Properties of superconducting S-I-N, S-I-S, and S-C-S structures with amorphous weak coupling

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The properties due to the presence of two-level structure systems in superconducting tunnel junctions with amorphous insulators, as well as in point and bridge Josephson junctions with amorphous surrounds, are investigated. Equations are obtained for tunneling with participation of the two-level systems for the cases of quasistatic tunneling in an S-I-N junction (N is the normal metal) and for the case of Josephson tunneling in an S-I-S junction. It is shown that inelastic tunneling makes an additional contribution to the nonlinearity of the current-voltage characteristic of an S-I-N junction. The specific phenomena of nonexponential relaxation in this junction (in particular, tunnel-current relaxation), which have a $1/t$ dependence, are discussed. Low-frequency noise in S-I-S and S-C-S structures, due to transitions in the two-level system and having a $1/f$ dependence at not too small junction sizes are considered. In the case of the stationary Josephson effect this noise has features of critical-current fluctuations that can manifest themselves, in particular as fluctuations of the magnetic flux linked with a weakly coupled superconducting ring. Under conditions of the nonstationary Josephson effect the two-level structures lead to broadening of the Josephson-generation line. It is proposed to use the nonstationary Josephson effect to observe the echo effect in glasses.

It has by now been reliably established that one of the properties of materials with amorphous structure is the presence in them of specific low-temperature excitations. These are two-level systems (TLS)[1] that can interact with conduction electrons.[2] We have recently considered the effect of such an interaction on the properties of a tunnel junction whose insulating film is amorphous,[3] in which case the tunneling electrons can interact with the TLS of the insulator. It has been shown that a host of specific phenomena can occur in this system. These include, in particular, inelastic tunneling (which makes possible tunnel spectroscopy of TLS) as well as low-frequency $1/f$ noise. We note that this specific noise in tunnel junction was independently investigated by Kogan and Nagaev,[4a] who were the first to point in a preceding paper[4b] to TLS as possible sources of the noise.

In Ref. 3 was investigated the case of a junction with normal metals. Particular interest attaches, however, to tunneling in superconduting sytems. The tunnel effect in superconductors has long been the subject of numerous theoretical and experimental investigations, and tunnel-spectroscopy methods yield important data on the electron and phonon systems of the superconductor.[5] As for the Josephson effect (which manifests itself not only in tunnel junctions but in weak junctions of other types, particularly in S-C-S structures, where C is a constriction, it is widely used in practice (and not in pure research alone[5,6]).

Analysis of the properties of superconducting tunnel junctions and of Josephson structures is of twofold importance. First, the assumption that in actual systems this region has amorphous properties is apparently perfectly valid—this applies both to the oxide film of the junction oxide film and to the region of point-contact Josphson structures (we note incidentally that total amorphization is not mandatory for the existence of TLS, and a strong disorganization of the crystal structure suffices[7]). It is therefore important to understand the degree to which TLS can influence the properties of real junctions. This pertains primarily to low-frequency noise connected with transitions to the TLS in the case of Josephson structure, since it is just on the use of the latter that modern extremely highly sensitive measurements that are based.

Second, in our opinion, it is of appreciable interest to use both the well-developed quasiparticle tunneling procedure and the exclusive opportunities that the Josephson effect offers for the study of the amorphous state in general and of TLS in particular.

A propos both circumstances we note a recent experiment[8] in which generation of monochromatic phonons in a Josephson junction was observed under conditions of the nonstationary Josephson effect at a phonon frequency equal to the Josephson frequency. One of the possible mechanisms proposed in Ref. 8 to explain this effect involved the assumption that the barrier is vitreous and that a contribution is made by the TLS. (We, however, have proposed another explanation,[9] believed to be somewhat more realistic,[1)] in which only a disordered structure of the barrier is assumed[9].)

