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Transport anomalies in highly doped conjugated polymers at low temperatures

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The anomalous low-temperature transport properties of highly doped conjugated polymers are considered. It is concluded that the behavior observed is in particular related to glassy properties of the polymer structure. The two-level system (TLS) model and more general soft potential model are applied to explain behavior of the conductivity at temperatures from several mK up to 50–70 K. The negative temperature resistance coefficient observed at low temperatures is attributed to weak localization which is at higher temperatures suppressed by a strong inelastic scattering of electrons by low-energy vibrational excitations typical for glasses. The semiconducting behavior of the resistance at high temperatures is ascribed to macroscopic disorder. Namely, we assume that the sample contains regions which are insulating at $T\rightarrow 0$ but contribute to the conductivity at elevated temperatures. [S0163-1829(99)03717-0]

I. INTRODUCTION

The nature of the metallic state in conducting polymers continues to be a subject of intensive research but, at the same time, is still a topic of significant controversy.[1,2] On the one hand, in some papers it is argued that the temperature dependence of the conductivity, thermopower, magnetic susceptibility, and magnetoconductance of heavily doped conjugated polymers demonstrates properties that are characteristic for disordered metals; the metal-insulator transition (MIT) would be described by conventional three-dimensional (3D) localization-interaction model for transport in disordered metals near the MIT.[1] On the other hand, according to another point of view, the transport properties are dominated by more macroscopic inhomogeneities and the MIT would be better described in terms of percolation between metallic islands.[2]

Recent progress in processing of stable conducting polymers has significantly improved the material quality with corresponding increase of the electric conductivity. A crossover from negative to positive temperature coefficient of resistance (TCR), a characteristic feature of metallic state, has been found recently at temperatures below 20 K for several heavily doped conjugated polymers, for example, for polyacetylene (CH)*x*,[3] polypyrrole (PPy),[4,5] and poly(phenylenevinylene) PPV.[6,7] The same crossover has been reported in our previous papers for variously doped (PF_6,CF_3SO_3, and BF_4) poly(3,4-ethylenedioxythiophene), PEDOT,[8] and ion-implanted polyaniline (PANI).[9]

It was concluded that the observed crossover from negative to positive TCR at low temperatures is not a feature specific for these systems but a phenomenon common for all highly conducting polymers.[8] It was considered within the framework of conventional 3D localization-interaction model for transport in disordered metals near MIT (Refs. 1 and 10) as a result of dominant contribution of electron-electron interaction.

However, for most of the metallic polymer samples the transport behavior at temperature below 1 K was found to be more complicated. It was shown that for such metallic polymers as $(CH)_x$, PPy, PEDOT, and ion-implanted PANI, the transition in $\rho(T)$ from positive to negative TCR with $\rho(T)\propto\ln(T)$ occurs at T below 1 K and the magnetoresistance $\Delta\rho(H,T)/\rho(0,T)$ started to depend significantly on the orientation of magnetic field, i.e., whether it was parallel or perpendicular to the film surface.[9–15] The conductivity decreased weakly with increasing parallel magnetic field while for perpendicular magnetic field the conductivity increased significantly with the magnetic-field increase.[15] The magnitude of negative magnetoresistance is larger than one expects for a simple interplay of weak localization and electron-electron interaction contributions to the 3D conductivity at millikelvin temperatures. The temperature dependence of the sheet resistivity for highly doped ion implanted PANI,[9] typical for other metallic polymers, is shown in Fig. 1. This transport behavior was interpreted as a manifestation of quasi-two-dimensional character of electron transport in highly doped PEDOT and ion-implanted PANI below 1 K because of possible existence of anisotropic plane structure.[9,15] Note that the manifestation of the quasi-2D transport in metallic polymer films at temperatures below 1 K was also reported in Ref. 18. As it is obvious, the investigated (CH)*x* , PPy, PANI, PEDOT films can not formally be considered as 2D systems, and an exact nature of possible quasi-two-dimensional character of electronic transport below 1 K in heavily doped metallic polymers is not so far completely clear.

