Conductance noise in interacting Anderson insulators driven far from equilibrium

V. Orlyanchik, and Z. Ovadyahu

Racah Institute of Physics, The Hebrew University, Jerusalem 91904, Israel

Abstract

The combination of strong disorder and many-body interactions in Anderson insulators lead to a variety of intriguing non-equilibrium transport phenomena. These include slow relaxation and a variety of memory effects characteristic of glasses. Here we show that when such systems are driven with sufficiently high current, and in liquid helium bath, a peculiar type of conductance noise can be observed. This noise appears in the conductance versus time traces as downward-going spikes. The characteristic features of the spikes (such as typical width) and the threshold current at which they appear are controlled by the sample parameters. We show that this phenomenon is peculiar to hopping transport and does not exist in the diffusive regime. Observation of conductance spikes hinges also on the sample being in direct contact with the normal phase of liquid helium; when this is not the case, the noise exhibits the usual 1/f characteristics independent of the current drive. A model based on the percolative nature of hopping conductance explains why the onset of the effect is controlled by current density. It also predicts the dependence on disorder as confirmed by our experiments. To account for the role of the bath, the hopping transport model is augmented by a heuristic assumption involving nucleation of cavities in the liquid helium in which the sample is immersed. The suggested scenario is analogous to the way high-energy particles are detected in a Glaser’s bubble chamber.

I. INTRODUCTION

Conduction noise is an inherent property of essentially all electronic systems. The most common form of this noise has a 1/fα power-spectrum (PS) with α of order unity. The ubiquity of the 1/f spectrum in the noise of Fermi-gas systems is quite intriguing in that it seems to be insensitive to the specific type of transport. For example, as the metal insulator transition is crossed, the transport mode changes from diffusive to a hopping process, and this is preceded by a dramatic increase of the noise magnitude. However, the power spectrum usually retains its power-law form with α changing only slightly.

In this paper, we describe results of noise experiments performed on Anderson localized indium-oxide films measured at liquid helium temperatures, deep in the hopping regime. Several versions of indium-oxide films were employed in these studies; crystalline samples as well as several variants of amorphous indium-oxide (e.g., different carrier concentrations). Electronic transport in these systems was extensively studied in the near-equilibrium regime as well as when driven far from equilibrium. In the latter case, peculiar glassy features were found in accordance with theoretical expectations for strongly interacting Anderson insulators. In particular, when excited far from equilibrium, the conductance of the system increases, and then relaxes slowly towards its equilibrium value. This relaxation time is controlled by several factors the most important of which is the strength of the inter-electron interaction - the stronger the interaction the more sluggish is the relaxation. Increasing the static disorder (characterized, say, by the sample resistivity at a given temperature) also slows down the relaxation and so does a high magnetic field. In addition, other non-equilibrium features were observed such as aging and related memory effects. The original motivation for studying conductance noise in these systems was to get more information on their glassy behavior.

The above glassy features are restricted to the strongly localized regime and they disappear when the system crosses over to the diffusive regime by either, reducing the disorder or raising the sample temperature. In this crossover from ergodic transport to a glassy transport regime, the noise characteristics of these samples are essentially unaffected. In particular, the noise is Gaussian and has a 1/fα power spectrum with α close to 1 on both sides of the ‘ergodic-to-glassy’ crossover, again illustrating the ubiquity of this type of noise. In addition, no “saturation” is observed in the PS down to f=10−3 Hz as shown in figure 1 contrary to theoretical expectations that assumes that the noise is due to fluctuations in carrier concentration.

It is important to emphasize however that the 1/f spectrum such as shown in figure 1 is observed when the noise is measured on macroscopic samples and under linear-response conditions. On the other hand, when the drive current through the sample exceeds a threshold value, and in addition, the sample is immersed in the normal phase of liquid helium, the noise characteristics change dramatically.

As shown in figure 2 the PS becomes flat (“saturated”) up to some corner frequency f* and above it the PS drops sharply, faster than a power-law. The transition is also characterized by a considerable increases in the noise magnitude, especially around f*. Time domain sweeps (figure 3) reveal that this excess noise is associated with the appearance of downward-going spikes in the conductance whose (average) frequency increases exponentially with the drive current.

The phenomenology associated with this new type of noise is described in section III below. It is demonstrated that such current spikes may be reproduced by artificially generated cooling-bursts. We then present a heuris-
tic picture that explains how spontaneous cooling-events may arise from the interplay between the sample being driven far from linear response, and its interaction with the liquid helium bath. Specifically, the cooling events are ascribed to the production of cavities at the sample/liquid interface. These are triggered by high-energy events that intermittently appear in the current-driven sample. It is shown that the statistical occurrence of such events is a natural consequence of hopping transport in a strongly interacting system. A model, based on the percolation picture of hopping transport, shows how such events arise and predicts their occurrence probability as function of current and sample parameters. This is detailed in section IV. The model considers the charge transport in the hopping system as a traffic-flow in a network that, upon strong enough drive, results in traffic-jam events in analogy with other physical situations. Flow of particles through a disordered system often leads to traffic congestion problems resulting from the interplay between disorder and interactions. The most familiar example for this phenomenon is traffic-jams that are part of modern life in urban areas. A common form of this problem occurs when the density of cars in a one-lane road exceeds a threshold value. As a rule, this will

FIG. 1: Typical Power Spectrum of conductance fluctuations in an insulating sample, measured under linear response condition at $T=4.11\,\text{K}$. Sample: $\text{InO}_x$, length=1mm, width=1mm, $R_\square=12\,\text{M}\Omega$. The lower frequency part of the spectrum was measured by averaging 36 1-hour pieces of $G(t)$ runs and that was spliced with two spectra taken with the HP35660A up to 800Hz. The dashed line represents a 1/f power-spectrum.

FIG. 2: Conductance noise Power Spectra at different bias currents. Note the abrupt change of the noise characteristics between $I=1.9\,\mu\text{A}$ and $I=4.4\,\mu\text{A}$. The dashed line is a 1/f PS for comparison. Sample: $\text{In}_2\text{O}_3-x$, length=3.5mm, width=1mm, $R_\square=20\,\text{M}\Omega$.

