Overdamped Phase Diffusion in hBN Encapsulated Graphene Josephson Junctions

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We investigate the zero-bias behavior of Josephson junctions made of encapsulated graphene boron nitride heterostructures in the long ballistic junction regime. For temperatures down to 2.7K, the junctions appear non-hysteretic with respect to the switching and retrapping currents $I_S$ and $I_R$. A small non-zero resistance is observed even around zero bias current, and scales with temperature as dictated by the phase diffusion mechanism. By varying the graphene carrier concentration we are able to confirm that the observed phase diffusion mechanism follows the trend for an overdamped Josephson junction. This is in contrast with the majority of graphene-based junctions which are underdamped and shorted by the environment at high frequencies.

Graphene-based superconductor-normal metal-superconductor (SNS) Josephson junctions have been a popular medium of choice for studying the fundamentals[11] as well as applications[12] of superconducting devices for more than a decade. However, the full spectrum and consequences of the interactions between the graphene Josephson junction and the environment have not been fully mapped. For example, the observed critical current $I_C$ of graphene Josephson junctions is consistently suppressed compared to theoretical predictions; leading to postulations that the junctions are severely underdamped[5,6,9,10,18], despite the relatively low hysteresis between the switching $I_S$ and the retrapping $I_R$ currents. These junction dynamics via it’s interaction with the environment can be directly investigated by looking at the statistical distribution of the switching current $I_S$ [7,9,19,20], or via the measurement of zero-bias resistance arising from the phase diffusion mechanism[5,18,21,25]. Indeed, often attributed to the large capacitance generated by the bonding pads and leads, previous works have shown that the vast majority of graphene-based JJs are underdamped[5,6,9,10].

In this work we report on Josephson junctions made from hexagonal boron-nitride (hBN) encapsulated graphene with Molybdenum-Rhenium (MoRe) alloy superconducting contacts [14]. These devices are governed by ballistic electron transport and have been found to be in the intermediate to long-junction regime[9]. Here, the MoRe contacts terminate shortly after the active region and are connected to the bonding pads via thin Gold leads. Moreover, the MoRe-Au interface seems to exhibit a significant contact resistance. The junctions were measured at temperatures between 2.7 and 7K, where a clear phase diffusion governed, zero-bias resistance can be observed[5,18,21,25]. However, for these devices, when changing the carrier concentration via the back gate, the zero-bias resistance follows the trend expected for overdamped junctions[5,25,27]. Thus, we conclude that we have demonstrated ballistic graphene Josephson junction in the overdamped regime.

Graphene is made using the exfoliation method[28] and is encapsulated in hexagonal boron-nitride (hBN) using the “pick-up” method[26]. Using CHF3/O2 plasma, the hBN-Graphene-hBN stack is etched thorough in order to make quasi-one dimensional electrical contacts with superconducting electrodes[14]. Molybdenum-Rhenium(MoRe) alloy electrodes are deposited onto the device using DC sputtering with the approximate thickness of $100 - 120nm$. The bonding pads and thin metal leads making contact to MoRe are made of Cr/Au (5nm/110nm). Here, we present data on the device of length $L = 500nm$ (the distance between MoRe contacts), and the width $W = 3µm$.

The device is measured in a home made crycooler with a base temperature of $\sim 2.5K$, and isolated via a heat shield and R-C filters placed at the low temperature stage. Josephson junction resistance is measured using the lockin method with a four-probe geometry. The junction is biased by a variable DC current with a small AC excitation of 5nA. The voltage across the junction is amplified using a custom differential pre-amp prior to being fed into the lockin. The gate voltage applied to the back of the 300nm SiO2 oxide layer is used to control the carrier density of graphene. Figure 1 presents the differential resistance $dV/dI$ as a function of applied DC bias current $I_{Bias}$. All the curves show two transition points. As the bias current is swept from a large negative value to a large positive value, the absolute value of the current on the negative side below which the junc-
tion becomes superconducting is the retrapping current $I_R$, i.e. $|I_{Bias}| = I_R$. On the positive side, the junction transitions from superconducting to normal state at the switching current $I_S$ [18]. Figure 1(a) shows resistance versus $I_{Bias}$ for different temperatures with the backgate voltage set to 35V. Figure 1(b) shows the resistance versus bias current for different gate voltages taken at 3K. In both cases the switching and retrapping currents $I_S$ and $I_R$ follow the expected trends: falling exponentially with increasing temperature, and increasing with gate voltage away from the Dirac point with the hole conduction regime exhibiting a suppressed critical current due to the effects of contact doping [5, 6, 9].