Thus one can discriminate between three well-defined regions of the $\rho(T)$ curve: at temperatures up to ~1 K one has a negative TCR $\rho(T)\propto\ln T$, at temperatures from ~1 K up to several tens of K TCR is positive, at still higher temperatures one again deals with negative TCR with apparent power dependence of $\rho(T)$.

We would like to note that the dramatic change of the TCR sign at $T\sim 1$ K accompanied by a change of the magnetoresistance behavior is not explained in the standard localization-interaction theory because the latter by no means predicts a crossover from a *negative* to *positive* TCR with increase of temperature. Thus one expects that some additional factor enters the stage. We would like to note that the behavior at the two low-temperature regions could be explained if one assumes that at temperatures up to ~1 K the resistance is dominated by the quantum corrections (weak localization) while at higher temperatures one deals with a

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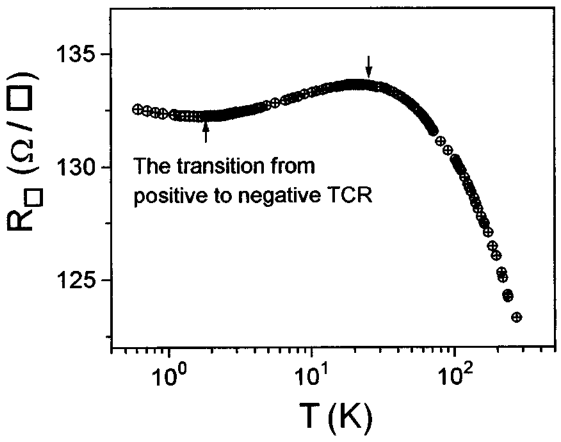

FIG. 1. The temperature dependence of the sheet resistance for polyaniline film implanted with Ar^+ ions,[9] typical for other highly conducting polymers (see, e.g., Ref. 13).

standard classical metal conductivity. Indeed, the weak localization would explain negative TCR and negative magnetoresistance (the details related to the logarithmic temperature dependence and the magnetoresistance anisotropy could be related to the layered structure discussed above). Then, the positive TCR and positive magnetoresistance are typical for classical transport in metals. The more so, the magnetoresistance isotropy at the ''classical'' region could be explained by the fact that the characteristic spatial scale for the classical magnetoresistance is the mean free path which is much smaller than the spatial scale typical for quantum interference effects.

However, such explanation inevitably meets a problem related to apparently small crossover temperature ($\sim 1-2$ K) because for standard dirty metals the crossover from the quantum correction regime to the classical one related to a suppression of the phase coherence by inelastic processes is expected for much higher temperatures.

In the present paper we analyze anomalous transport properties of highly conducting metallic polymers taking into account that the glassy properties of the polymer structure can be of importance at very low temperatures. The well-known two-level-system model as well as a more general soft potential model are applied to explain the behavior of conductivity below 1 K and in the region of positive TCR. We will show that the glassy character of the vibrational density of states (characterized by a pronounced shift to low-frequency region) can explain both the suppression of weak localization at small temperatures and the significant contribution of inelastic scattering to resistance at higher temperatures.

Then, we suggest that the semiconductor behavior (negative TCR) for temperature dependence of conductivity observed at high temperatures is related to ''macroscopic'' disorder leading to spatial inhomogeneities of the mobility edge position and, correspondingly, to a percolation conductivity. In this case a raise of temperature increases a mass of the percolation cluster thus increasing the conductivity.

