Hot electron effects and electric field scaling near the metal-insulator transition in multilayer MoS_2

Byoung Hee Moon,[1,2,*] Gang Hee Han,[1,3] and Young Hee Lee[1,2,†]

[1]*Center for Integrated Nanostructure Physics, Institute for Basic Science (IBS), Suwon 16419, Republic of Korea*

[2]*Department of Energy Science, Sungkyunkwan University, Suwon 16419, Republic of Korea*

[3]*Department of Physics, Incheon National University, Incheon 22012, Republic of Korea*

(Received 5 November 2019; revised manuscript received 9 January 2020; published 24 January 2020)

The layered transition metal dichalcogenides have emerged as valuable platforms to study the challenging problem of metal-insulator transition in two dimensions. It was demonstrated that multilayer MoS_2 exhibits clearly distinctive metallic and insulating behaviors in conductivity in response to both temperature and the electric field. Here, we report on the scaling analyses of conductivity for the electric field in addition to the temperature, which is performed with the consideration of electron-electron interactions for multilayer MoS_2. Based on the analysis of hot electron effects in the electric field, we find that scaling for the electric field is relevant for the metallic phase in the high-field regime, enabling one to extract the dynamical critical exponent z close to 1. This result supports that the metal-insulator transition in multilayer MoS_2 is a true quantum critical phenomenon, in which strong interactions induce the transition.

DOI: 10.1103/PhysRevB.101.035421

Transition metal dichalcogenides (TMDs) have distinctive features that enable the study of metal-insulator transition (MIT) in disordered two-dimensional (2D) systems [1]. When compared to conventional systems, such as Si metal-oxide-semiconductor field-effect transistors (MOSFETs) or the heterojunctions of different semiconductors [2,3], TMDs exhibit the characteristic advantages for the investigation of the interplay between disorder and carrier-carrier interactions, which is a key aspect to understanding the mechanism of the MIT in 2D systems. In layered TMDs, this interplay can be modulated more easily by varying the thickness or passivating the surface of the samples. For example, MIT phenomena in molybdenum disulfide, MoS_2, were analyzed for various thicknesses using a scaling law. It was concluded that the interaction effects in thicker samples become more significant and might even drive the localization [1].

The scaling behavior for an appropriate physical variable is characteristic of a second-order phase transition. In the localization problem, the correlation lengths ξ in the spatial direction, i.e., the localization length, and ξ_τ in the time direction are the only characteristic length scales that diverge at transition, leading to a scaling behavior of the conductivity (σ) with temperature (T) near the MIT [4]. In addition to temperature, a useful parameter for the scaling behavior is the current or the applied electric field (E). If the nonlinear conductivity for the electric field near MIT is from critical fluctuations, electric field scaling (E scaling) together with temperature scaling (T scaling) provides valuable information on the critical exponents; in particular, the values of the correlation length exponent ν and the dynamic critical exponent z can be obtained separately. However, in many cases, the nonlinearity in conductivity for the electric field arises due to heating effects rather than field effects, causing E scaling to be irrelevant [5–7].

The electric field–dependent conductivity (σ_E) of a few-layered MoS_2 near the MIT and an early attempt of E scaling in TMDs were reported in Ref. [1], in which σ_E exhibits clearly separate metallic and insulating behaviors also. However, the relevance of the heating effect was checked only at one carrier density for each metallic and insulating side, leading to an incomplete conclusion. In this study, we discuss this heating effect in the electric field–dependent conductivity in more detail for a broader range of carrier densities (or back-gate voltages, V_{BG}) around the MIT. We find that the heating effect is significant for both insulating and metallic phases at low-field regimes, but it becomes negligible just near the transition in the insulating side. This discontinuity provides support for the critical phenomenon of MIT in this system. Further, we find that the conductivity for the higher field in the metallic phase is not explained by the heating effect. Accordingly, we apply E scaling for this regime to extract the values of ν and z separately, which reveals that Coulomb interactions between carriers play a significant role in the transition in the few-layered MoS_2.

