# 1/f noise due to motion of dislocations and intergranular boundaries

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It is shown that thermally activated motion of extended defects such as dislocations and intergrain boundaries is a microscopic mechanism that ensures an exponentially broad relaxation spectrum for point defects, and leads thus to the onset of a $1/f$ noise in the sample resistance. Estimates that agree with experiment are obtained for the amplitude of $1/f$ noise induced by dislocation motion in massive samples and in thin metallic films. The amplitude of the current $1/f$ noise is found to be independent of temperature in a broad temperature interval. The mechanism of generation of magnetic $1/f$ noise by motion of domain walls in magnets is discussed briefly.

## 1. INTRODUCTION

Current noise, whose spectral density increases with decrease of frequency approximately as $f^{-1}$ (called $f$-noise), is observed in an unusually large class of objects and can apparently be regarded as a universal property of conductors.[1,2] Moreover, recent experimental investigations have shown that a similar spectrum of fluctuations in response to an external excitation are observed in practically any real disordered system. Examples are magnetic noise in ferromagnets[3] and in spin glasses,[4] current fluctuations in systems with charge-density waves,[5] and others.

It is customary now to assume that there is no single microscopic mechanism leading to the onset of $1/f$ noise. A suggested common cause of this effect in diverse systems is the existence of an exponentially broad distribution of the relaxation times of the fluctuating parameter. In fact,[1,2] let $x(t)$ be a fluctuating quantity with zero mean value. In the presence of only one characteristic relaxation time $\tau$, the correlation function is

$$\langle x(t)x(0)\rangle = Ae^{-t/\tau},$$

and the corresponding spectral density $S_x(\omega)$ ($\omega = 2\pi f$) is

$$S_x(\omega)=4\int_0^\infty dt\langle x(t)x(0)\rangle\cos\omega t = A\frac{4\tau}{1+(\omega\tau)^2}. \quad (1)$$

Let the system have a set of relaxation times, each determined by some activation process, so that

$$\tau_E=\tau_0e^{E/T}. \quad (2)$$

The spectral density $S_x(\omega)$ is now defined as

$$S_x(\omega)=4A\int dE\,P(E)\frac{\tau_E}{1+(\omega\tau_E)^2}, \quad (3)$$

where $P(E)$ is the distribution function of the energy barriers. If the characteristic scale of variation of the distribution function is $E_0\gg T$ and $P(E)$ is a sufficiently smooth function, we obtain by substituting (2) in (3)

$$S_x(\omega)\sim\frac{4AT}{\omega}P\left(\ln\frac{1}{\omega\tau_0}\right). \quad (4)$$

Thus, a broad distribution of the relaxation time leads indeed to $1/f$ noise (with logarithmic corrections). A great variety of microscopic mechanisms ensure such a distribution. In the existing models, an exponentially broad set of times is usually related to relaxation of microscopic local defects. The substantial difficulty encountered in such models is how to obtain a uniform distribution of the activation energies. In fact, if the local defects consists of a small number (two, three, etc.) of elements, only a limited choice of activation energies is possible, while to observe a continuous $1/f$ spectrum it is necessary [see Eqs. (2) and (3)] that the difference between the activation energies be much smaller than $T$ ($\approx 0.03$ eV at room temperature). As the temperature is lowered, one should observe in such systems a decay of the noise spectrum into individual Lorentz peaks of type (1); the fact that no such phenomenon exists imposes even more stringent requirements on models with local defects.

Another difficulty of such a model is that it predicts a linear temperature dependence [see Eq. (4)] of the $1/f$ noise, which appears upon averaging over a broad distribution of $P(E)$. The additional temperature dependence of the dimensionless quantity $S_R(\omega)/R^2$, where $R$ is the sample resistance, can contain only even powers of $R$ (for metals, $R\propto T$ at high temperatures and $R\approx$const at low ones). In most experiments on metallic films under conditions when $R\propto T$, however, $S_R(\omega)/R^2$ has turned out to be independent of temperature.

We shall show in this paper that a continuous activation-energy spectrum is produced in natural fashion in models that include extended defects. These can be dislocations, boundaries between grains or phases, domain walls in magnets, etc. Extended defects have a large number of interacting degrees of freedom. In addition, they interact with randomly disposed local defects and with one another, and this leads to a practically continuous spectrum of the relaxation times in finite-size samples.