FIG. 3: Conductance as a function of time in ac (using a sine drive at 23Hz, upper plot) and dc (lower plot) measurement configurations. Note the downward-going spikes in both cases. Sample: $\text{InO}_x$, length=1mm, width=1mm, $R_\square=12.5\,\text{M}\Omega$. 

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result in a ‘stop-and-go’ traffic flow although a continuous motion at the average speed of the slowest driver is theoretically possible. Granular flow is another example of a similar nature, which due to its technological implications has received wide attention. Here, we present experimental results that are consistent with traffic-jam behavior in disordered electronic systems. It is argued that this phenomenon is generic to Anderson insulators where transport is by hopping and where Coulomb interaction is significant. These are the same two ingredients that lead to electron-glass behavior. Note however that the effects associated with electron glass are measured in the linear response regime, and they actually disappear when driven too far into the non-ohmic regime. In addition, none of the glassy effects requires the presence of liquid helium for their observation. The feature that is common to the two phenomena is that hopping conductivity is a pre-requisite as both phenomena disappear in the diffusive regime.

In section V we describe a heuristic picture that purports to explain the way gaseous cavities at the sample/helium interface are generated in response to the electronic jams, and further results and discussion is given in section VI.

II. SAMPLES DESCRIPTION AND MEASUREMENT TECHNIQUES

Our experiments were performed using either crystalline or amorphous indium-oxide films (referred to in this paper as In$_2$O$_{3-x}$ and InO$_x$ respectively). These were prepared and their disorder fine-tuned to be in the insulating regime by the methods described elsewhere. Most of the samples were deposited on microscope glass-slides. However, identical results were obtained using samples deposited on alumina or Si-wafers as substrates. The samples thickness was typically 50Å, and 200Å, for In$_2$O$_{3-x}$ and InO$_x$ respectively, and various lateral dimensions of 20µm to 3cm were studied. Except when otherwise noted, all measurements reported here were done with the sample immersed in liquid helium at $T=4.11$K (under 690 Torr). The samples were mounted on a stage at the end of a probe and were loosely wrapped with a Teflon tape. This was done for protecting the sample and for securing the connecting wires. However, we found that the Teflon tape also contributed to the consistency of the spikes appearance, probably by protecting the sample from spurious thermal shocks.

Conductance measurements were made by biasing the sample with a constant voltage source while measuring the resulting current (voltage drop across a series resistor). In addition to this two-terminal configuration, we employed in some cases a four-terminal technique to verify that the measured noise originates from the sample (rather than from the contacts). Time traces and spectra were recorded by HP35660A or HP35670A Spectrum Analyzers using EG&G 5113 as pre-amplifier.

In most of the experiments the helium bath was a storage dewar, which was convenient for long measurements at a constant stable temperature. Based on previous studies using these materials we know that insulating In$_2$O$_{3-x}$ films exhibit the Mott’s VRH law: $R(T) \propto \exp \left( \frac{C}{T^2} \right)$. For samples with $R_{\infty}$ between 10MΩ and 10GΩ at 4K, which is the range studied here, $T_0$ was in the range of 3000-8000K. The hopping law in the amorphous films depends, among other things, on the carrier concentration $n$. Samples with $n > 10^{20} \text{cm}^{-3}$ (a range of $n$ which exhibited spikes most clearly), tend to exhibit $R(T) \propto \exp \left( \frac{C}{T} \right)$ with $T^*$ ranging between 150 to 600K for the above range in $R_{\infty}$.

III. III - THE BASIC FEATURES OF THE PHENOMENON

The phenomenon we wish to focus can be observed in the time dependence of the sample conductance. This is illustrated in figures 4 and 5 for typical In$_2$O$_{3-x}$ and InO$_x$ samples respectively. For sufficiently small drive current $I_T$, the conductance $G$ versus time shows only small fluctuations, which, upon Fourier transforming turn out to exhibit the common 1/f noise power-spectrum. However, once $I_T$ exceeds a threshold value, a sudden change in the time traces is observed as illustrated in figures 4 and 5.

Note that above a certain $I_T$ the conductance versus time traces $G(t)$ acquire an additional component in the form of downward-going spikes that are rare near the threshold $I_T$ but their number per unit time increases dramatically with current. A mere 3-4% increase in $I_T$ past the threshold typically increases the spikes frequency by two orders of magnitude (e.g., figures 4&5 and the discussion regarding this issue in section V).

Spikes do not appear (in the experimental time-window of up to 100sec.) unless the sample is driven with sufficiently high current. As a rule, this occurs when the sample is far into the non-ohmic regime as shown in figure 6. The first observation is then that this is manifestly a far-from-linear response phenomenon.

In the majority of samples, (more than 60 studied here) the spikes had the same generic shape; a fast “attack” followed by a longer “recovery-tail”. The relative magnitude of the spike and its duration however, are dependent on specific sample parameters as discussed later. Naturally, the appearance of these spikes modifies the noise characteristics of the signal in terms of both, magnitude and spectrum: The power spectrum turns out to be flat (“white-noise”) at low frequencies, followed by a fast decline above a certain frequency, both features are obviously related to the peculiar shape of the individual spike. In many cases, the spikes had the same unique shape (as, e.g., in figure 4) namely; each spike is an exact replica of any other spike. In most other instances, two or three different shapes could be identified. The detailed shape
of a particular spike is sample specific; Samples from the same batch with nearly identical $R_\parallel$ and $I-V$ characteristics can be distinguished by their spike shape.

These remarks apply to the immediate vicinity of the threshold where individual events could be resolved in the time domain, which was an extremely narrow range of bias current. At higher currents, the power spectrum change with $I_T$ indicates that the spikes duration may become shorter (c.f., the shift of $f^*$ to higher frequencies in figure 2).

Figure 7 shows time traces at different currents employing a 4-terminal measurement on an insulating In$_2$O$_{3-x}$ sample. The spikes in this case exhibit the same shape as those in figures 4 and 5, and their number per unit time increase with the drive current in the same manner as before. However, these spikes are now upward-going. One may therefore conclude that the events associated with the spikes are temporary drops of the sample conductance occurring in the bulk of the film (as opposed to contacts problems).