We consider in the present paper certain properties of superconducting structures with amorphous weak coupling. The basic analysis will be made for a tunnel junction whose dielectric film has amorphous properties, but many deductions will be made also concerning the properties of point and bridge Josephson junctions.

We obtain first an expression for the tunnel current, applicable both to quasiparticle tunneling in an S-I-N junc-

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tion (N—normal metal) and to Josephson tunneling in an S-I-S system. It turns out, in particular, that the critical Josephson current I_c depends on the TLS occupation numbers even in the absence of a normal component (at $T \ll T_c$), so that its values can be different for different configurations in a TLS system. This important circumstance is typical also of other Josephson structures with amorphous weak coupling (S-C-S and bridge junctions). In the latter case its behavior is governed by the direct dependence of I_c on the junction resistance in the normal state,[6] which depends in turn on the TLS occupation numbers.

We turn next to the case of quasiparticle tunneling. The most interesting here is the contribution of inelastic processes, in which the tunneling is acompanied by excitation (or de-excitation) of TLS. This contribution introduces an additional nonlinearity of the current-voltage characteristic of the junction (compared with that due to the singularity of the density of states of the semiconductor). Although under ordinary conditions (for large junctions, at any rate) it is apparently difficult to separate the contribution due to the TLS, the presence or the TLS can lead to specific relaxation phenomena. In particular, if a sufficiently high voltage V (such that at any rate $eV > \Delta$, where Δ is the superconducting gap) is applied to the junction at $T = 0$ and is subsequently decreased abruptly to values such that $eV < \Delta$ (at which there is no quasiparticle current under ordinary conditions), a nonexponentially decreasing ($\propto t^{-1}$) current can be observed in the system and is due to inelastic processes. Furthermore, if the external voltage is turned off, there exists an independent mechanism that ensures the presence of a slowly relaxing voltage on the junction (similar to that considered in Ref. 3).

The dependence of the critical current of the Josephson structures on the TLS ensures a specific low-frequency fluctuation mechanism which is considered in a separate section (for large junctions, these fluctuations have an $1/f$ frequency spectrum). These fluctuations of I_c can manifest themselves, in particular through fluctuations of the magnetic flux that is linked with the weakly coupled superconducting ring. In the case of the nonstationary Josephson effect, the fluctuations connected with the TLS can lead to broadening of the Josephson-generation spectrum.

We consider finally certain special phenomena that can be observed in in the systems considered. We shall discuss, in particular, the possibility of using the nonstationary Josephson effect to study the echo due to TLS.

1. EXPRESSION FOR TUNNEL CURRENT IN S-I-N AND S-I-S JUNCTIONS

As in Ref. 3, we use the tunnel-Hamiltonian method. It provides, on the one hand, a sufficiently accurate description of the tunneling processes. On the other, it presents a clear physical picture of the phenomenon. We write the total Hamiltonian of the system in the form

$$H=\sum_{\mathbf{p}} (W_{\mathbf{p}}-\mu) a_{\mathbf{p}l}^{+}a_{\mathbf{p}l}+\sum_{\mathbf{p}} (W_{\mathbf{p}}+eV-\mu) a_{\mathbf{p}r}^{+}a_{\mathbf{p}r}$$
$$+\sum_{\substack{\mathbf{p},\sigma\\ \alpha=l,r}} (\Delta_\alpha a^{+}_{\mathbf{p}\sigma,\alpha}a^{+}_{-\mathbf{p}-\sigma,\alpha}+\Delta^{*} a_{\mathbf{p}\sigma,\alpha}a_{-\mathbf{p}-\sigma,\alpha})$$
$$+\sum_{i} E_i \hat{S}_z^{\,i}+H_T^{\,0}+H_T^{\,1}+H_{is}+\tilde{H}. \tag{1}$$