We restrict ourselves mainly to $1/f$ noise due to dislocation motion in metals and in thin metallic films. This restriction is due mainly to the circumstance that in $1/f$-noise experiments this motion is the subject of most investigations, so that we can compare our numerical estimates with available data. It is precisely in thin films, as we shall see below, that it is possible to have a regime in which the noise amplitude is independent of temperature in a sufficiently wide range.

In the next section we shall discuss the existing models that relate the microscope motion of defects with sample-resistance fluctuations. We shall show next that extended defects (dislocations) have an exponentially wide spread of relaxation times and serve as switches of sort for local defects and for the related microscopic interference pattern of electron scattering. The amplitude and temperature dependence

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of $1/f$ noise generated in this manner agree with available experimental data. At the end of the article we shall discuss briefly noise connected with domain-wall motion in magnets.

## 2. DEFECT MOTION AND RESISTANCE FLUCTUATIONS IN METALS

A direct connection was established in recent experiments[6–12] between the defect content of a metallic film and the $1/f$-noise amplitude. It is indicated in Refs. 13 and 14 (see also Kogan's review[1]) that the most probable noise-generation mechanism is reorientation of defects having a symmetry lower than that of the crystal lattice. An example of such a defect is a dislocation or a split interstitial defect. Random reorientation of the defects changes the anisotropy of the local microscopic resistance and leads in final analysis to fluctuations of the sample resistance.

Defects can have also a more complicated structure and can consist of three and more elements (impurities). Note that the change of the electric resistance in this model, called the local-interference model, depends only on the change of the mutual orientation or of the distance between isotropic scattering centers that make up the complex defect. Simple displacement of a single isotropic defect in an ordered lattice makes no contribution to the sample resistance.

An alternate model of resistance fluctuations was proposed by Feng, Lee, and Stone.[15] This model can be used in the limit of low temperatures, when the electron phase coherence length due too electron–phonon processes is much longer than the mean free path for interaction with impurities. In this case an important role is played by multiple interference of electrons as they are scattered by impurities. As shown in Ref. 15, and also in an independent paper by Al'tshuler and Spivak,[16] displacement of even one defect under multiple-interference conditions alters the sample resistance substantially. Thus, according to Ref. 15, thermally activated motion of impurity atoms (with a suitably chosen distribution of the relaxation times) can explain the $1/f$ noise phenomenon. With rise of temperature, when the electron phase-coherence length is decreased by thermal processes, the amplitude of the resistance fluctuations also decreases.

The relative efficiencies of the local- and multiple-interference mechanisms are compared in Ref. 14. It is shown that when the impurity mean free path $l$ is varied, the amplitude ratio $\langle(\delta\rho)^2\rangle_{\rm li}/\langle(\delta\rho)^2\rangle_{\rm mi}$ of the noise produced by these mechanisms of local and multiple interference (for displacement of one defect) changes greatly (like $l^{7/2}$) and reaches a value on the order of unity for $l = l^* \approx 25$ Å. Thus, in the opinion of the authors of Ref. 14, the multiple-interference mechanism is effective only in strongly disordered metals with $l < l^*$. Note that it is assumed in this estimate that the densities of the mobile defects that cause the resistance change are the same in the two models. In fact, however, only complex local defects contribute to the local interference, whereas a part in the multiple-interference mechanisms is played by all mobile defects, whose number can be larger by several orders. Therefore, within the framework of the model of thermally activated impurities, the question of predominance of one mechanism or the other when the mean free path is determined by scattering from impurities remains open. If the mean free path is determined by phonons, the local-interference mechanism predominates.

## 3. MOTION OF A DISLOCATION IN A MEDIUM WITH DEFECTS

As noted in the Introduction, the weak spot of both models is the question of the exponentially broad distribution of the times connected with the rearrangement and displacement of the defects. We propose here for the onset of a wide distribution of the relaxation times a microscopic mechanism connected with the presence of extended defects (dislocations and intergranular boundaries) in a metallic film. In a joint paper by one of us and Ioffe[17] it is shown that the fluctuation motion of extended defects is logarithmically slow. The reason is the large spread of the activation energies corresponding to transitions of segments of a dislocation line (or of intergranular boundaries) between equilibrium (metastable) positions. To be specific, we consider here only dislocation motion. A generalization to include a two-dimensional interphase boundary will be made in Sec. 8.