The vast majority of previously reported graphene

Josephson junctions, exhibit hysteresis between the switching and retrapping currents, with $I_S > I_R$ even for temperatures above 3K. While certain works attribute this hysteresis to self-heating of the junction [13, 29], it
has also been shown that most of the graphene Josephson junctions exhibit underdamped behavior \[5, 21–25\]. (Graphene-based Josephson junctions a typically overdamped due to the large capacitance caused by the presence of the backgate.) Here, however, for all tested gates and temperatures, we do not observe a difference between \( I_S \) and \( I_R \). This leads to an initial indication that the device may be overdamped.

We now further investigate the junction dynamics via the characterization of device behavior in the phase diffusion regime. Phase diffusion manifests itself as an observable non-zero resistance below the critical current (even at \( I_{Bias} = 0 \), arising from phase slips that are caused by thermal noise. The rate of these phase slips down the prototypical titled washboard potential, and therefore, the measured zero-bias resistance is governed by the junction dynamics which dictate the energy dissipation rate \[18\]. Indeed, we are able to observe a measurable resistance in our devices, even at zero bias, and down to 3K in temperature. We define the measured zero-bias differential resistance as \( R_0 \). In order to confirm that \( R_0 \) arises from the phase diffusion mechanism, we study the evolution of this resistance with respect to temperature. Regardless of the junction damping dynamics the trend behavior of \( R_0 \) with respect to temperature should have the following dependence:\[5, 21–25\]:

\[
R_0(T) \propto T^{-1} e^{-2E_J/k_B T} \tag{1}
\]

Here, \( E_J = \hbar I_C / 2e \) is the Josephson energy. Reworking equation (1), we can arrive at the proportionality relationship: \( \log (R_0(T) \cdot T) \propto -2E_J/k_B T \). In Figure 2a we plot the value \( R_0 \cdot T \) versus inverse temperature \( T \) on a semi-log scale. Indeed, we find that the relationship is nearly linear, consistent with theory. From here, knowing the temperature, we can extract \( E_J \).

We now define the prefactor to the exponential in Equation 1 as a variable \( R'_0 \) with \( R'_0 \equiv R_0 e^{(2E_J/k_B T)} \) \[5\]. The relationship between \( R'_0 \) and \( E_J \) is different for depending on the damping dynamics of the junction. Previous theoretical works have defined three different regimes: For overdamped junctions, the expected relationship is \( R'_0 \sim R_N E_J/k_B T \) \[25, 27\]. Here \( R_N \) is the normal resistance, i.e. the resistance the junction when it is in normal state. For underdamped junctions, \( R'_0 \sim (\hbar/e^2) \hbar \omega_p/k_B T \) \[23\], where \( \omega_p \propto \sqrt{E_J/C} \) is the plasma frequency. (And \( C \) is the capacitance of the junction.) Finally, if the junction is underdamped at low frequencies, but becomes overdamped at the plasma frequency (due to being shorted by the environment), we have \( R'_0 \sim Z_0 E_J/k_B T \) \[21\]. Here \( Z_0 \) is the real part of the high frequency impedance caused by the junction’s environment \[21\]. (Typically \( Z_0 \) is found to be \( \sim 200 \sim 250 \Omega \).)

The calculated \( R'_0 \) versus the Josephson energy is shown in Figure 2(b) data taken at 3.5K. Alongside, we plot the fitting results obtained for each of the three cases mentioned above. Note that for the underdamped, and dumping the environment case the Josephson energy \( E_J \) is treated at the independent parameter in the fit. For the overdamped case we treat the gate voltage \( V_G \) as the independent parameter and the Josephson energy \( E_J \) is calculated from the from the critical current \( I_C \) measured at that \( V_G \). This was done in order to represent a more precise trend for \( R_N \). Due to contact resistance (resistance caused by imperfect electrical contact between MoRe and graphene), or device processing errors the measured resistance may be different from the true normal resistance \( R_N \). However, for graphene governed by ballistic electron transport, the inverse resistance is proportional to the square root of gate voltage, i.e. \( R_N \propto 1/\sqrt{|V_G|} \[30\]. It can be clearly seen that the measured data conforms well to the trend expected for an overdamped junction.