II. SPECIFIC FEATURES OF LOW-TEMPERATURE TRANSPORT IN METALLIC POLYMERS: MODEL AND DISCUSSION

The typical explanation of low-temperature resistance behavior in polymers is based on the well-known interaction-localization model combining effects of electron-electron interactions and weak localization which gives for 3D case the following expression for the correction to the zero-temperature conductivity:

$$\sigma=\sigma_0+mT^{1/2}+BT^{p/2}, \qquad (1)$$

where the second and third terms on rhs correspond to interaction and localization contributions, respectively (see, e.g., Ref. 16); the exponent p is determined by the phase-breaking mechanism. It is important that $B>0$ while typically $p\geqslant 1$. Thus this formula can in principle describe a crossover from positive to negative TCR (see Ref. 10) for samples deep in metallic regime (where $m<0$). At the same time, it can hardly explain the crossover from *negative* to *positive* TCR discussed above. This fact throws a doubt in the relation between the positive TCR observed at intermediate temperatures and the interaction contribution suggested, in particular, in Ref. 10. We would like also to note that the negative TCR at higher temperatures persists up to rather high temperatures and for some samples corresponds to a significant decrease of resistance (up to several times); thus it hardly can be related to the weak-localization contribution. However, the negative TCR extending into the high-temperature region is rather typical for metallic glasses (see, e.g., Refs. 23 and 24).

The facts mentioned above stimulate us to look for another explanation for the low-temperature resistance behavior in the conducting polymers exploiting the fact that these materials exhibit not only an electronic disorder, but also a rather strong structural disorder. Indeed, as is generally believed, the structure of polymer materials exhibits—at least at low temperatures—the properties typical for glassy materials. Moreover, the polymer materials are often used as model materials to study the glassy structures[17]. The glassy properties can be even intensified in highly doped metallic polymers, where the dopant leads to increase of disorder, the weak bonding between the conjugated polymer chains supporting ''soft'' atomic configurations.

A well-known feature of glasses is the fact that the low-energy excitations in these materials differ significantly from those typical for crystals. In particular, the low-temperature heat capacity is dominated by the two-level systems (TLS) related to atoms or atomic groups which move in double-well interatomic potentials with ''soft'' barriers[19,20] (Fig. 2). While for ''strong'' barriers the states corresponding to each of the wells are decoupled, the tunneling through the barrier leads to collectivization of the lowest states. These excitations are characterized by a constant density of states

$$P(E,r)=P_0r^{-1}(1-r)^{-1/2}, \qquad (2)$$

where $r=(\Delta_0/E)$, $E=(\Delta^2+\Delta_0^2)^{1/2}$ is the TLS interlevel splitting, Δ is the atomic potential asymmetry (that is the potential difference for the two of the TLS configurations), while Δ_0 is the tunneling matrix element.

However, the TLS model was shown (see, e.g., Ref. 20) to hold only for low energies E up to some characteristic energy which for typical glasses is about 1–10 K. The limitation for the TLS model is related to the fact that for the soft double-well potentials presence of upper levels (in addition to the two lowest which form the TLS) becomes important with increase of the energy. Then, at high enough tempera-

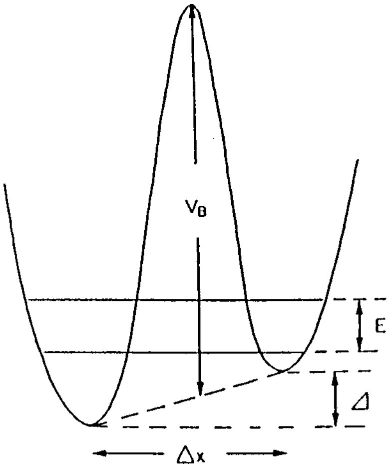

FIG. 2. Double-well potential giving rise to the two-level system at low temperatures. The two lowest levels with spacing E are shown.

tures the activation over the barrier starts to dominate the tunneling. To consider the excitations with higher energy the so-called soft potential model was suggested in Ref. 21.

The model assumes soft local modes with anharmonic potentials

$$V(x)=\mathcal{E}[\eta(x/a)^2+\xi(x/a)^3+(x/a)^4] \tag{3}$$

to exist in glasses; $\mathcal{E}$ is an energy of atomic scale, x is a displacement in terms of some generalized coordinate, a is a distance of the order of the interatomic spacing, while η and ξ are random parameters; η characterizes the effective elastic modulus while ξ characterizes the potential asymmetry. Here we will not go into details concerning this model referring, to e.g., the review paper[20].