To fabricate the device, MoS_2 was mechanically exfoliated on a SiO_2 (300 nm)/Si (heavily p-doped) wafer. The thickness of a chosen flake was found to be ~7 nm through atomic force microscopy. The electrodes were patterned by e-beam lithography, followed by Cr (1 nm)/Au (60 nm) deposition. The optical image of a four-probe device is presented in the inset of Fig. 1(a). Figure 1(a) displays the conductivity (σ_T) as a function of temperature down to 2 K for back-gate biases (V_{BG}) from 0 to 60 V at intervals of 3 V. In this report, the conductivity was taken from the inverse of the sheet resistance. The conductivity was extracted in the zero limit of drain-source voltage (V_{ds}) (discussed later with regard to Fig. 2). The conductivity behaves differently at low temperature: $\Delta\sigma/\Delta T > 0$ and $\Delta\sigma/\Delta T < 0$. This suggests the

*ibhmoon@skku.edu

†leeyoung@skku.edu

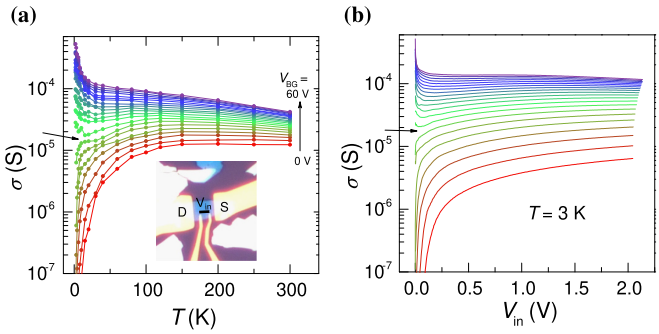

FIG. 1. (a) Electric conductivity (σ) as a function of temperature (T) for back-gate bias V_{BG} from 0 (bottom) to 60 V (top) with 3-V steps. The inset shows an optical image of a multilayer MoS_2. (b) σ as a function of voltage V_{in} ($\propto E$) at $T = 3$ K for the same V_{BG} range as in (a). Arrows indicate the separation of metallic and insulating conductivity behaviors.

occurrence of an MIT around $V_{BG} = 21$ V which corresponds to the critical carrier density $n_c \sim 3.95 \times 10^{12}$ cm^{-2} at room temperature, estimated from the simple approximation, $n_c = C_{ox}(V_{BG} - V_{TH})/q$, where C_{ox} is the geometrical oxide capacitance, V_{TH} is the threshold voltage, and q is the elementary charge. At this carrier density, the Fermi temperature is $T_F \sim 220$ K and the dimensionless interaction parameter, the ratio of Coulomb energy to kinetic energy at the Fermi level, is $r_s = m^* q^2 g_V/4\pi\varepsilon\varepsilon_0\hbar^2\sqrt{n_c} \sim 8$. Here, the effective mass of electron $m^* \sim 0.5m_0$; the valley degeneracy $g_V = 6$ [8]; the relative dielectric constant $\varepsilon \sim 8$ for multilayer MoS_2 [9], where m_0 is the free electron mass and ε_0 is the vacuum dielectric constant. The large value of $r_s \gg 1$ indicates the importance of carrier-carrier interactions in this system.

The identification of the critical bias V_c is supported by the drain-source voltage-dependent conductivity measured at $T = 3$ K, as illustrated in Fig. 1(b), where the behaviors $\Delta\sigma/\Delta V_{in} > 0$ and $\Delta\sigma/\Delta V_{in} < 0$ are clearly distinguished at the same V_{BG}. Here, V_{in} represents the voltage drop between the two inner probes. This study initially discusses the origin of this contrasting nonlinear conductivity behavior with varying electric field around the MIT and consequently demonstrates the scaling behavior based on a conventional scaling scheme for 2D systems [2,10]. The nonmonotonic temperature dependence of conductivity in the metallic phase near the transition in Fig. 1(a) has been a significant and still controversial issue in revealing the MIT mechanism in 2D systems [11–13]. This issue will be discussed in another publication.