Here a^{+} and a are the electronic second-quantization operators, σ is the electron spin, the subscripts l and r refer to the left and right banks of the junction, $W_{\mathbf{p}}$ is the dispersion law of the normal electrons (which was assumed for simplicity to be the same for both banks), and $\Delta_{l,r}$ is the order parameter (generally speaking, complex). The subscript i numbers the TLS, which are described by the "spin" operators $\hat{S}_z^{\,i}$, $\hat{S}_x^{\,i}$ (Refs. 1 and 2), and $E_i = (\Delta_i^2 + \Delta_{0i}^2)^{1/2}$; Δ_i and Δ_{0i} are the TLS parameters (the asymmetry of the two-well potential and the tunneling matrix element, respectively). The operator H_{iS} describes the TLS interaction with the phonons, as well as the non-tunneling interactions with the bank electrons; the latter interactions serve as an additional mechanism of the TLS relaxation.[3] The operator

$$H_T^{\,0}=\sum_{\mathbf{p}\mathbf{p}',\alpha\neq\alpha'} T_{\mathbf{p}\mathbf{p}',\alpha\alpha'}a_{\mathbf{p}\alpha}^{+}a_{\mathbf{p}\alpha'} \tag{2}$$

corresponds to the usual tunneling processes, while the operator H_T^{1} corresponds to tunneling with participation of TLS. The operator $\tilde{H}$ includes the remaining parts of the total Hamiltonian (the phonon Hamiltonian, the electron–phonon interaction, etc.).

The operator H_T^{1} can be described by the expression

$$H_T^{\,1}=\sum_{i,\mathbf{p},\mathbf{p}',\alpha\neq\alpha'} \left[K_{\alpha\alpha',\mathbf{p}\mathbf{p}'}\left(\frac{\Delta_{0i}}{E_i}\hat{S}_x^{\,i}+\frac{\Delta_i}{E_i}\hat{S}_z^{\,i}\right)+\hat{J}U_{\alpha\alpha',\mathbf{p}\mathbf{p}'}\right] a_{\mathbf{p}\alpha}^{+}a_{\mathbf{p}'\alpha'}, \tag{3}$$

where $\hat{J}$ is a unit operator in "spin" space, and the matrix elements K and U can be estimated at[3]

$$U_{\mathbf{p}\mathbf{p}',\alpha\alpha'}=U^{*}_{\mathbf{p}'\mathbf{p},\alpha'\alpha}\sim -T_{\mathbf{p}\mathbf{p}'}\frac{\varkappa d}{W_0}\left(\frac{V_{1\mathbf{q}}+V_{2\mathbf{q}}}{2}\right),\qquad K_{\mathbf{p}\mathbf{p}',\alpha\alpha'}=K^{*}_{\mathbf{p}'\mathbf{p},\alpha'\alpha}\sim -T_{\mathbf{p}\mathbf{p}'}\frac{\varkappa d}{W_0}(V_{1\mathbf{q}}-V_{2\mathbf{q}}),$$
$$\varkappa=\frac{1}{\hbar}\{2m[U_0-(W_{\mathbf{p}}-W_{\perp})]\}^{1/2}, \tag{4}$$
$$W_0=\frac{\hbar\varkappa^2}{2m},\qquad \mathbf{q}=\frac{1}{\hbar}(\mathbf{p}-\mathbf{p}'),$$

$W_\perp$ is the kinetic energy connected with motion parallel to the barrier, U_0 is the barrier height, d is the barrier thickness, and V_{1q} and V_{2q} are the Fourier components of the TLS in the barrier and are labeled in accordance with the two-hump potential well in which they are located (they are normalized to the unit-cell volume).