As is seen from Eq. (3), the double-well potentials correspond to negative values of η while the ‘‘extremely soft’’ potentials correspond to $\eta \rightarrow 0$. As was shown (see, e.g., Ref. 20), the natural energy scale for the low-energy excitation in glasses is given by a characteristic energy $W \sim 3-10$ K of the order of the interlevel spacing in the quartic potential, Eq. (3), with $\eta=0$, $\xi=0$:

$$W=\mathcal{E}\eta_L^2,\ \eta_L=(\hbar^2/2Ma^2\mathcal{E})^{1/3}.$$

Here M is the effective mass of the mode. The analysis of the density of states for the excitations characterized by the potential (3) (Ref. 20) shows that for the interlevel spacings E smaller than W the density of states $P(E)$ is dominated by the two-level systems while for higher energies one mostly deals with some different sort of localized excitations: at $E<3W$ they are strongly anharmonic single-well and weak double-well effective potentials and for $E>3W$ they are quasilocal harmonic vibrations. As for the latter region, one expects $P \propto E^4$ (Ref. 20) while in the intermediate region the $P(E)$ dependence is also expected to be strong enough.

The excitations corresponding to potentials of the sort in Eq. (3) are obviously localized ones. In addition, the vibrational excitations in glasses also include delocalized excitations—phonons; however, until now, it has only been proved that the delocalized excitations exist in the low-energy region of the spectrum where their coupling with the localized vibrations is rather weak. At high frequencies the strong coupling between excitations of different sorts makes the problem difficult enough and it is still unclear whether the delocalized excitations exist at the high-energy part of the spectrum. However, it is known that the spectrum is in any case dominated by the localized excitations discussed above, the total density of states being shifted to lower energies as compared to the one typical for crystalline materials.

Having these considerations in mind one expects that the inelastic scattering of electrons is dominated by the localized excitations. In particular, for low temperatures the effective inelastic relaxation time is controlled by the TLS’s and can be estimated as

$$\frac{1}{\tau(E)} \simeq \frac{\hbar^2}{mp_F}\frac{|V(E)|^2}{\varepsilon_F^2}P(E)E. \tag{4}$$

Here $V(E)$ is the matrix element for the inelastic electron scattering by the excitation. For the two-level-system (TLS) model one can specify $V(E)$ as $V=\tilde{V}(\Delta_0/E)$ where $P(E)$ is given by Eq. (2), while $\tilde{V}$ is the difference of the potentials felt by an electron for the two TLS configurations. Actually, in this case one should integrate over r for a given E, but it only gives a factor of the order of unity. Note that for the polymer system in question with relatively small values of ε_F one may expect $\tilde{V} \sim \varepsilon_F$. Taking for $\overline{P}$ the value 10^{33} erg^{-1} cm^{-3} (which is somewhat larger than for typical glasses but may be valid for a doped polymer) and for $\hbar/p_F$ (characterizing the spatial scale of the electron wave function) the value $\sim 10^{-7}$ cm one obtains for $E \sim 1$ K $\tau \sim 10^{-10}$ s assuming the electron mass m to be equal to its ‘‘bare’’ value 10^{-27} g. This estimate is in agreement with experiment. Note that actually one can expect a larger value of m for the polymers thus even increasing the estimated value of τ which allows us to have lower values of P to ensure the experimental values of τ. As for the temperature dependence of the inelastic relaxation times, one easily sees that for the TLS model $\tau^{-1} \propto T$.