According to conventional hopping theory, the nonlinear response of the conductivity to the electric field in the insulating state is, in general, due to the field-enhanced hopping with the exchange of phonons [14,15]. In contrast, electron heating is considered to be the primary cause of the nonlinear σ_E in the metallic state, particularly at low temperature, owing to the weak coupling to phonons [16]. However, near the MIT, the situation becomes more intricate. In the insulating side, the localization length ξ increases as it approaches the transition as $\xi \sim |n_{2D} - n_c|^{-\nu}$, where n_{2D} is the carrier density and n_c is the critical carrier density at $V_{BG} = V_c$. In this case, the diffusive character of electrons can be developed such that heating effects become non-negligible [5]. On the other hand, in the metallic phase, the timescale τ_c, which controls the response to the electric field, is slow and follows $\tau_c \sim \xi^z$ near the transition, so that a cooling bottleneck for hot electrons may be an insignificant factor [4]. In other words, the critical fluctuations by a field rather than the thermal effects may be responsible for the nonlinear conductivity in the metallic phase near the MIT. For the electric field scaling analysis to be valid, it is crucial to check that the thermal effect is negligible compared to the field effect in the region of the nonlinear σ_E.

In our previous report on a thinner MoS_2 (~3.5 nm thickness) [1], we concluded that the electron heating effect is significant in the metallic phase but not in the insulating phase. However, this conclusion was based on an analysis with only one carrier density in each metallic and insulating side. Here, we conducted the same analysis as in Ref. [1], but on a thicker MoS_2 (~7 nm) with various carrier densities near the MIT. We found that the above conclusion is not valid for a wider range of carrier densities in the insulating phase. We also found that the heating effect in the metallic phase is significant only in the low-field regime. For the analysis, we used the following heat-balance equation [17],

$$P = I_{ds}V_{in} = \Sigma\left(T_e^\delta - T_L^\delta\right) = \Sigma\left(T_e^\delta - T^\delta\right), \qquad (1)$$

where I_{ds} is the drain-source current, Σ is the coupling constant, T_e is the effective electron temperature, and δ is the exponent of the power law. Here, T_L is replaced by the system temperature T by assuming that the acoustic phonon is the dominant cooling mode at low temperatures and that it reaches thermal equilibrium with the substrate (system) much faster than the phonon-electron relaxation [18]. Because of this slow phonon-electron relaxation, when the power P is introduced into the system, it heats the electrons to a higher temperature than the lattice temperature T_L at low temperature. In this case, scaling is controlled by the temperature, indicating the failure of E scaling. This can be verified by checking δ in Eq. (1), which was predicted to be $\delta = 4$–6, and has been observed for most cases of disordered thin films including monolayer MoS_2 at low temperature [18–20]. If the experimental value of δ is much larger than this, the cooling rate by phonon emission is sufficient, and the field effect is more responsible for the nonlinear σ_E [4].

To apply this scenario, we first convert σ_E to the effective electron temperature T_e based on σ_T. Figures 2(a)-2(d) present the examples of extracting σ_T in the zero voltage (V_{in}) limit for chosen values of T and V_{BG}. Figure 2(e) displays σ_T as a function of T for various values of V_{BG}. The red lines are fittings obtained using mathematical functions and will be utilized to convert σ_E into T_e. Figures 3(a) and 3(b) illustrate the resulting plots of P vs T_e for various values of V_{BG} in the insulating and the metallic phases, respectively. First, the slope change occurs at $T_e \sim 10$ K in both phases, which is expected to be close to the Bloch-Grüneisen temperature T_{BG}. Above T_{BG}, δ tends to decrease because the phonon emission is not restricted by Pauli blocking [18]. This trend ceases to hold in the metallic phase far from the transition. In addition, δ is much larger than 6, indicating that the field effect is already dominant over the heating effect. Below $T_e \sim 10$ K in the metallic phase, δ is close to or slightly smaller than 6 for all back-gate biases. Thus, we conclude that the heating