Similar spikes and associated power spectrum had been observed in flow of granular particles through a pipe. We believe that this similarity is not coincidental - both phenomena involve an intermittent dropout in traffic-throughput due to a forced-flow of strongly interacting, discrete particles. However, the physical processes that lead to the observation of conductance spikes are more intricate than the one-dimensional granular flow problem. The complex nature of the this phenomenon may be appreciated by considering the following key observations:

1. The bath in which the sample is immersed plays a crucial role. Spikes do not appear unless the sample is in contact with liquid helium. Also, no spikes are observed when the sample is vacuum-loaded (while being kept at $T=4.11$K via a copper cold-finger). In addition, coating the sample with a thick layer of a photo-resist film eliminated the spikes even when the sample was immersed in liquid helium. A more dramatic demonstration of the role of the bath is illustrated in figure 8. This shows that the spikes vanish abruptly upon cooling the helium bath below $T=2.174$K. Note that this temperature is the 'λ-point' below which the helium becomes a superfluid.

2. The phenomenon is peculiar to samples that are deep into the hopping regime having sheet resistance of 1MΩ to 10GΩ. No spikes appeared in any of our In$_2$O$_{3-x}$ samples that were in the diffusive regime (i.e., when the sample conductance $G$ is higher than the quantum-value

FIG. 4: Conductance fluctuations as a function of time for a typical In$_2$O$_{3-x}$ sample measured at different bias currents (the current values in µA are shown for each trace) near the threshold. Sample length=1.1mm, width=3mm, $R_\parallel=68$MΩ (at the threshold drive).

FIG. 5: Conductance fluctuations as a function of time for a typical InO$_x$ sample measured at different bias currents (the current values in µA are shown for each trace) near the threshold. Sample: length=0.4mm, width=2mm, thickness=175Å, $R_\parallel=75$MΩ (at the threshold drive).
FIG. 6: Resistance as a function of the bias field for amorphous (InO$_x$-solid circles, length=0.5mm, width=5mm) and crystalline (In$_2$O$_{3-x}$-open squares, length=0.15mm, width=1mm) samples. The arrows mark the threshold for the appearance of spikes.

FIG. 7: Four probe measurement of resistance as a function of time below and above the threshold current (the current values in µA are shown for each trace). Sample: In$_2$O$_{3-x}$, length=0.15mm, width=15mm, $R_{\square}=3.7\Omega$ (at the threshold drive).

FIG. 8: $\delta G(t)$ for different temperatures near the $\lambda$-point. Note the abrupt disappearance of the spikes at the $\lambda$-point. Sample: In$_2$O$_{3-x}$, length=0.5mm, width=0.5mm, $R_{\square}=16\Omega$ (at the threshold drive).

$e^2/h$ even at currents that were so high as to cause helium boiling (six orders of magnitude higher power than that used in figures 4 or 5). Therefore, the mode of transport in the sample plays an important role in giving rise to the conductance spikes.

These empirical observations suggest that the spikes involve interplay between the sample and the liquid bath, and in the following, we attempt to elucidate their respective role.

The disappearance of the spikes at the transition into the superfluid phase led us to suspect that the spikes may be associated with thermal events, which are then annihilated (or considerably weakened) by the superfluid component kicking-in as the temperature falls below the $\lambda$-point. To account for the observed downward-going spikes, these have to be cooling-events given the fact that the conductance in the hopping regime increases with temperature (namely, $dG/dT > 0$). The fractional change of the conductance $\Delta G/G$ associated with a spike is 0.01-0.2% which could be affected by a cooling-burst of 0.1-2mK for a typical sample. To check on such a cooling conjecture, several samples were subjected to cooling-bursts produced by cutting-off the power, for brief period ($\approx 2.5\text{mS}$), from an external heat source to which the sample was otherwise constantly exposed. Typical results are illustrated in figure 9 using two different ways to create cooling events. Adjusting the duty-cycle such that power into the heat source is applied only during 2.5mS intervals resulted instead in upward-going spikes (figure 9). The heat source in these examples was ei-
FIG. 9: Spikes caused by artificial cooling and heating events: (a) IR LED cooling; (b) IR LED heating; (c) Micro-resistor heating; (d) Micro-resistor cooling. The long conductance tail observed after the IR excitation in (b) is an inherent glassy effect; This slow relaxation occurs whenever the sample is excited with sufficiently high quantum energy source (c.f., Ben-Chorin et al in reference 6). Note the absence of such tail in plate c. Samples: (a) and (b) InO$_x$, length=0.6mm, width=1.5mm, $R_{\square}$=9.4MΩ; (c) and (d) InO$_x$, length=2.4mm, width=1.5mm, $R_{\square}$=20.1MΩ. In these measurements, the conductance was monitored under linear response conditions.

ther a micro-resistor or small IR light-emitting diode coupled thermally to the sample. The micro-resistor was a 350Å gold film 20µm wide and 10mm long deposited on a mylar film which was thermally anchored to the sample.

Note that, in response to thermal shocks, current-spikes are generated in either case, and they are similar in shape to those produced ‘naturally’ in that they mimic the fast “attack” and slow recovery-tail form of the natural spikes (c.f., figure 10). However, independent of the method by which they are produced, the artificial spikes were much shorter than the ‘natural’ ones typically are. In particular, the recovery tail of the artificial spikes was always shorter than 3mS, independent of the magnitude of the thermal shock or its sign. This 3mS is presumably the time set by the thermal-inertia of the combination of the substrate and sample-stage. By contrast, the duration of spikes produced by above-threshold currents depend on sample parameters, most notably on its carrier concentration $n$, and could be as large as 200mS. An example for the dependence of the spike duration $\tau$ on $n$ is illustrated in figure 11. Note that for the sample with the largest $n$, $\tau$ exceeds the duration of the artificial cooling-spikes by almost two orders of magnitude.