However—as was discussed above—the TLS’s are effective only up to energies $\sim W$. At higher energies one mostly deals with some different sort of localized excitations, i.e., at $W<E<3W$ with strongly anharmonic single-well and weak double-well effective potentials and for $E>3W$ with quasilocal harmonic vibrations. In contrast to the TLS, for these excitations $P(E)$ dependence is expected to be strong. Then, one may also expect very large probabilities of the inelastic processes for strongly anharmonic ‘‘three-level systems’’ related to double-well potentials with a barrier height $\sim W$ (Ref. 22) where the third level is above the barrier because of the most pronounced difference between the different states in such systems. So one concludes that with increase of the temperature above W the probability of inelastic scattering of electrons by excitations can strongly increase.

The inelastic electron relaxation affects electron transport in two ways. First, it contributes to the phase relaxation rate and thus to the weak-localization corrections. Here one expects the scattering by the localized excitations with relaxation rates as high as $\sim 10^{-10}$ s^{-1} at energies ~ 1 K to dominate over the electron-electron scattering contribution.

One notes that due to the fact that the relation $T\tau_{in} \gg \hbar$ holds for the mechanism in question, the value of the dephasing relaxation time coincides with τ_{in}. Correspondingly, the weak-localization corrections can be estimated as

$$\Delta R_{wl} \sim -R\frac{\hbar^2}{p_F^2 a l}\ln\frac{\tau_{in}}{\tau}(2\mathrm{D}); \quad \Delta R_{wl} \sim -R\frac{\hbar^2}{p_F^2 l^2}\sqrt{\tau/\tau_{in}}(3\mathrm{D}). \tag{5}$$

Here τ is the elastic relaxation time, l is the free path while a is the thickness of quasi-2D layer.

Then, there is a direct contribution of inelastic processes to the ''classical'' resistance. Assuming the Mattisen rule to hold, one writes for the contribution in question

$$\Delta R_{cl} \sim R\frac{\tau}{\tau_{in}}. \tag{6}$$

Note that in contrast to the low-temperature phonon scattering the inelastic channel driven by the localized excitation provides as large momentum relaxation as the elastic one thus significantly increasing the inelastic contribution to the resistance.

One expects that the classical contribution dominates over the weak-localization one if $d\Delta R_{cl}/dT > d\Delta R_{wl}/dT$. This unequality holds if

$$\tau_{in} < \tau\frac{\hbar^2}{p_F^2 a l}(2\mathrm{D}); \quad \tau_{in} < 2\tau\frac{\hbar^4}{(p_F l)^4}(3\mathrm{D}). \tag{7}$$

Thus one sees that the crossover in the temperature resistance coefficient for the pure conducting polymers can occur at small temperatures still obeying the relation $\tau_{in} \gg \tau$. Having in mind these facts one can attribute the behavior observed for $T < 1$ K to the weak-localization effects (which is also evidenced from the magnetoresistance studies).

The positive temperature resistance coefficient at higher temperatures can be explained by a domination of the classical contribution. As we have seen, for the pure enough materials this can occur even for relatively large τ_{in}. However, we would like to emphasize that at $T > W$ one expects a strong decrease of τ_{in}. Thus one can relate the crossover in question partly to this decrease. The estimate of W of the order of 1 K does not contradict strongly to the corresponding values for typical glasses; some smaller value can be explained by a more pronounced fractal structure, larger free volume and, correspondingly, more degrees of freedom for the dopant atoms.

Unfortunately, the presence of an electron system in the conducting polymers does not allow independent measurements of the parameters of the low-energy excitations like density of states because the corresponding contributions to both heat capacity and heat conductivity (which are the base of such measurements for insulating polymers) are masked by the electron contributions. In contrast to many metallic glasses which exhibit superconductivity, the electron contributions in the polymers cannot be freezed out. Thus we cannot go beyond the qualitative arguments given above.

Comparing the situation for the conducting polymers discussed above with the one known for metallic glasses (where the low-temperature features mentioned earlier, to our knowledge, were not reported) we would like to note much larger values of $p_F l$ for the typical (although dirty) metals. This factor can significantly diminish the weak-localization contribution and even prevent its observation (we mean the energy dependence of transport in metallic glasses, which will be discussed below).