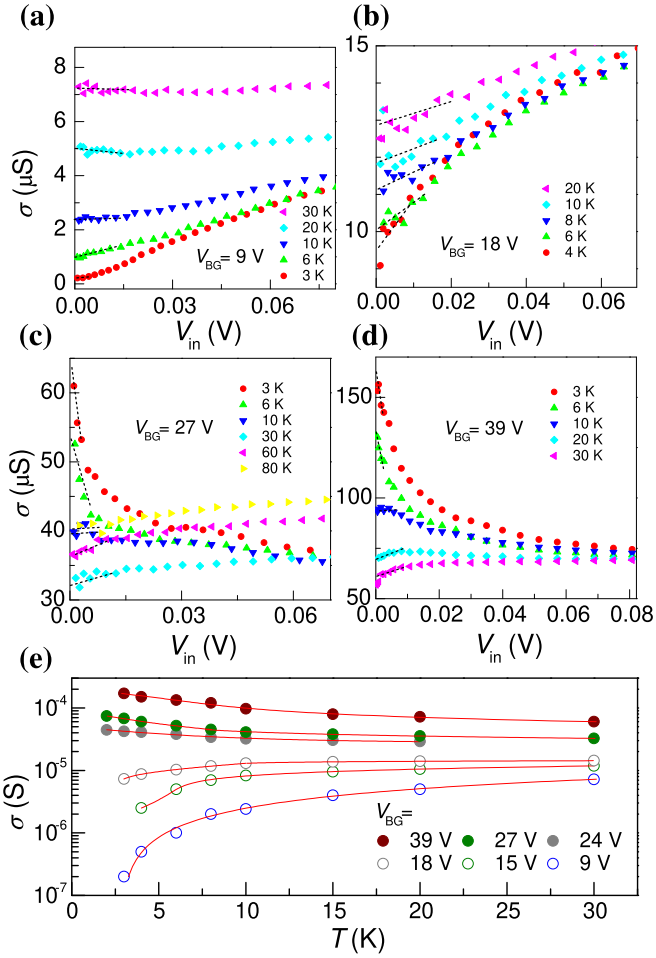

FIG. 2. (a)–(d) σ vs $V_{\rm in}$ ($\propto E$) at several temperatures T for back-gate biases $V_{\rm BG} = 9$, 18, 27, and 39 V, respectively, showing the extractions of σ in the zero limit of $V_{\rm in}$. (e) Extracted σ as a function of T for various $V_{\rm BG}$ values. Red lines are fittings obtained from mathematical functions for the conversion of E into the effective electron temperature T_e.

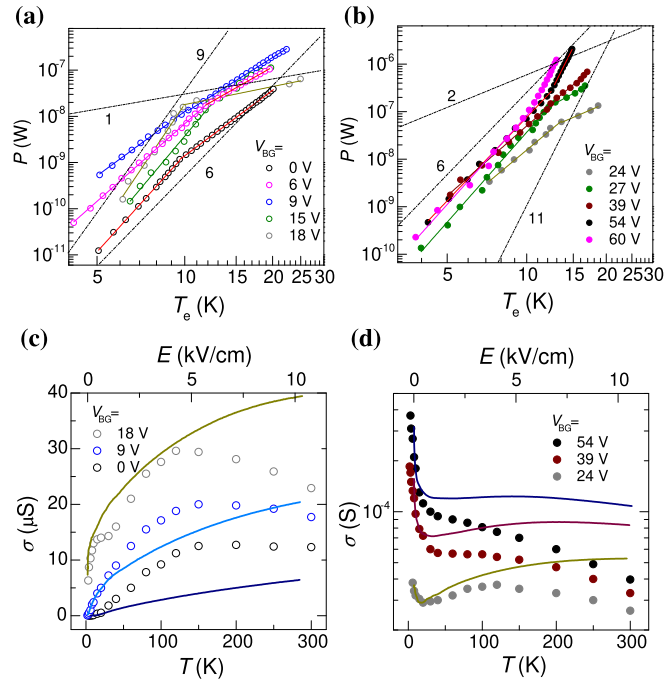

FIG. 3. (a), (b) Power P as a function of T_e for various $V_{\rm BG}$ values for the insulating and metallic phases. Dashed lines correspond to the slopes δ of the values next to each line. (c), (d) σ for temperature (bottom axes) and electric field (top axes) for the insulating and metallic phases.