We shall argue below that the cooling events result from the formation of (gaseous) cavities at the sample/helium interface. In this picture, the long relaxation tail is associated with the re-liquefaction of the cavities. Before addressing these issues in more detail however, we need to consider the mechanism by which cavities are generated in the first place and in particular the role of the sample parameters in triggering these cavities.

Led by the observation that spikes appear only when the system is in the hopping regime (point 2 above), we review in the next section some of the salient features of hopping conduction that distinguish it from diffusive transport using the conventional percolation picture and taking into account the discreteness of the hopping process. This treatment reveals the natural reason for crossover behavior at a threshold current density; a feature required by one of the empirical observations alluded to above. Moreover, the dependence on the sample parameters agrees with the consequences of this scenario, which also accounts for the exponential dependence of the spikes frequency on current.
FIG. 11: Spike shapes for three InO$_x$ samples with different carrier concentration. Note the difference in the spike duration $\tau$.

IV. IV – SOME RELEVANT CONSEQUENCES OF HOPPING CONDUCTIVITY

The system we are dealing with, an Anderson insulator, is a degenerate Fermi gas with spatial disorder sufficiently strong to cause localization of the wave functions. Namely, the amplitude of the wave function is appreciably only around a certain point in space $r_0$ and decays as $\exp\left[-\frac{|r-r_0|}{\xi}\right]$ away from it. Charge transport in such a system is controlled by the quantum mechanical transition probability between localized states. The transition rate between sites $i$ and $j$ is given by:

$$\omega_{i,j} = \omega_0 \exp\left[-\frac{r_{ij}}{\xi} - \frac{\Delta E_{ij}}{k_B T}\right]$$  \hspace{1cm} (1)

Here $\omega_0$ is an attempt frequency (typically, $\omega_0 \approx 10^{12}-10^{13}$ sec$^{-1}$), $\xi$ is the localization radius, $r_{ij}$ is the sites spatial separation, and $\Delta E_{ij}$ is their energy difference.

Since the system is disordered, $r_{ij}$ and $\Delta E_{ij}$ are random variables distributed over some range. The exponential dependence of $\omega_{i,j}$ on these variables (equation 1) leads to an extremely wide distribution of transition rates. The macroscopic system may then be viewed as a random-resistor-network in which each pair of sites $i, j$ is connected by a Miller-Abrahams resistor

$$R_{ij} \propto \omega_{i,j}^{-1}$$

The wide distribution of the $R_{ij}$’s in the random-resistor-network leads to several unique features of electronic transport in such a medium. The most familiar feature is that the current in the system is carried by a percolation-network (CCN) and form pockets that are effectively “dead-wood” and some dangling-off branches (“dead-ends”) that do not contribute to dc current. This is illustrated in figure 12. Each branch of the CCN is composed of a series of $R_{ij}$’s with different values. The current through each branch is controlled by the largest $R_{ij}$ that percolates in the system (being inter-connected by smaller resistors). This is called the critical-resistor $R_C$ or the “bottleneck”-resistor. The sheet resistance $R_{0}$ of a macroscopic two-dimensional sample is of the order of $R_C$. Upon a change of temperature, electric field, magnetic field etc., the $R_{ij}$’s of the system will in general change too, resulting in a modified CCN, with different $R_C$ and $R_{ij}$. The typical distance between $R_C$’s on neighboring branches is called the percolation radius $L_C$ and it is the measure of the CCN mesh-size.

The other feature that distinguishes hopping transport from that of metallic conductivity is its discrete nature. The electron spends most of the time in a localized state attempting to cross an effective barrier controlled by the factor $\frac{r_{ij}}{\xi} + \frac{\Delta E_{ij}}{k_B T}$ (which is typically $\gg 1$) with an exponentially small probability while the actual transition-time is very short. In other words, on a microscopic scale, charge motion is a ‘stop-and-go’ process marking the electron as a “bad driver”. A corollary of this picture is the emergence of a “critical current” $i_c$ related to the transition rate through a bottleneck resistor. Using $e$ for the electronic charge, $i_c$ may be expressed by:

$$i_c = e \cdot \omega_0 \cdot \exp[-X_C] \text{ where } X_C = \frac{r_{ij}}{\xi} + \frac{\Delta E_{ij}}{k_B T}$$  \hspace{1cm} (2)

To exceed $i_c$ an electron must be pushed towards site $i$ before the electron already in this site has a chance to hop to site $j$. This situation will be referred to as a jamming event.

FIG. 12: Schematic representation of the current carrying network in an Anderson insulator. The lines are the current carrying paths. Stars represent "bottleneck" resistors in the network.
increases at a larger rate than the rate of the production of new roads, leading to the annoying traffic-jam during rush hours.

The total current at which the likelihood of a jam event becomes appreciable can be estimated as follows. Assume a two-dimensional system, namely, a film of thickness $d \ll L_C$, and width $W$. The critical current $i_c$ through a typical branch will then be reached when $I_T \approx i_c \cdot W/L_C$. Using equation 2 and noting that $R_\square \approx R_C \approx R_0 \cdot \exp[|X_C|]$, this happens when the (2D) current density $J_C^S$ is $\approx \frac{\sigma_{\text{bulk}}}{R_C L_C}$. It is often found empirically $\delta \approx \frac{\tau}{\kappa}$ in two-dimensional hopping systems that $R_0$ is of the order of the quantum-resistance $\frac{\hbar}{e^2}$. Using this relation one gets:

$$J_C^S = \frac{\hbar \nu}{e \kappa L_C} \approx \frac{\hbar \nu}{e \kappa L_C}$$