![FIG. 3. Gate voltage dependence of measured normal resistance \( R_N \) (black line) compared to the resistance calculated from the theoretically predicted relationship between \( R_N \) and \( V_G \): \( R_N = R_0/[I_0(E_J/2k_B T)]^{-2} \) (red line). (\( I_0 \) is the modified Bessel function.) The measured \( R_N \) here is taken at 4K. The data and calculated result match well, supporting the claim that the measured device is and overdamped junction.](image)

Finally, we confirm the above observation by comparing the measured normal resistance \( R_N \) with that back calculated from the the zero-bias resistance \( R_0 \). Following the expression in Ref. \[27\] we have:

\[
\lim_{I_{Bias} \rightarrow 0} \frac{V/I_{Bias}}{R_N} = [I_0(\frac{1}{2})]^2 \tag{2}
\]

Here, \( I_0 \) is the modified Bessel function, and \( \gamma = \sqrt{I_0(E_J/2k_B T)} \). For \( I_{Bias} \) approaching zero, the equation simplifies to \( R_N = R_0/[I_0(E_J/2k_B T)]^{-2} \). Figure 3 shows the measured normal resistance \( R_N \) versus gate voltage \( V_G \) taken at 4K (black line): while the resistance calculated from Equation 2 is plotted in red. It can be seen that we have a good match between the measured and the theoretical result, in particular for high values of \( R_N \) (close
to the Dirac point). This further supports our claim that we are observing overdamped dynamics.

Whether the junction is underdamped, or overdamped is typically determined by the quality factor $Q$, with $Q = R_N (2eI_C C/h)^{1/2} = \pi R_N \sqrt{E_J C}$. (C being the capacitance of the junction.) A $Q < 0.85$ results in an overdamped junction. However, the difficulty in calculating an accurate quality factor $Q$ is two fold: One, the measured critical current $I_C$ has been consistently less than the value predicted by theory leading to the question of whether the true critical current is being measured. Two, it is unclear which areas of the device play an active role in the junction capacitance $C$. (Previous works have suggested that the capacitance all the way up to the capacitive coupling between the device bonding pads plays a significant role. Thus the junction capacitance is always several orders of magnitude larger compared to that expected from the dimensions of the SNS region itself.) Assuming that the junction capacitance includes the contribution of the $100 \times 100 \mu m$ bonding pads that couple to each other via the backgate below the $300 nm$ thick SiO$_2$ layer, we arrive at $C \approx 600 F$. Taking $E_J = 0.3 meV$ and $R_N = 100 \Omega$, we calculate a quality factor $Q = 1.6$: the overdamped regime. Hence, we are led to conclude that the tested junction has been sufficiently isolated from the capacitive contribution of the bonding pads. Indeed, we estimate that the interface between MoRe superconducting contacts and the gold leads suffers from a contact resistance in the order of tens of $\Omega$, which is sufficient to reduce the quality factor below $1$.

In conclusion, we have investigated the phase diffusion regime in hBN encapsulated graphene Josephson junction governed by ballistic electron transport. The observed trend of the measured zero-bias resistance $R_0$ with respect to carrier concentration conforms well to theory describing phase diffusion in an overdamped junction regime. This is the first conclusive confirmation of overdamped behavior in graphene-based Josephson junction. We attribute this behavior to effective isolation of the Josephson junction from the capacitive contribution of the bonding pads. The isolation arises from a resistive connection with the device layout.