effect is dominant in the metallic phase at $T_e \lesssim 10$ K. On the other hand, in the insulating phase, $\delta \sim 9.4$ is significantly larger than 6 at $V_{\rm BG} = 18$ V just near the transition, as shown in Fig. 3(a), which is consistent with our previous report [1]. However, deeper into the insulating phase, the value of δ becomes closer to 6. This trend indicates the significance of the field effect near the transition and the heating effect deep in the insulating phase.

To confirm these conclusions, we plot σ_T against σ_E in Figs. 3(c) and 3(d) for the insulating and the metallic phases, respectively. In these figures, data points and solid lines indicate σ_T (bottom axis) and σ_E (top axis), respectively. Near the transition in the insulating side, σ_E is larger than the maximum of σ_T for $E > 5$ kV/cm at $V_{\rm BG} = 18$ V [Fig. 3(c)]. Thus, the thermal effect cannot explain this large value of σ_E. In addition, the shoulder feature at $T \sim 10$ K in σ_T is absent in σ_E, confirming the significance of the field effect on the nonlinearity of σ_E. The feature of a larger value of σ_E than the maximum of σ_T becomes weaker and disappears deeper in the insulating phase (i.e., smaller $V_{\rm BG}$), which suggests a higher possibility that, even with strong fields, the heating effect is the main origin of the nonlinear σ_E deep in the insulating phase. In contrast, σ_E is always larger than σ_T except at very low-field regimes in the metallic phase, and exhibits a nonmonotonic behavior even at the highest $V_{\rm BG} = 54$ V (deep metallic phase), where σ_T is monotonic. Thus, we can deduce that the main source of the nonlinear σ_E in the metallic phase is the heating effect in the low-field regime and the field effect in the high-field regime. Regarding the nonmonotonicity in σ_E, it was theoretically recognized that the static electric field can induce the nonlinear response in conductivity via the dephasing effect in quantum interference of interacting electrons. In particular, when the scattering amplitude in the particle-hole triplet channel is considerable, the large field is predicted to have an antilocalizing effect; i.e., the correction to the conductivity is positive [21,22]. We suspect that this nonthermal effect possibly plays a significant role in nonmonotonic σ_E in the metallic phase.

Figure 4 summarizes the values of δ in the low-field regime (or small T_e). As previously discussed, δ fluctuates around 6 for the entire $V_{\rm BG}$ range except just near the transition in the insulating side. This discontinuity strongly suggests that, in this system, the MIT is a critical behavior rather than the crossover between the strong and weak localizations [23,24]. It is worth noting that the significant electron heating by the electric field deep in the insulating phase indicates that the resistivity (or conductivity) is determined by the electron temperature rather than the phonon temperature at low temperatures. This suggests that hopping is assisted by electron-electron interactions [25] rather than by phonons as usually assumed in conventional hopping theories [26–28].

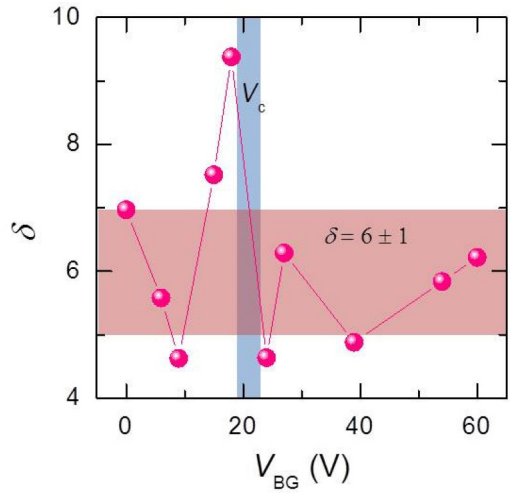

FIG. 4. Exponent δ for various values of V_{BG}.