A jamming event is then likely to occur once the current density through the system exceeds a value, which is inversely proportional to the sheet-resistance $R_C$. Alternatively, the threshold condition can be expressed as a critical voltage drop $V_c$ across a bottleneck resistor:

$$V_c = \omega_0 R_0 \approx \frac{\hbar \nu}{e \kappa L_C}$$

We now show that the onset for the conductance-spike in our samples, as well as their frequency as function of the local voltage (or the global current-density), follow the behavior expected by the above considerations. First, consider the correlation between the threshold-current density $J_C^S$ at which the spikes appear and the sample $R_C$, shown in figure 14. This figure summarizes results of more than 40 In$_2$O$_{3-x}$ and InO$_x$ samples with vastly different geometries; in particular, it includes samples with width $W$ in the range 20$\mu$m to 3cm. This three orders of magnitude spread should be compared with a mere factor of ±4 scatter in the data in the figure. Thus, the correlation between $J_C^S$ and $R_C$ is quite suggestive. The dashed line in the figure is calculated by equation 3 using $\omega_0 = 6 \cdot 10^{12}$ sec$^{-1}$ and $L_C=1\mu$m, both are reasonable values for a hopping system. The use of a disorder-independent percolation radius needs justification. For a Mott VRH in 2D (which is the case for the In$_2$O$_{3-x}$ samples) one expects $L_C$ to scale as $\left[\frac{r}{(\frac{\tau}{\kappa})^{\frac{1}{2}}} \right]^\nu$, where $\nu$ is of order $128$. This would then lead to $L_C$ going like $\xi \cdot \left(\frac{\tau}{\kappa}\right)^{\frac{1}{2}}$ and combined with $T_0 \propto \xi^{-2}$, and at given temperature, $L_C \propto \xi^{-\nu}$. For the range of $R_C$ studied here (10MΩ-1GΩ), the associated $\xi$’s range from 20Å to 5Å respectively, which means that the variation of $L_C$ is a mere factor of 1.6, which is smaller than the scatter in the data in figure 14. Even this factor is compressed by the voltage; Note (equation 4) that the threshold is characterized by a constant local voltage. Since $L_C$ is reduced by voltage, the more so the larger ‘equilibrium $L_C$’ is, the effective range of $L_C$ is smaller than the estimate based on near-equilibrium conditions. It seems harder to understand why the results for the InO$_x$ samples do not fall on the same line as those of the crystalline samples. As mentioned above (section II), InO$_x$ samples (especially those with high $n$) show a different
hopping law than In$_2$O$_{3-x}$ samples of similar $R_{\square}$. There is thus no a-priori reason to assume that In$_2$O$_{3-x}$ and InO$_x$ samples of equal resistance should have the same $L_C$. Nevertheless, independent estimates of $L_C$ in these systems appear to be in the same range of values. For In$_2$O$_{3-x}$ films in this range of disorder, $L_C$ values ranging between 0.3 to 0.5 $\mu$m were found based on the magnitude of the conductance fluctuations as a function of sample size. In high carrier concentration InO$_x$ films, a value of 0.3 $\mu$m $L_C$ may be estimated by extrapolating the results of Frydman et al. The latter are based on measuring the resistance versus thickness. Interestingly, a similar value for $L_C$ of 1 $\mu$m was estimated based on the magnitude of mesoscopic fluctuations for insulating films of granular nickel, a system that differs markedly from both, crystalline and amorphous indium-oxide.

The local voltages $V_c$ associated with the data points in figure 14 can be estimated using the expression: $\frac{I_T R_{\square} L_C}{W}$ and they vary in the range 12-70mV while the sheet resistance $R_{\square}$ (at threshold) changes by more than two orders of magnitude. This range for $V_c$ compares favorably with equation 4 that for $\omega_0 = 6 \cdot 10^{12}$ sec$^{-1}$ yields $V_c = 25$ mV.

Furthermore, the same $V_c$ seems to be in control of the dependence of the spikes (average) frequency $\omega$ on the current near threshold shown above (figures 4 and 5). The extremely fast $\omega(I_T)$ is suggestive of an activated process, which turns out to be of the Arrhenius type. This is illustrated in figure 15 for several of the studied samples where we plot $\omega$ as function of $V = \frac{I_T R_{\square} L_C}{W}$ (normalized by the respective $V_c$ for each sample). It is tempting to analyze these results using a relation of the form:

$$\omega = \omega_0 \exp \left[ -\frac{e(V - V_c)}{k_B T^*} \right]$$

(5)

with the rationale that the local voltage $V$ increases the probability to cause a jam by reducing the barrier $V_c$. Note that the $T^*$ used in equation 5 is unlikely to be the same as the bath temperature $T$ in the present case. Recall that we are dealing here with a far from equilibrium situation and there is no unique way to estimate an effective temperature $T^*$. As remarked above, the spikes occur when the field $F$ brings the system near (but somewhat below) the activationless conductance regime where the ‘temperature’ of the electrons is controlled by $F$. For the lack of anything better, it seems plausible to estimate $T^*$ from the resistance at the threshold field by comparing it to the ‘ohmic’ $R(T)$ data (namely, using the resistance of the sample as a ‘thermometer’). For the samples used in figure 15 we get $k_B T^*$ in the range 0.5meV to 0.9meV and using these data with $\omega_0 = 6 \cdot 10^{12}$ sec$^{-1}$ in equation 5, gives a range for $V_c$ of 25-100mV. Note that these values of $V_c$ overlap with the range of energies associated with $\delta \varepsilon = \frac{e^2}{\kappa r}$ discussed above.

There are then several independent indications that the jam-scenario may be involved in controlling the onset as well as the frequency of appearance of the conductance spikes.

The question is what is then the role of the He bath:
FIG. 16: A schematic description of the localized states in the CCN near a bottleneck resistor (that for $i \ll i_c$ is assumed to be $i, j$), for small currents ($i \ll i_c$ – lower configuration) and for a typical jam event ($i \geq i_c$ – upper configuration). Light circles are unoccupied sites and dark circles are occupied sites. δε is symbolized here as the sum of the displacements of the sites $k$ and $i$ from the baseline.