As for the increase of the conductance at temperatures above 20–50 K, we would like to mention that it agrees qualitatively with the corresponding resistance behavior known for some metallic glasses (see, e.g., Refs. 23 and 24), where the negative temperature resistance coefficient is observed. This behavior is related in Ref. 24 to energy dependence of the partial conductivity defined as

$$\sigma = \int \sigma_E\left(-\frac{\partial F_0}{\partial E}\right)dE \tag{8}$$

(F_0 being the Fermi distribution). Assuming that $\sigma_E(E-E_F) = \sigma_E(E_F) + C(E-E_F)^\alpha$ (E_F being the Fermi energy) and modeling $\partial F_0/\partial E$ as $\exp - |E-E_F|/T$, one obtains

$$\Delta\sigma(T) = CT^\alpha\Gamma(1+\alpha) \tag{9}$$

(Γ being the gamma function).

Note, however, that the magnitude of the resistance variation known for metallic glasses is not large—$<10\%$ (see e.g., Ref. 24). At the same time, for some conducting polymers the decrease of R observed up to room temperature is much larger and can be approximately fitted by the law $R \propto T^{-0.3}$. One possible explanation is related to the fact that for the conducting polymers the mobility edge is close to E_F and thus an increase of σ_E with increase of E is more pronounced.

We would like, however, to discuss also another possibility related to the fact that the conductivity in highly doped polymers is expected to exhibit a significant spatial inhomogeneity; the evidence of this fact was recently reported in Refs. 15 and 18. As a simple model, we will assume that the effects of disorder make the mobility edge spatially dependent and thus the increase of the conductivity can be ascribed to additional conduction paths activated at the elevated temperature. Indeed, the anisotropy of magnetoresistance within the weak-localization temperature region below 1 K evidences the inhomogeneous character of the current distribution. Thus one expects that the ''metallic'' conductivity at small temperatures is related only to some cluster within the bulk sample. For the rest of the sample one expects the chemical potential μ to be lower than the mobility edge $E_m(r)$ and thus this part is in insulating state. Due to a strong disorder one expects a wide spatial scatter of the gap between the chemical potential and the mobility edge starting from its zero value. Thus the regions where $E_m(r)$ appears to be lower than the chemical potential μ support the metallic conductivity; the metallic character of the conductivity at small T means a percolation through these regions. Taking the simplest assumption that the conductivity within the metallic region does not depend on $\mathbf{r}$ one expects a critical behavior of the conductivity as a function of the relative weight of the conducting region x as

$$\sigma \propto (x-x_c)^{\alpha},$$

where x_c corresponds to the percolation threshold while the index α is equal to 2 for the 3D case and 1 for the 2D case. One has in mind that for an insulating region the conductivity is $\propto \exp -(E_m-\mu)/T$; thus one expects that the regions where $E_m(\mathbf{r})-\mu \leqslant T$ exhibit nearly metallic conductivity and should be included into the metallic cluster with an increase of T. Thus one expects the value of x to increase with T. If one introduces the distribution function $F(E_m)$ [defined in a way that $F(E_m)dE_m$ gives the relative volume covered by regions where the value of E_m is within the interval dE_m] one obtains for the x enhancement with temperature increase as

$$\Delta x = \int_{\mu}^{\mu+T} dE_m F(E_m) \sim TF(\mu+T).$$

Thus the ''semiconducting'' behavior of the conductivity starts to be pronounced when $\Delta x \sim (x_m - x_c)$ where x_m is the mass of the metallic cluster at $T=0$. Correspondingly, when $\Delta x > (x_m - x_c)$ one has $\sigma \propto [(TF(\mu+T))^{\alpha}]$. For the experimental behavior $\rho \propto T^{-0.3}$ $(\sigma \propto T^{0.3})$ one obtains $F \propto (E-\mu)^{-1+0.3/\alpha}$.