The significance of electron-electron interactions in TMDs has also been noticed for MIT phenomena [1,29], as the delocalizing effect by strong interactions has been believed to be the origin of MIT as a quantum phase transition in 2D systems [11,30,31]. In this case, the conductivity near the MIT satisfies the scaling law for the temperature (T scaling) and the electric field (E scaling) [1,4]. It is described by the equations $\sigma(T,\ \delta n) = \sigma_c(T)F_T[T/T_0(\delta n)]$ and $\sigma(E,\ \delta n) = \sigma_c(E)F_E[E/E_0(\delta n)]$ for T and E scaling, respectively, where F is the universal scaling function, $\delta n \equiv (n_{2D}/n_c - 1)$, $T_0(\delta n) \sim |\delta n|^{z\nu}$, and $E_0(\delta n) \sim |\delta n|^{(1+z)\nu}$. From these two scaling analyses, the exponents z and ν can, in principle, be obtained separately. It is evident that the $\sigma_{T(E)}$ for various carrier densities cannot be collapsed into a single curve after rescaling them by the scaling parameters T_0 (E_0) because they are nonmonotonic with respect to T (E). If instead of such a one-parameter scaling scheme we consider one more scaling variable for interactions, which is equivalent to scaling the conductivity normalized by a temperature (electric field)-dependent critical conductivity $\sigma_c(T)[[\sigma_c(E)]$ [32], i.e., σ/σ_c, well behaved scaling results are observed, as illustrated in Fig. 5. Figures 5(a) and 5(d) present the renormalized σ_T and σ_E and Figs. 5(b) and 5(e) demonstrate the collapses of the renormalized σ_T and σ_E after rescaling T and E. Figures 5(c) and 5(f) depict the plots of the scaling parameters T_0 and E_0, respectively. We note that the collapse of the rescaled σ/σ_c for the temperature is fairly reasonable for the wide temperature range from the ballistic to the diffusive regime below the Dingle temperature $T_D = \hbar q/2k_B m^*\mu \sim 9$ K, where the field effect mobility $\mu \sim 1500\,\mathrm{cm^2\,V^{-1}\,s^{-1}}$ at 2 K.

In E scaling, we disregard the insulating phase part because of the heating effect, as we discussed earlier. For the metallic phase, we exclude the low-field data in which the heating effect is significant. From the power-law behaviors of T_0 for δn as in Fig. 5(c), we obtain $z\nu = 2.24$ and 1.71 for the insulating and metallic phases, respectively. Similarly, from E_0 for δn, as depicted in Fig. 5(f), we obtain $(1+z)\nu = 3.65$ for the metallic phase. These two scaling results yield the separate values of $z = 0.88$ and $\nu = 1.94$. The result $z \approx 1$ was also observed in a low-disorder system [10], which was interpreted as the result of strong long-range electron-electron interactions [4]. We stress that the z value in our