Let us describe briefly the results of Ref. 17 and the model employed. Dislocation motion in a field of random defects can be regarded as a sequence of thermally activated transitions of dislocation segments between metastable states of the dislocation. These metastable states correspond to different positions of the dislocation in a random field of defects, and the most effective are transitions between states of close energy ($\Delta E \lesssim T$). The characteristic energy scale that determines the distribution of the barriers between such states and also the spread of their energies is called the pinning energy $E_p$. The pinning energy is connected with the length $L$ of the fluctuating segments and increases with it.

A circumstance of importance to us is that over large scales (and hence also over long times) the dislocation motion is determined by the defects, which form a random-force field. The random force acting on the dislocation is produced by defects that change their energy state when the dislocation passes through them. An example of defects of this type is a divacancy, whose components are loated on opposite sides of a slip plane (see Fig. 1). When an edge dislocation passes between the vacancies making up the pair, the "upper" vacancy shifts together with the "extra" atomic half-plane and the distance between vacancies, and hence the divacancy energy, changes. It can thus be assumed that a dislocation passing through a divacancy is acted upon by a force whose random character is connected with the random disposition of the divacancies in the slip plane. Note that a similar effect arises when a dislocation interacts with interstices; the latter are known to form, in metals having an fcc lattice, complex defects with dumbell-shaped configuration—split interstices.

To estimate the characteristic times of the dislocation and the parameters of the distribution function is it neces-

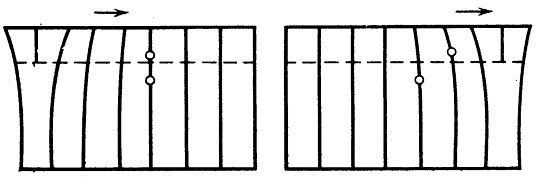

FIG. 1.

sary to estimate the pinning energy of a dislocation segment of length $L$. This estimate was obtained in Ref. 17 by using the method of Ymry and Ma,[18] in which the energy connected with the flexural fluctuations is optimized. The loss of elastic energy by flexure is of the order of $\varkappa u^2/L$, where $\varkappa$ is the elastic energy per unit dislocation length and $u$ is the transverse displacement of the fluctuating segment. The compensating gain due to the interaction with the impurities is $V(cLu)^{1/2}$, where $V$ is the energy of interaction with an individual defect and $c$ is the density of the complex defects in the slip plane. Recall that we are discussing motion in a random-force field, therefore a contribution is made by all complex defects found on the swept area $Lu$. Minimizing the total energy, we obtain for the displacement

$$u \approx L(c^{1/2}V/\varkappa)^{2/3}, \tag{5}$$

and the corresponding energy barrier is

$$E(L) \approx L\varkappa(c^{1/2}V/\varkappa)^{4/3}. \tag{6}$$

For $E(L) \gg T$ (this is precisely the situation we consider) the characteristic hop time is

$$t(L) \approx \tau_0 e^{E(L)/T}, \tag{7}$$

where $\tau_0^{-1} \approx 10^{12}$–$10^{14}$ is the characteristic trial frequency, which is of the same order as the Debye frequency.

Note that the real density of the defects in the slip plane is substantially higher than the average over the sample (we have in mind here the bulk density recalculated for the density $c = n_i a$ in an atomic plane, where $n_i$ is the bulk density of the defects and $a$ is the lattice constant). The point is that a moving dislocation is by itself a powerful source of vacancies. In addition, energetic diffusion of point defects towards the dislocation core takes place in the dislocation strain field. As a result, the moving dislocation is surrounded by an atmosphere of defects[19] (Cottrell atmosphere) and the effective density of complex defects in the slip plane can be estimated at $c = 10^{12}$–$10^{14}$ cm$^{-2}$.