We define the tunnel current as

$$I=A\frac{dN_l}{dt},\qquad N_l=\sum\langle\hat{\rho}a_{\mathbf{p}l}^{+}a_{\mathbf{p}l}\rangle, \tag{4a}$$

where A is the junction area. The problem reduces thus to calculation of the time derivative of the correlator $\langle \hat{\rho} a^+ a \rangle$, where $\hat{\rho}$ is the statistical operator of the system. We apply for this purpose the Bogolyubov transformation, which eliminates the anomalous terms of the zeroth-approximation Hamiltonian:

$$a^+_{\mathbf{p},\sigma}=u_p\alpha^+_{\mathbf{p},\sigma}+v_p^*\alpha_{-\mathbf{p},-\sigma}, \qquad a^+_{\mathbf{p},-\sigma}=u_p\alpha^+_{\mathbf{p},-\sigma}-v_p^*\alpha_{-\mathbf{p},\sigma}, \tag{5}$$

where

$$u_{p\alpha}=|u_{p\alpha}|e^{i\chi_\alpha/2},\quad v_{p\alpha}=|v_{p\alpha}|e^{-i\chi_\alpha/2},\quad |u_{p\alpha}|^2={}^1/_2(1+\xi_{p\alpha}/\varepsilon_{p\alpha}),$$
$$|v_p|^2={}^1/_2(1-\xi_{p\alpha}/\varepsilon_{p\alpha}),\quad \xi_{pr}=(W_p-\mu+eV), \tag{5a}$$
$$\xi_{pl}=(W_p-\mu),\quad \varepsilon_{p\alpha}=(\xi_{p\alpha}^2+|\Delta_\alpha|^2)^{1/2},$$

and χ is the phase of the order parameter. We then obtain directly

$$\langle a_{p\sigma}^+a_{p\sigma}\rangle=|u_p|^2\langle\alpha_{p\sigma}^+\alpha_{p\sigma}\rangle+|v_p|^2(1-\langle\alpha^+_{-p-\sigma}\alpha_{-p-\sigma}\rangle)$$
$$+\sigma u_pv_p(\langle\alpha_{p\sigma}^+\alpha^+_{-p-\sigma}\rangle-\langle\alpha_{p\sigma}\alpha_{-p-\sigma}\rangle). \tag{6}$$

We confine ourselves for simplicity to an analysis of two situations:

1) $\Delta_r=0$ (*S-I-N* junction);

2) $V=0$ (pure Josephson tunneling).

As for the general case $V\neq0$, $\Delta_r\neq0$, it can be easily seen that the energy spectrum defined in (5a) does not agree in this case with the true spectrum of the system, since the position of the gap in the real spectrum of the right bank corresponds to $\xi'_{pr}\equiv W_p-\mu=0$, and not to $\xi_{pr}=W_p-\mu+eV=0$ (as would follow from (5a)). Thus, the general case of an *S-I-S* junction calls generally speaking for a more complicated analysis (in which even the zeroth-approximation Hamiltonian contains terms of the type $\alpha^+\alpha^+$ and $\alpha\alpha$), which leads to rather cumbersome calculations. No such problems arise for an *S-I-N* junction, since the normal-metal spectrum is insensitive to the location of the Fermi level and the matter reduces to modification of the distribution function. Nonetheless, the main distinguishing features of the quasiparticle tunneling in an *S-I-S* junction have their counterparts in the simpler analysis of the *S-I-N* function.

Taking the foregoing into account, transforming the Hamiltonian (1) with the aid of (5), and using the standard correlator-separation procedure to find the time derivatives of the quantities, we get after laborious but essentially straightforward transformations the contribution due to the presence of the TLS, in the form

1) *S-I-N* junction

$$I_n^1=\frac{A}{\hbar}\sum_{\mathbf{p},\mathbf{p}',i}[B_1(\mathscr{L}_1+\mathscr{L}_2)+B_2\mathscr{L}_3],$$

$$\mathscr{L}_1=|K_{\mathbf{pp}'}|^2(\pi/4)(\Delta_{0i}/E_i)^2\{F_{p'r}(1-F_{pl})[\delta(\varepsilon_{pl}-\varepsilon_{p'r}+E_i)n_-^i$$
$$+\delta(\varepsilon_{pl}-\varepsilon_{p'r}-E_i)n_+^i]-F_{pl}(1-F_{p'r})[\delta(\varepsilon_{pl}-\varepsilon_{p'r}+E_i)n_+^i$$
$$+\delta(\varepsilon_{pl}-\varepsilon_{p'r}-E_i)n_-^i]\},$$

