$ consistent with periodic boundary conditions, which limits us to extremely strong magnetic fields. However, the obtained thermoelectric transport coefficients exhibit universal behaviors, as long as $M$ is not too small (greater than 10).

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Different literatures may have a sign difference due to different conventions.