Comparison of the model discussed above with experimental data leads to the conclusion that the controversy between the two approaches currently used to describe the charge-carrier transport in conducting polymers (Anderson transition and percolation theory) can be, to some extent, artificial. Namely, we believe that the disorder is not related to some given spatial scale (microscopic or macroscopic), but has no given scale, more exactly, it exists for any given spatial scale (exponential scatter of correlation scales). In particular, the ''microscopic'' disorder leading to Anderson transition can be accompanied by macroscopic inhomogeneity of the parameters of the transition. This factor can in particular lead to quasi-2D transport giving its signature to the low-temperature resistance behavior. Then, in this case, one can expect the mobility edge to be spatially dependent whereas the metallic cluster is formed by the regions where it is lower than the chemical potential.

We believe that this factor can explain the increase of the conductivity experimentally observed at relatively high temperatures. Indeed, the activation of carriers to the mobility edge in nonmetallic regions can increase the mass of the percolation cluster. An analysis of temperature dependencies of the conductivity of metallic polymers shows that in the temperature range 80–300 K experimental data follow with sufficient accuracy the power low $\rho(T) \propto T^{-x}$, where $x = 0.2$–0.3. Here we would like to mention that the power law of the resistance decrease with temperature increase was also predicted by Larkin and Khmelnitskii[25] for the metal-insulator transition. However, this model would lead to a monotonous increase of ρ with a temperature decrease which is in contradiction with the experimental data.

Then, we have mentioned the discrepancies existing between theoretical picture existing for a dirty metal and the behavior experimentally observed for polymers at low temperatures. In particular, the weak-localization picture (characterized by negative resistance temperature coefficient and negative magnetoresistance) is restricted by surprisingly low temperatures (<1 K) while at higher temperatures (up to tens of K) the behavior is more or less standard for classical metals. At the same time for very dirty conductors (which the polymers probably are) the suppression of weak localization by inelastic processes is typically expected for much higher temperatures. We believe that the observed behavior can be explained if one takes into account the noncrystalline structure of polymers (i.e., the microscopic disorder). The latter leads to the presence of soft vibrational modes with a large density of states which leads to a shift of the vibrational density of states to lower frequency region as compared with crystalline structures. One of the properties of these modes is their strong anharmonicity which is by many orders of magnitude stronger than the anharmonicity typical for the low-frequency phonons in crystals. Then, for these localized excitations the inelastic channel of the electron scattering—in contrast to low-frequency phonons—ensures the momentum relaxation as efficient as the elastic one. We believe that taking into account the electron coupling with the modes in question one can explain the observed behavior. Namely, one expects both the effective momentum relaxation and much larger inelastic scattering rates than those known for crystals to cause a crossover from the weak-localization limit to a standard inelastic contribution to resistance at temperatures as small as 1–3 K.

Generally, the comparison of different predictions of the soft potential model with experimental data for glasses and, in our case, for metallic polymers leads to the conclusion that this model successfully describes the behavior observed up to the temperatures ~ 50–70 K that is up to temperatures at which the ''macroscopic'' disorder ensuring the ''semiconducting'' behavior of the sample resistance starts to dominate. It could be considered as a result of superposition of transport processes with various activation energies in the same polymer sample.

III. CONCLUSIONS

The anomalous temperature behavior observed for highly conducting polymers is considered. We believe that the behavior is explained by a combination of macroscopic and microscopic structural disorder. In particular we ascribe the semiconducting behavior observed at high enough temperatures to a contribution of the sample regions which continue to be insulating at $T \rightarrow 0$. The nonmonotonous behavior with a minimum at $T \sim 1$–3 K is considered to be resulting from a combination of different aspects of the microscopic disorder. The negative resistance coefficient at low T is attributed to the weak localization which is at higher temperatures suppressed by a strong inelastic scattering resulting from the low-energy vibrational excitations typical for glassy structures. The soft potential model is shown to be in qualitative agreement with this behavior.

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