$$\mathscr{L}_2=2[|T_{\mathbf{pp}'}+U_{\mathbf{pp}'}+{}^1/_2K_{\mathbf{pp}'}\Delta_i/E_i|^2n_+^i$$
$$+|T_{\mathbf{pp}'}+U_{\mathbf{pp}'}-{}^1/_2K_{\mathbf{pp}'}\Delta_i/E_i|^2n_-^i$$
$$-|T_{\mathbf{pp}'}|^2](F_{p'r}-F_{pl})\delta(\varepsilon_{pl}-\varepsilon_{p'r}),$$

$$\mathscr{L}_3=|K_{\mathbf{pp}'}|^2(\pi/2)(\Delta_{0i}/E_i)^2$$
$$\times\{(1-F_{-p'r})(1-F_{pl})[\delta(\varepsilon_{pl}+\varepsilon_{p'r}+E_i)n_-^i \tag{7}$$
$$+\delta(\varepsilon_{pl}+\varepsilon_{p'r}-E_i)n_+^i]$$
$$-F_{pl}F_{-p'r}[\delta(\varepsilon_{pl}+\varepsilon_{p'r}-E_i)n_-^i+\delta(\varepsilon_{pl}+\varepsilon_{p'r}+E_i)n_+^i]\}.$$

Here n_+^i and n_-^i are the occupation numbers of the upper and lower levels of the ith TLS, while F_{pl} and F_{pr} are the equilibrium distribution functions of the quasiparticles in the bank. Further, $F_{pl}=F_0(\varepsilon_l)$, where F_0 is the Fermi function. As for F_{pr}, it can be easily seen that within the framework of our representation (5) and (5a) it is given by

$$F_{pr}=F_0(|\xi_{pr}|-eV\,\mathrm{sign}\,\xi_{pr}). \tag{8}$$

For the coherence factors B_1 and B_2 we obtain

$$B_1=(|u_{pl}|^2-|v_{pl}|^2)|u_{pl}u_{p'r}-v_{pl}v_{p'r}|^2+2(u_{pl}v_{pl})^2(|u_{p'r}|^2$$
$$-|v_{p'r}|^2)+2u_{pl}v_{pl}u_{p'r}v_{p'r}(|u_{pl}|^2-|v_{pl}|^2)$$
$$={}^1/_2(\xi_{pl}/\varepsilon_{pl}+\mathrm{sign}\,\xi_{pr}),$$
$$B_2=(|u_{pl}|^2-|v_{pl}|^2)|u_{pl}v_{p'r}+u_{p'r}v_{pl}|^2-2u_{pl}^2v_{pl}^2(|u_{p'r}|^2$$
$$-|v_{p'r}|^2)-2u_{pl}v_{pl}u_{p'r}v_{p'r}(|u_{pl}|^2-|v_{pl}|^2)$$
$$={}^1/_2(\xi_{pl}/\varepsilon_{pl}-\mathrm{sign}\,\xi_{pr}).$$

2) Josephson tunneling at $V=0$

$$I_s^1=(-A/\hbar)\sin(\chi_l-\chi_r)\sum_{\mathbf{pp}'i}B_3(\mathscr{L}_4+\mathscr{L}_5+\mathscr{L}_6),$$