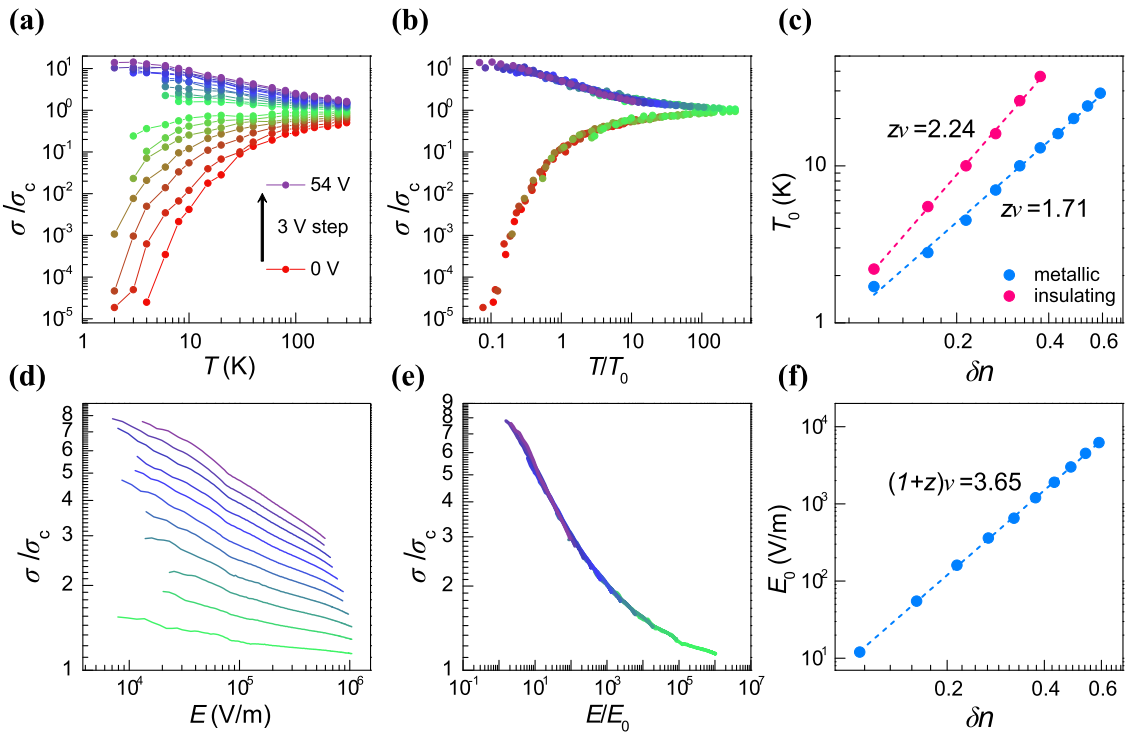

FIG. 5. (a), (d) Normalized conductivity σ/σ_c as a function of temperature and electric field for V_{BG} ranging from 0 to 54 V in steps of 3 V. (b), (e) σ/σ_c as a function of normalized temperature T/T_0 and electric field E/E_0. (c), (f) Scaling parameters T_0 and E_0 as a function of δn.

multilayer MoS_2 is close to 1, which is consistent with the idea of scaling for two parameters related to disorder and interactions.

In summary, we measured the conductivity of multilayer MoS_2 for temperatures and electric fields around the MIT. In this system, the nonlinearity in the electric field–dependent conductivity (σ_E) emerges due to electron heating in the metallic phase in the low-field regime and in the insulating phase but not in the region just near the transition. Based on this, we derived several conclusions about this system: (i) Hhopping deep in the insulating phase is induced by electron-electron interactions rather than phonons. (ii) Discontinuity in the exponent δ across the transition corroborates the critical phenomenon of MIT. (iii) The dynamic critical exponent $z \approx 1$ is consistent with a quantum phase transition in strongly correlated systems.

This work was supported by the Institute for Basic Science (IBS-R011-D1).

B.H.M. and G.H.H. contributed equally to this work.