## 4. RESISTANCE FLUCTUATIONS. HIGH TEMPERATURES

The restructuring of a complex defect is accompanied by a change of the cross section $\sigma$ for scattering of the conduction electrons, by a value of the order of the lattice unit-cell area. The relative change of the resistivity can be estimated at

$$\delta\rho/\rho \approx \sigma l/\Omega, \tag{8}$$

where $\Omega$ is the volume of the sample. The relative mean squared change of the conductivity for a random restructuring of many defects is

$$\langle(\delta\rho)^2\rangle/\rho^2 \approx (\sigma l)^2 n(t)/\Omega, \tag{9}$$

where $n(t)$ is the density of the defects restructured or reoriented in a time $t$.

Most experimental results on $1/f$ noise are known to be well described by Hooge's empirical equation[20]

$$S_R(\omega)/R^2 \approx \alpha/N\omega, \tag{10}$$

where $N$ is the total number of carriers in the sample and $\alpha$ is known as the Hooge constant and ranges in various cases from $10^{-1}$ to $10^{-4}$. The most widely used is the value (obtained by Hooge) $\alpha \approx 2\cdot 10^{-3}$; as a rule, $\alpha$ depends here little on temperature.

Since the noise is inversely proportional to the number $N$ of carriers, samples of minimum possible volume are usually employed to increase the measurement accuracy. As a rule these are thin films, not thicker than 1000 Å. Depending on the impurity density and on the relations between the constants contained in (6), the characteristic length $L(t)$ of the fluctuating dislocation segment can be smaller as well as larger than the film thickness $d$. In the experimental situations known to us, $L(t) > d$ (for typical values $\varkappa a \approx 1$ eV,. $T \approx 100$ K, $ca^2 \approx 0.01$, measurement frequencies $\omega \approx 10$ Hz, and $V \approx 10^{-2}$ eV we obtain the estimate $L \gtrsim 10^3$ Å), but nonetheless we consider for the sake of completeness both cases.

### A. Massive samples, $L(t) < d$

Using (4) and (5) we easily obtain

$$n(t) \approx N_D\left(\frac{\varkappa}{c^{1/2}V}\right)^{4/3}\frac{T}{\varkappa}\ln\frac{t}{\tau_0}, \tag{11}$$

where $N_D$ is the dislocation density. Substituting $n(t)$ in (9) we obtain after taking a Fourier transform

$$S(\omega) \approx N_D\left(\frac{\varkappa}{c^{1/2}V}\right)^{4/3}\frac{T}{\varkappa}\left(\frac{\sigma l}{V}\right)^2\frac{\ln(1/\omega\tau_0)}{\omega}. \tag{12}$$

### B. Thin films, $L(t) > d$

It can be concluded from dislocation-mobility data that the minimum length of a flexible dislocation segment is of the order of several hundred angstrom. Therefore for $d < 1000$ Å and for times $10^{-3}$–1 s, which exceed $\tau_0$ by nine to twelve orders, $L(t)$ is limited by the sample thickness (we consider dislocations with edge components, which broach the film in a transverse direction). In this case we can neglect the internal degrees of freedom and the elastic energy connected with the dislocation deformation, and regard the dislocation as a particle diffusing one-dimensionally in the field of a random force. A typical energy barrier surmounted in a time $t$ is

$$E \approx T\ln(t/\tau_0). \tag{13}$$

Its value [cf. the derivative of (5)] is determined by the typical fluctuation of the defect field on the swept area:

$$V(dauc)^{1/2} \approx T\ln(t/\tau_0), \tag{14}$$

and, since $dauc$ is the average number of defects on the swept area, the average density of defects reoriented in a time $t$ is

$$n(t) \approx N_D\frac{T^2}{V^2 d}\ln^2\frac{t}{\tau_0}. \tag{15}$$

Substituting (15) in (9) and using for the mean free path the estimate $l \approx v_F\tau_{\text{in}} \approx v_F\hbar/T$, which is valid at room temperature and for not too dirty samples, we get

$$\frac{S(\omega)}{R^2} \approx \frac{1}{\Omega}(\sigma v_F\hbar)^2\frac{N_D}{V^2 d}\frac{\ln^2(1/\omega\tau_0)}{\omega}. \tag{16}$$

It is remarkable that this expression contains neither a temperature dependence nor a dependence on the density of the complex defects. With the exception of $N_D$ and $d$, all the microscopic parameters in (16) can be regarded as almost constant.