Why the electronic mechanism does not generate spikes unless the sample is in contact with liquid helium. The answer is – it presumably does. However, the electronic processes that take place due to a traffic-jam should usually die out very quickly and escape observation. Consider that a jamming event has occurred at some point in the sample. The population of electronic states near a jammed site may then look as in figure 16. Note that, for $i \geq i_c$ an effective energy barrier $\delta \varepsilon$ is created due to the repulsion between the electrons occupying sites $k$ and $i$. Such a configuration is a high-energy state and it will quickly dissipate itself, say, by the electron in $i$ hopping away to another site. The transition time associated with a single-particle hop is quite short. This may be estimated from $\omega^{-1} \approx \frac{\omega_D^{-1} R_{33} e^2}{\hbar}$ to be $\approx 10^{-9}$ sec for $R_{33} \approx 100 \text{M} \Omega$. Therefore, the duration of an electronic-jam ought to be much shorter than the spike widths we typically observe (c.f., figure 11). The latter therefore must involve another agent such as the one suggested next.

V. A MECHANISM FOR COOLING EVENTS TRIGGERED BY THE ELECTRONIC TRAFFIC-JAMS

We suggest that the spikes are the response of the sample to cooling events associated with nucleation of He cavities at the sample/liquid interface. These are triggered by the electronic jams, which thus control the onset and frequency of the events, while the duration of a spike is determined by the lifetime of the formed cavities. Gaseous helium cavities may be produced by heat-pulses, or by acoustic-wave-bursts, and in particular, the threshold for a cavity formation on glass surfaces (heterogeneous cavitation) is quite low. Both heating-pulses and acoustic-waves are potential products of energy-bursts involved in the jam events. That a brief heat-pulse may emanate from a jam event seems obvious. Emission of acoustic wave needs more elaboration. To see how this might happen note that when a jam is “on” an electric field appears for a brief moment across a bottleneck resistor. This field is of the order of $10^5$ V/cm, which is quite large. While this field is on, the medium will polarize to some degree. This in turn will be accompanied by a mechanical deformation (electrostriction). When the jam goes “off” the stress on the medium is relieved, and this cycle of ‘pull-and-let-go’ is very likely to generate acoustic waves much like plucking a string does. This effect is possible since the “on” time of the jam is longer than $\omega_D^{-1}$ ($\omega_D$ is the Debye frequency of the medium which is of the order of the range of the $10^{12}-10^{15}$ sec$^{-1}$.) even for $R_C = 1 \text{M} \Omega$. It is hard to tell which of these two agents is more dominant in our experiments and perhaps they are complementary. In fact, our scenario is somewhat analogous to the way Glaser’s bubble chamber detects high-energy particles. The operation of the bubble chamber involves super-heating the liquid followed by a sudden drop of pressure to make visible the trajectory of a fast moving particle. Super-heating in our picture is facilitated by the accentuated local heating inherent to the inhomogeneous mode of hopping transport. The local temperature near a bottleneck resistor (on scale of a thermal phonon wavelength of $1000 \text{A}$) could conceivably be a fraction of a degree $K$ above the temperature of the non-conducting part of the sample. The effective barrier for cavitation will be lower at these hot spots and bubble nucleation will most likely start at one of these spots. Below the $\lambda-$ point, such hot spots will be eliminated by the superfluid counter-flow and this presumably is the main reason for the disappearance of the spikes (c.f., figure 8).

Once a cavity nucleates and starts to expand, it sucks energy from the medium (the latent heat of transforming liquid into gas) thus producing local cooling. This is detected as a drop of the conductance – the fast “attack” of the spike. The trailing edge of the spike is controlled by the re-condensation of the gaseous cavity. Hence, the duration $\tau$ of a spike depends on the volume of the cavity. A simple estimate for re-condensation of a bubble with diameter $D$ in cm yields $\tau \approx 10^8 D^3$ sec. To account for the spike duration-times observed in figure 11 (varying in the range 1-200mS) the respective cavity diameters has to be in the range of 2-10µm. Even the largest bubbles in this range are too small to be optically imaged by our current techniques yet they possess large enough thermal-mass to sustain the cooling effect for a long time.

Further support for the relevance of cavitation is the dependence of the spike frequency on the vapor pressure of the helium bath. As illustrated in figure 17a, the average spike frequency $\omega$ decreases with pressure $P$. This is consistent with the assumption that cavities are facilitated by the negative pressure of an acoustic wave. To keep $\omega$ constant when $P$ increases $I_T$ has
FIG. 17: (a) The frequency of the spikes as a function of the vapor pressure of the helium bath. (b) The relative magnitude of the spikes as a function of the vapor pressure of the helium bath. This experiment was done by pressurizing the storage dewar, and monitoring the results ‘on the fly’. The entire experiment took about 7 minutes to accomplish, during which the temperature rise was negligible as indicated by the change in $G$. $\Delta P=0$ is related to the atmospheric pressure (693Torr). Sample: InO$_x$, length=0.7mm, width=1mm, $R_{\square}=1.7M\Omega$ (at the threshold drive).

However, under these conditions, the magnitude of the spikes decreases extremely fast with $P$ as shown in figure 17b. Over the range of $\Delta P$ shown in the figure, the sample conductance changes by a negligible amount, and by applying artificial cooling-bursts in this region it was found that the sample’s sensitivity as a bolometer is independent of $P$. It must therefore be concluded that the decrease of $\Delta G/G$ with pressure results from the diminished cooling power of the cavities. We interpret this decrease as indication that the heat of vaporization occurs near the region where the ambient temperature is near the critical temperature.

We have considered another scenario for cooling, based on the notion that bubbles, produced by local Joule-heating, grow and eventually float up thereby the sample is re-cooled by helium counter flow. There are a number of problems with such a scenario. To leave the surface of the sample a gaseous bubble has to be rather big, of the order of few hundred microns, and the energy required to produce it at the rate observed is incompatible with the input Joule-energy and clearly inconsistent with the exponential dependence of this rate. In addition, we made a number of tests to explore the possibility that bubbles may form to be big enough do leave the sample. Two of these are described below.

In the first experiment, two samples were lowered into the helium bath and were attached to the probe parallel to one another and to the floor. The lower one was biased with current so that it produced spikes. The upper sample (3-4mm above the “active” one) was biased with sub-threshold current and showed no hint of any thermal shock that ought to have been recorded had a gaseous bubble emanating from the lower sample hit its surface.