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[1] H. B. Heersche, P. Jarillo-Herrero, J. B. Oostinga, L. M. K. Vandersypen, and A. F. Morpurgo, Nature 446, 56 (2007)
[2] F. Miao, S. Wijeratne, Y. Zhang, U. C. Coskun, W. Bao, and C. N. Lau, Science 317, 1530 (2007)
[3] X. Du, I. Sliachko, and E. Y. Andrei, Phys. Rev. B 77, 184507 (2008)
[4] C. Ojeda-Aristizabal, M. Ferrier, S. Guéron, and H. Bouchiat, Phys. Rev. B 79, 165436 (2009)
[5] I. V. Borzenets, U. C. Coskun, S. J. Jones, and G. Finkelstein, Phys. Rev. Lett. 107, 137005 (2011)
[6] C. T. Ke, I. V. Borzenets, A. W. Draelos, F. Amet, Y. Bomze, G. Jones, M. Craciun, S. Russo, M. Yamamoto, S. Tarucha, and G. Finkelstein, Nano Letters 16, 4788 (2016)
[7] G.-H. Lee, D. Jeong, J.-H. Choi, Y.-J. Doh, and H.-J. Lee, Phys. Rev. Lett. 107, 146605 (2011)
[8] U. C. Coskun, M. Brenner, T. Hynel, V. Yakaryuk, A. Levchenko, and A. Bezyadin, Phys. Rev. Lett. 108, 097003 (2012)
[9] I. V. Borzenets, F. Amet, C. T. Ke, A. W. Draelos, M. T. Wei, A. Seredinski, K. Watanabe, T. Taniguchi, Y. Bomze, M. Yamamoto, S. Tarucha, and G. Finkelstein, Phys. Rev. Lett. 117, 237002 (2016)
[10] M. B. Shalom, M. J. Zhu, V. I. Fal’ko, A. Mishchenko, A. V. Kretinin, K. S. Novoselov, C. R. Woods, K. Watanabe, T. Taniguchi, A. K. Geim, and J. R. France, Nature Physics 12, 318 (2015)
[11] T. F. Q. Larson, L. Zhao, E. G. Arnault, M.-T. Wei, A. Seredinski, H. Li, K. Watanabe, T. Taniguchi, F. Amet, and G. Finkelstein, Nano Letters 20, 6998 (2020)
[12] I. V. Borzenets, Y. Shimazaki, G. F. Jones, M. F. Craciun, S. Russo, M. Yamamoto, and S. Tarucha, Scientific Reports 6 (2016), 10.1038/srep23051
[13] I. V. Borzenets, U. C. Coskun, H. T. Mebrahtu, Y. V. Bomze, A. I. Smirnov, and G. Finkelstein, Phys. Rev. Lett. 111, 027001 (2013)
[14] F. Amet, C. T. Ke, I. V. Borzenets, J. Wang, K. Watanabe, T. Taniguchi, R. S. Deacon, M. Yamamoto, Y. Bomze, S. Tarucha, and G. Finkelstein, Science 352, 966 (2016)
[15] M. T. Wei, A. W. Draelos, A. Seredinski, C. T. Ke, H. Li, Y. Mehta, K. Watanabe, T. Taniguchi, M. Yamamoto, S. Tarucha, G. Finkelstein, F. Amet, and I.V. Borzenets, Physical Review B 100 (2019), 10.1103/physrevb.100.121403
[16] G.-H. Lee, K.-F. Huang, D. K. Efetov, D. S. Wei, S. Hart, T. Taniguchi, K. Watanabe, A. Yacoby, and P. Kim, Nature Physics 13, 693 (2017)
[17] G.-H. Park, M. Kim, K. Watanabe, T. Taniguchi, and H.-J. Lee, Scientific Reports 7 (2017), 10.1038/s41598-017-11209-w
[18] M. Tinkham, Introduction to Superconductivity Dover Books on Physics Series (Dover Publications, 2004).
[19] T. A. Fulton and L. N. Dunkleberger, Phys. Rev. B 9, 4760 (1974).
[20] J. Clarke, A. N. Cleland, M. H. Devoret, D. Esteve, and J. M. Martinis, Science 239, 992 (1988)
[21] G.-L. Ingold, H. Grabert, and U. Eberhardt, Phys. Rev. B 50, 395 (1994)
[22] J. M. Martinis and R. L. Kautz, Phys. Rev. Lett. 63, 1507 (1989).
[23] R. L. Kautz and J. M. Martinis, Phys. Rev. B 42, 9903 (1990).
[24] Y. M. Ivanchenko and L. A. Zil’berman, JETP Lett 8, 113 (1968).
[25] Y. M. Ivanchenko and L. A. Zil’berman, Sov. Phys. JETP 28, 1272 (1969).
[26] L. Wang, I. Meric, P. Y. Huang, Q. Gao, Y. Gao, H. Tran, T. Taniguchi, K. Watanabe, L. M. Campos, D. A. Muller, J. Guo, P. Kim, J. Hone, K. L. Shepard, and C. R. Dean, Science 342, 614 (2013).
[27] V. Ambegaokar and B. I. Halperin, Phys. Rev. Lett. 22, 1364 (1969).
[28] K. S. Novoselov, D. Jiang, F. Schedin, T. J. Booth, V. V. Khotkevich, S. V. Morozov, and A. K. Geim, Proceedings of the National Academy of Sciences 102, 10451 (2005).
[29] H. Courtois, M. Meschke, J. T. Peltonen, and J. P. Pekola, Phys. Rev. Lett. 101, 067002 (2008).
[30] M. Müller, M. Bräuninger, and B. Trauzettel, Phys. Rev. Lett. 103, 196801 (2009).
[31] K. K. Likharev, Dynamics of Josephson Junctions and Circuits (Gordon and Breach, 1991).