The dislocation density in films is usually $\sim 10^8$ cm$^{-2}$, the electron density is $n \approx 8.5 \cdot 10^{22}$, $\sigma \approx 4\pi/k_F^2 \approx 6.5 \cdot 10^{-16}$ cm$^2$, $\varepsilon_F/V \approx 10^2$, and $d \approx 10^3$ Å. The noise is measured in the range 1–$10^4$ Hz. Assuming $\ln(1/\omega\tau_0) \approx 25$, we obtain the estimate

$$\alpha \approx 6 \cdot 10^{-3}, \tag{17}$$

which turns out to be surprisingly close to the typical observed value of Hooge's constant.

## 5. LOW TEMPERATURES

At low temperatures, when the electron coherence length $\tau_{in} v_F$ (where $\tau_{in} \approx \hbar/T$) is much larger than the impurity mean free path, multiple interference processes become significant. In this case the amplitude of the sample resistance fluctuations for a displacement of one impurity is determined by a fraction of all the random broken trajectories with spacing $l$ and total length $\tau_{in} v_F$ which are affected by the displacement of this impurity. The corresponding estimate, given by Al'tshuler and Spivak[16] and also by Feng, Lee, and Stone,[15] takes in our notation the form

$$N\frac{\langle(\delta\rho)^2\rangle}{\rho^2} \approx \left(\frac{mv_F}{\hbar}\right)^2 \left(\frac{L_{in}}{l}\right)^3 \frac{n_m}{n}\sigma. \tag{18}$$

Here $n_m$ is the density of the mobile defects (i.e., those not changing their position), and $L_{in} = (1/3 v_F \tau_{in})^{1/2}$ is the diffusion length connected with the characteristic total trajectory length $v_F \tau_{in}$. Equation (18) is written under the assumption that the number of displaced defects is small. If the density $n_m$ of the displaced defects reaches the value $1/\sigma L_{in}^2$ at which one atom is displaced on the average on each random trajectory, the noise reaches saturation.

Relation (18) was obtained for massive samples with linear dimensions exceeding the diffusion length $L_{in}$. In the case of a thin film of thickness $d < L_{in}$ the noise amplitude is determined by the relation

$$N\frac{\langle(\delta\rho)^2\rangle}{\rho^2} \approx \left(\frac{mv_F}{\hbar}\right)^2 \frac{L_{in}^2 d}{l^3} \left(\frac{L_{in}}{d}\right)^2 \sigma\frac{n_m}{n}. \tag{19}$$

Note the following unique situation: since the characteristic size of the fluctuating segment is $L(t) \propto T$ and the length $L_{in} \propto T^{-1/2}$, the samples used in experiment are usually films from the standpoint of dislocation motion ($L(t) > 10^3$ Å) and three-dimensional objects from the standpoint of multiple interference ($L_{in} \lesssim 100$ Å). At extremely low temperature the situation may be reversed. We confine ourselves, however, to the more realistic situation $L(t) > d > L_{in}$.

As the dislocation moves, the atomic planes in the samples become reconnected (see Fig. 1). This introduces distortions in all the trajectories that pass through the sites swept by the dislocation motion. If the dislocation displacement $u$ is small compared with the mean free path $l$, the number of "rearranged" sites contained in (18) is

$$n_m \approx N_D u/\sigma, \quad u < l. \tag{20}$$

If the dislocation displacement exceeds $l$, a second, third, etc. intersection of the random-walk trajectories takes place; the fraction of new trajectories that are crossed only the first time is then small. This number can be estimated by starting from the circumstance that for a displacement $\approx L_{in}$ a complete replacement of the ensemble of crossed trajectories must take place. Thus, for $u > l$ the density $n_m$ is given by

$$n_m \approx N_D \frac{l}{\sigma}\left(1 + \frac{u}{L_{in}}\right). \tag{21}$$

## 6. TEMPERATURE DEPENDENCE OF NOISE AMPLITUDE

Using Eqs. (18) and (16) we can estimate the temperature $T_c$ at which the local-interference mechanism is replaced by the multiple-interference mechanism:

$$\frac{\alpha_{li}}{\alpha_{mi}} \approx 0{,}16 ca^2 \frac{\varepsilon_F}{T}\frac{l}{a} \approx 1 \tag{22}$$