In another experiment, an electromechanical “woodpecker” was employed with the idea that mechanical shocks may dislodge some “ripe” bubble and thus synchronize spike appearance. The device used was a small loudspeaker cone to which a small plastic beak was glued and the rig was attached on top of the probe such that it could periodically transmit a mechanical kick down to the sample. The latter was biased with near threshold current that naturally produced occasional spikes. Various combinations of the “woodpecker” frequency and amplitude were tried as well as several choices of bias current. No correlation was found between the mechanical shocks and the spikes appearance. The negative results of these experiments give further reasons to rule out the cooling mechanism based on “bubble-emission”.

VI. VI - FURTHER RESULTS AND DISCUSSION

To shed more light on the cooling events and the way they are produced we performed the following experiments. In the first, the idea was to further check on the thermal nature of the effect by using a sample (biased with a sub-threshold current) as a bolometer to detect the pulse that, by assumption emanates from the ‘active’ sample. For that purpose, two samples were deposited on the same substrate as two cross strips separated by a thick insulating barrier (5000Å of SiO$_2$). As shown in figure 18 the ‘bolometer’ records a signal that is synchronized with and mimics the main features of the spike observed in the $G(t)$ of the active sample. The magnitude of the signal at the detector sample is considerably reduced (note that the sensitivity of the ‘bolometer’ given by $dG/dT$ is greater than that of the active sample because it is used with a smaller bias current). In addition, the spike appears somewhat delayed and broader relative to the original, which is consistent with a diffusive propagation process.

While the appearance of the spikes is not reflected in the (averaged) $R$ vs. $F$ plots (c.f., figures 6 & 13), it is readily observed in derivative measurements. The advantage of this type of measurement is that it makes it possible to check for hysteresis in the threshold current.
by changing the sweep rate. We studied the differential conductance \( \frac{dI}{dV} \) of some samples as a function of the bias voltage \( V \) focusing on a limited range straddling the threshold for spikes appearance. A typical \( \frac{dI}{dV}(V) \) plot is shown for a In\(_2\)O\(_3-x\) sample in figure 19. Note that at \( V \approx V^* \), where \( V^* \) is the threshold value for spike appearance, \( \frac{dI}{dV} \) exhibits a sharp change and above \( V^* \) the conductance is smaller than the theoretical curve obtained as follows. As shown in figure 13, the spikes in this sample appear in the voltage regime somewhat below the transition to activationless hopping. The current-voltage characteristics in 2D take the form:

\[
\log[I] \propto -\left(\frac{V}{V_0}\right)^{\frac{1}{3}} \tag{6}
\]

from which the differential conductance is obtained as:

\[
\frac{dI}{dV} \propto \exp[-(\frac{V}{V_0})^{\frac{1}{3}}] \cdot V^{-\frac{4}{3}} \tag{7}
\]

this is plotted in figure 19 with the value of \( V_0 \) extracted by fitting equation 6 to the data in figure 13. The difference between this theoretical curve and the experimental \( \frac{dI}{dV} \) is 2.1-3.4\% while the maximum amplitude of the spikes for this sample is 0.02\%. In other samples, we noticed that \( \frac{dI}{dV} \) above threshold increase at a slower rate than the natural trend observed below \( V^* \) leading in some samples to even bigger discrepancies. Therefore, the diminished conductance above threshold cannot be simply explained by the fact that the spikes (that are associated with negative \( dI \)) “pull-down” the signal. The reason for this discrepancy might be due to the way the cavitation processes, that are fed by the applied field, interfere with the photon-assisted-hopping processes involving the same field. This problem is currently under further investigation.

As a further test of the conjecture that the spike phenomenon involves hopping transport, we made a preliminary study of the effect of a magnetic field on the spikes. For this purpose, we chose to work with the crystalline version of indium-oxide that has been extensively studied and where the magneto-resistance (MR) is well understood as being due to orbital and spin effects. The sign of the MR is negative or positive when the dominant mechanism is quantum-interference effect (orbital) or spin-alignment (isotropic) respectively. The relative contribution of the orbital part can be controlled by changing the angle between the sample plane and the field direction. It is thus possible to compare the results at high field \( H_1 \) with the results at small field \( H_2 \) while adjusting things such that \( R_{\square}(H_1) = R_{\square}(H_2) \). Note that while the sample is the same from the point of view of the measurement circuit, it must be microscopically different; both configurations have similar bottleneck resistors \( R_C \)’s but at least some of these are placed at different locations relative to the sample axes. In

FIG. 18: \( \delta G(t) \) generated by the active sample (open squares) and recorded by the detecting sample (solid circles) (see text for details). Samples: generator - InO\(_x\), length=2.4mm, width=0.7mm, \( R_{\square} = 0.8\)M\(\Omega \) (at the threshold drive); detector - InO\(_x\), length=1.5mm, width=1mm, \( R_{\square} = 6.7\)G\(\Omega \).

FIG. 19: Dynamic conductance \( (dI/dV) \) versus voltage over a region including the threshold for spikes appearance. Open and solid triangles represent data taken upon sweeping the applied bias \( V \) up and down respectively. Sweep rate is 0.08V/sec. Note the small hysteresis. The dashed line shows the differential conductance expected by equation 7. Note that the experimental \( dI/dV \) is lower than the theoretical value even at voltages much larger than the voltage where individual spikes cannot be resolved in the time domain. This sample is the same as in figure 13.
other words, the current carrying network must be different. To see why that must be so it is enough to note that the spin-alignment mechanism eliminates from the CCN doubly occupied states, a process that is complete for $g\mu_B H > k_B T$. The result of this magnetic field stratagem is illustrated in figure 20 for one of the two samples studied so far. In this particular case, the 9T field has presumably shifted the nucleation-site of the cavity resulting in a different looking spike. This is a plausible scenario; the hottest spot on the sample would be the preferred site for cavity nucleation and this spot is probably at or near one of the critical resistors in the CCN. Therefore, when the CCN is changed by the field the spike might originate at a different location. And exhibit somewhat different shape much like when using another sample from the same batch.