[1] B. H. Moon, J. J. Bae, M.-K. Joo, H. Choi, G. H. Han, H. Lim, and Y. H. Lee, Nat. Commun. **9**, 2052 (2018).
[2] S. V. Kravchenko, G. V. Kravchenko, J. E. Furneaux, V. M. Pudalov, and M. DIorio, Phys. Rev. B **50**, 8039 (1994).
[3] S. V. Kravchenko and M. P. Sarachik, Rep. Prog. Phys. **67**, 1 (2004).
[4] S. L. Sondhi, S. M. Girvin, J. P. Carini, and D. Shahar, Rev. Mod. Phys. **69**, 315 (1997).
[5] M. E. Gershenson, Y. B. Khavin, D. Reuter, P. Schafmeister, and A. D. Wieck, Phys. Rev. Lett. **85**, 1718 (2000).
[6] R. Leturcq, D. L'hote, R. Tourbot, V. Senz, U. Gennser, T. Ihn, K. Ensslin, G. Dehlinger, and D. Grützmacher, Europhys. Lett. **61**, 499 (2003).
[7] M. Osofsky, M. LaMadrid, J. B. Bieri, W. Contrata, J. Gavilano, and J. M. Mochel, Phys. Rev. B **38**, 8486 (1988).
[8] J. Hong, C. Lee, J. S. Park, and J. H. Shim, Phys. Rev. B **93**, 035445 (2016).
[9] X. Chen, Z. Wu, S. Xu, L. Wang, R. Huang, Y. Han, W. Ye, W. Xiong, T. Han, and G. Long *et al.*, Nat. Commun. **6**, 6088 (2015).
[10] S. V. Kravchenko, D. Simonian, M. P. Sarachik, W. Mason, and J. E. Furneaux, Phys. Rev. Lett. **77**, 4938 (1996).
[11] A. Punnoose and A. M. Finkel'stein, Phys. Rev. Lett. **88**, 016802 (2001).
[12] S. Anissimova, S. V. Kravchenko, A. Punnoose, A. M. Finkel'stein, and T. M. Klapwijk, Nat. Phys. **3**, 707 (2007).
[13] M. M. Radonjic, D. Tanaskovic, V. Dobrosavljevic, K. Haule, and G. Kotliar, Phys. Rev. B **85**, 085133 (2012).
[14] A. Larkin and D. Khmelnitskiĭ, Sov. Phys. JETP **56**, 647 (1982).
[15] A. Dvurechenskiĭ, V. Dravin, and A. Yakimov, JETP Lett. **48**, 155 (1988).
[16] M. L. Roukes, M. R. Freeman, R. S. Germain, R. C. Richardson, and M. B. Ketchen, Phys. Rev. Lett. **55**, 422 (1985).
[17] J. K. Viljas and T. T. Heikkila, Phys. Rev. B **81**, 245404 (2010).
[18] K. Kaasbjerg, K. S. Bhargavi, and S. S. Kubakaddi, Phys. Rev. B **90**, 165436 (2014).
[19] J. T. Karvonen, L. J. Taskinen, and I. J. Maasilta, Phys. Rev. B **72**, 012302 (2005).
[20] M. Ovadia, B. Sacépé, and D. Shahar, Phys. Rev. Lett. **102**, 176802 (2009).
[21] R. Raimondi, P. Schwab, and C. Castellani, Phys. Rev. B **60**, 5818 (1999).
[22] M. Leadbeater, R. Raimondi, P. Schwab, and C. Castellani, Eur. Phys. J. B **15**, 277 (2000).
[23] M. Y. Simmons, A. R. Hamilton, M. Pepper, E. H. Linfield, P. D. Rose, and D. A. Ritchie, Phys. Rev. Lett. **84**, 2489 (2000).
[24] S. Das Sarma and E. H. Hwang, Phys. Rev. B **89**, 235423 (2014).
[25] S. Marnieros, L. Berge, A. Juillard, and L. Dumoulin, Phys. Rev. Lett. **84**, 2469 (2000).
[26] N. Mott, J. Non-Cryst. Solids **1**, 1 (1968).
[27] A. Efros and B. I. Shklovskii, J. Phys. C **8**, L49 (1975).
[28] V. Ambegaokar, B. Halperin, and J. Langer, Phys. Rev. B **4**, 2612 (1971).
[29] B. Radisavljevic and A. Kis, Nat. Mater. **12**, 815 (2013).
[30] A. Finkel'shteĭn, Sov. Phys. JETP **57**, 97 (1983).
[31] A. Punnoose and A. M. Finkel'stein, Science **310**, 289 (2005).
[32] D. A. Knyazev, O. E. Omel'yanovskii, V. M. Pudalov, and I. S. Burmistrov, Phys. Rev. Lett. **100**, 046405 (2008).