[in the derivation of this estimate we used $n_m \approx N_D(lu/\delta L_{in})$], hence

$$T_c \approx 0.16 ca^2 \varepsilon_F (l/a) \tag{23}$$

and at typical values of the parameters ($ca^2 \approx 10^{-2}$, $l/a \approx 5$) the value of $T_c$ is in the room-temperature range, $T_c \approx 10^2$. In addition, two other characteristic temperatures occur: the temperature at which the film turns into a bulky object from the standpoint of dislocation motion ($L(t) \approx d$):

$$T_d \approx d\varkappa\left(\frac{cV^2}{\varkappa^2}\right)^{1/3} \frac{1}{\ln(1/\omega\tau_0)}, \tag{24}$$

and the temperature at which the dislocation-detachment regime changes:

$$T_i \approx \frac{V}{\ln(1/\omega\tau_0)}. \tag{25}$$

If the basic assumption that the pinning is weak holds, $ca^2\varepsilon_F \gg V$, then $T_c/T_i \gg 1$. Next,

$$\frac{T_c}{T_d} \approx 0{,}16 \frac{l}{d}(ca^2)^{2/3}\left(\frac{\varkappa a}{V}\right)^{4/3} \ln\frac{1}{\omega\tau_0}, \tag{26}$$

$$\frac{T_d}{T_i} \approx \frac{d}{a}(ca^2)^{1/3}\left(\frac{V}{\varkappa a}\right)^{1/3} \tag{27}$$

and the most realistic is the relation $T_c \gg T_d \gtrsim T_i$, with $T_i \approx 5$ K.

Thus, the most probable temperature dependence will be a smooth ($\alpha_{mi} \propto T$) growth of the amplitude up to the temperature $T_c$, followed by a flattening at $T \approx T_c$, when the noise amplitude is determined by the local interference and is independent of the temperature. Note that expressions (24) and (25) contain in the right-hand sides the logarithm of the frequency. Therefore, by measuring the shifts of the singular points of the temperature dependence of the noise amplitude as the frequency is varied, it is possible in principle to separate the contributions containing the functional dependence $f(T\ln(1/\omega\tau_0))$ connected with the change of the dislocation regime, from the $\varphi(T)$ dependence that sets in when the electric-resistance mechanism is changed. This makes it possible to assess the dislocation-motion regime and the mechanism of noise generation in a resistor.

## 7. EXTREME LOW TEMPERATURES

The initial assumption of weak dislocation pinning by impurities no longer holds at temperatures on the order of

$T \approx V/\ln(1/\omega\tau_0)$, where $\omega$ is the characteristic frequency at which the measurements are made. At lower temperatures it can be assumed that during the measurement time no substantial change takes place in the picture of dislocation pinning to the impurities. All that is possible in this case is taken to be motion of dislocation segments between strong-pinning points. At high temperatures this motion could be regarded as diffusion, but at low temperatures account must be taken of the dislocation interaction with the frozen stress fields due to microscopic defects and the boundaries of grains with impurities located alongside the disolocation slip plane.

Thus, for arbitrarily low temperatures, transitions with arbitrarily low energies exist in a system with dislocations. Even small changes of the dislocation positions lead (via the multiple-interference mechanism) to substantial resistance fluctuations.

## 8. MAGNETIC NOISE

The logarithmic (in time) character of thermally activated motion of domain walls can also be a source of magnetic noise. It was recently reliably established by experiment (see Ref. 3) that the noise component of the magnetic moment of a ferromagnetic sample placed in an alternating field is connected with dynamics of the domain-wall motion in the course of remagnetization in the field of random inhomogeneities of the sample. In particular, an increase of the spectral density of magnetic noise near the remagnetization frequency was established and investigated in detail. It was observed[3] that the spectral density of magnetic noise near the remagnetization frequency and its harmonics varied approximately as $1/f^1$ (the pump frequency in this case was of the order of several megahertz and the detuning reckoned from the frequency of the harmonic varied in the range $10^{-2}$–$10^4$ Hz). This noise is attributed to slow wandering of the domain walls,[3] which was directly observed in thin films with the aid of the magneto-optic Kerr effect. The characteristic times of restructuring of such hard-to-remagnetize domains were of the order of $10^3$ remagnetization cycles.