As noted in section III, in quite a few cases, the spikes near threshold seem to have an identical shape. This “mesoscopic” behavior in macroscopic samples is intriguing: Both, $L_C$ and the cavities diameter, which are the relevant parameters, are much smaller than the size of the sample (which is typically 1mm across). Several factors may be involved in bringing about this behavior. In our model, it was implicitly assumed that the bottleneck resistors in the CNN are all equal. This may be a reasonable assumption for assessing the threshold current but the $R_C$’s are probably distributed over a considerable range, especially when the disorder is large. The jam events near threshold presumably involve the $R_C$’s at the far tail of this distribution. The question is how many ‘active’ $R_C$, are there over the ≈3% range of current, where individual spikes are resolved in 15-20% of the samples. Alternatively, there may be a ‘preferred’ site for cavity nucleation, which is triggered by any of a number of (simultaneous) jam events. In other words, we see no compelling reason that the cavity must be nucleated at the generating $R_C$ site though probably not too far from it.

An important issue that needs further elucidation is what determines the spike duration in a given sample. The spike duration $\tau$ and the relative amplitude of the spikes $\delta G/G$ are usually correlated as illustrated in figure 21. These results were obtained using InO$_x$ samples where $n$ could be changed over a large range by varying the In/O ratio. To allow comparison, all samples in this figure have similar dimensions of 1x1mm. No correlation between $\tau$ and disorder of a given sample was found. Note that longer $\tau$’s are statistically accompanied by bigger $\delta G/G$, and both decrease rapidly below a carrier-concentration of $\approx 6 \cdot 10^{19}$cm$^{-3}$. The correlation between $\tau$ and $\delta G/G$ is natural in our scenario; larger cavities presumably produce more cooling and last longer. The dependence of both on $n$ however is more intriguing in that it suggests that the phenomenon may disappear in the limit of small carrier concentration. Indeed, in samples with $n<5\cdot 10^{19}$cm$^{-3}$ the spikes are so
features of systems that were recently shown to exhibit electron-glass behavior, we tested a granular aluminum sample, a hopping mechanism that resembles a phase transition. The spikes appear and disappear with conductance spikes. In some respects, the phenomenon resembles a real (non-equilibrium) phase transition than what has been termed here a traffic jam.

A relevant question is the role of many-particle transitions that are expected to be important in a hopping system with strong electron-electron interactions. This aspect of hopping transport was omitted from our treatment on the assumption that many particle transitions become less important in the non-ohmic regime. Obviously, transport by many particles, correlated transitions could obviate the traffic jam events that are depicted here as resulting from single-particle transitions. In principle, extending the study of the dependence of the spike frequency on drive current to the low frequency regime might shed some light on this question.

In summary, we presented data on a new kind of conductance noise characterized by the appearance of downward-going spikes in the conductance versus time traces. This was shown to be a non-equilibrium phenomenon peculiar to the hopping regime. A model taking into account the specific features of hopping transport is shown to be consistent with several aspects of the phenomenon. In particular, the model explains the observation that the onset for the spike appearance and their average frequency is controlled by the current density. It also predicts certain dependencies on sample parameters in agreement with experiments on many different samples. The purely electronic consequences of this electronic-jam picture are expected to occur at microwave frequencies and would be difficult to observe. We ascribe the spikes to cooling events produced by cavitation phenomena. In our scenario, cavities are nucleated in response to pulses of energy (heat and/or acoustic waves) generated and synchronized by the electronic jam events. The time scales associated with the recondensation of the cavities is long enough to facilitate observation of individual events in much the same vein that bubble chambers are used in tracking fast moving elementary particles.

FIG. 22: $\delta G(t)$ for a granular Al sample measured at different currents near the threshold (the current values, in $\mu$A, are indicated for each trace). Sample: length=1mm, width=1mm, thickness=200Å, $R_\square=33.7\,\text{M}\Omega$ (at the threshold drive).

short and have so small relative magnitude as to make their detection very difficult as compared with the prominent spikes produced by samples with larger $n$ (c.f., figure 11). A plausible way to understand this trend is to note that samples with larger $n$ are more likely to produce cavities due to their stronger electron-electron interactions; An Anderson insulator lacks metallic screening, which in turn means that the inter-particle interaction is stronger in a system with higher carrier concentration. Hence, such a system has more "kick" to produce a bigger cavity by the mechanism described above.

This mechanism for spike production is generic; it should apply to any system provided that two conditions are met; a) transport is by hopping, and b) interactions are strong. Therefore, natural candidates for such behavior are electron-glasses that are glassy by virtue of the very same two ingredients. Led by this consideration, we tested a granular aluminum sample, a hopping system that was recently shown to exhibit electron-glass features. The results of such an experiment are shown in figure 22 demonstrating conductance spikes and bias dependence that are similar to the respective behavior in crystalline and amorphous indium-oxide films. It would be of interest to test our picture in other hopping systems whether or not they exhibit glassy features, and in other dimensionalities as well.

Finally, we wish to comment on the nature of the transition that is characterized by the appearance of the conductance spikes. In some respects, the phenomenon resembles a phase transition. The spikes appear and disappear rather suddenly over an extremely narrow range of current, and in several cases, hysteretic behavior was observed. It should however be noted that the underlying picture of traffic jam is essentially an extremely sharp crossover rather than a phase transition. The fact that a threshold current is assigned to a given sample is merely a result of the exponential dependence expressed by equation 5 and the finite time-window of the observation. The hysteretic behavior is probably related to the formation of cavities, a process that is more likely to be a real (non-equilibrium) phase transition than what has been termed here a traffic jam.

From the point of view of cavitation phenomena it is worth noting the similarity of our picture and the work of Sinha et al. Sinha et al used in their experiment electrical pulses fed into a Bi crystal to generate cavities in the helium. The Bi sample doubled as a bolometer detecting the thermal effects that resulted from cavitation, just as our samples are presumed to do. Their process differs from ours mainly in that in the hopping system electrical pulses are generated naturally by the sample, which essentially acts like a current-controlled pulse-
generator. This inherent property of the hopping regime should make these systems attractive for the study of cavitation phenomena.

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