By a method similar to the one used for the estimates in Sec. 3 (see also Ref. 17) it is possible to obtain for the average displacement of a domain wall expressions containing logarithmic dependences on the time. If the displacement is produced by an external driving force, the minimum time $\tau_0$ entering in the expressions logarithmic in $t$ is the period of this force. Owing to the nonlinearity of the system, a spectrum of the $1/f$ type can be observed not only in the low-frequency limit, but also for detuning from frequencies that are multiples of the pump frequency. In the preceding section we have considered the case of thermally activated motion, when the frequency $\omega_0$ was of the order of $10^{11}$ Hz, and therefore the quantity $\ln(\omega/\omega_0)$ in all the final expressions was practically constant in the measurement range $10^{-1}$–$10^3$ Hz. In the case of magnetic noise, when the driving frequency is higher by only two to four orders than the measurement frequency, the change of $\ln(\omega_0/\omega)$ can be substantial. This apparently explains the tendency of the observed exponent $\gamma$ to exceed unity. We present below the values of the exponent $\delta$ in the frequency dependence of the noise

$$S(\omega) \approx \omega^{-1} \ln^{\delta}(\omega_0/\omega) \tag{28}$$

for several cases[1)]: a) $\delta = 3/2$ for a bulky sample with point defects that assume the role of a random force acting on the domain wall; b) $\delta = 2$ for thin films, when the domain wall can be regarded as a string moving in a random-force field; c) $\delta = 6$ for a bulky sample and point defects acting as a random potential (e.g., microinhomogeneities interacting directly with the wall).

It must be noted that since the major inhomogeneities that control the motion of domain walls in magnets are dislocations and grain boundaries, a unique combined mechanism of $1/f$-noise generation is possible and is connected both with the logarithmically slow displacements of the domain walls and with the logarithmically slow change of the configuration of the random fields that pin the domain walls.

## 9. CONCLUSION

We have shown that the microscopic mechanism producing $1/f$ noise in metallic samples consists of local-defect displacements that are stimulated by dislocation motion. A dislocation line passing near a local defect restructures the latter (causes a small relative displacement of the defects or their reorientation relative to one another). This leads in turn to noticeable local fluctuations of the electric resistance. The exponentially broad spectrum of the local-defect relaxation times, which is needed for the onset of the $1/f$ noise, is ensured by the logarithmically slow character of the thermally activated motion of the dislocations in the random field of these defects. We have considered local interference of electrons,[13,14] a mechanism connected with the change of the cross section for electron scattering when the given defect changes orientation, and multiple interference,[15,16] a mechanism that leads to resistance fluctuations when the relative arrangement of the point defects is changed. The first mechanism is basic at room temperature, whereas the second determines the resistance fluctuations at low temperatures or in strongly disordered samples. The observed anomalies in the temperature dependence of the noise amplitude in thin films can thus be attributed to replacement of one mechanism by the other, as well as to a change of the effective dimensionality of the sample both relative to the dislocation motion and relative to the electron diffusion.

The connection between dislocations and $1/f$ noise in thin metallic film is attested to by data obtained in the series of studies by Fleetwood and Giordano.[7,8,21–23] They have observed in experiment that deformation of the sample leads to an appreciable (by an order of magnitude) increase of the $1/f$-noise amplitude. An appreciable spread in the noise amplitudes of nominally identical samples was also observed. After approximately six months the noise amplitude decreased to a certain minimum value. An analogous effect is produced also by annealing the samples. These results agree with our assumptions concerning the role of dislocations in the generation of $1/f$ noise.

Note that the mechanism of the onset of a random force acting on a dislocation, introduced in Ref. 17 and used in our paper, can determine also the frequency and temperature dependences of internal friction of crystals with dislocations.[17,24] This mechanism can explain, in particular, the high-temperature peaks in the temperature dependence of internal friction of plastically deformed metals.[24]

To verify the hypothesis set forth in our paper concerning the origin of the $1/f$ noise, measurements of the noise amplitude at various controllable dislocation densities are necessary. Important data can be obtained also from the shifts of the positions of the singular points of the temperature dependence of the noise amplitude when the measurement frequency is varied.

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[1)]We do not discuss here the quantitative theory, since the expressions for noise near different pump-frequency harmonics differ greatly in form and the values of the defect-interaction parameters contained in these expressions have not been sufficiently well investigated.

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