$$\mathscr{L}_4=\left\{\frac{1}{\varepsilon_{pl}+\varepsilon_{p'r}}[F_{-pl}-(1-F_{p'r})]+\frac{1}{\varepsilon_{pl}-\varepsilon_{p'r}}(F_{p'r}-F_{pl})\right\}$$
$$\times\left[\left|T_{\mathbf{pp}'}+U_{\mathbf{pp}'}+\frac{1}{2}K_{\mathbf{pp}'}\frac{\Delta_i}{E_i}\right|^2n_+^i\right.$$
$$\left.+\left|T_{\mathbf{pp}'}+U_{\mathbf{pp}'}-\frac{1}{2}K_{\mathbf{pp}'}\frac{\Delta_i}{E_i}\right|^2n_-^i-|T_{\mathbf{pp}'}|^2\right], \tag{9}$$

$$\mathscr{L}_5=\frac{1}{4}|K_{\mathbf{pp}'}|^2\left(\frac{\Delta_{0i}}{E_i}\right)(n_+^i-n_-^i)\left\{\frac{E_i}{(\varepsilon_{pl}+\varepsilon_{p'r})^2-E_i^2}[F_{pl}F_{-p'r}\right.$$
$$+(1-F_{pl})(1-F_{p'r})]$$
$$\left.+\frac{E_i}{(\varepsilon_{pl}-\varepsilon_{p'r})^2-E_i^2}(F_{pl}+F_{p'r}-2F_{pl}F_{p'r})\right\}.$$

Here

$$B_3=4u_{p'r}v_{p'r}u_{pl}v_{pl}=\left(\frac{\Delta_{pl}}{\varepsilon_{pl}}\frac{\Delta_{pr}}{\varepsilon_{pr}}\right).$$

The term $\mathscr{L}_6$ (which describes the effect of the TLS on the superconducting pairing) is independent of the TLS occupation numbers, and we shall not write it out explicitly, since it is of no importance for the phenomena of interest to us. We point out that the quantity

$$\frac{A}{\hbar}\sum_{\mathbf{pp}'}B_3\left[\frac{1}{\varepsilon_{pl}+\varepsilon_{p'r}}(1-F_{-pl}-F_{p'r})+\frac{1}{\varepsilon_{pl}-\varepsilon_{p'r}}(F_{pl}-F_{p'r})\right]$$
$$\times|T_{\mathbf{pp}'}|^2 \tag{10}$$
$$=I_{s0}=\frac{\pi}{2}\frac{\Delta(T)}{eR}\,\mathrm{th}\left(\frac{\Delta(T)}{T}\right)$$

(where R is the tunnel resistance) is none other than the critical current of the Josephson junction (cf. Ref. 10).

2. QUASIPARTICLE TUNNELING

We note that although the calculation presented implies $\Delta_r = 0$ (S-I-N junction), it seems to us that the main qualitative deductions and the order-of-magnitude estimates remain the same also for the general picture of quasiparticle tunneling in superconducting tunnel structures with amorphous insulators.

Let us analyze expression (7). We turn first to the contribution of the elastic tunneling (defined by $\mathscr{L}_2$). With allowance for the explicit expression (8) for F_{pr}, it is transformed into

$$I_{ne}{}^1 \sim j_0 \sum_i (\sigma_i{}^+ n_i{}^+ + \sigma_i{}^- n_i{}^-), \tag{11}$$

$$j_0 = \frac{2\pi e |T|^2}{\hbar} N_l N_r \int_\Delta^\infty d\varepsilon \nu_l(\varepsilon) [F_0(\varepsilon - eV) - F_0(\varepsilon + eV)],$$

where j_0 is none other than the density of the current through the S-I-N junction in the absence of TLS (se, e.g., Ref. 5); here N_l and N_r are the state densities of the normal electrons on the Fermi surface, and $\nu = \varepsilon/|\xi|$ is the reduced density of states of the quasiparticles. The quantities $|\sigma_i^+|$ and $|\sigma_i^-|$ have the meaning of the cross sections for elastic tunneling participation of a TLS located on the upper and lower levels, respectively. We have for these cross sections the estimate (cf. Ref. 3)

$$\sigma_i^+ \sim \frac{\pi a^2}{2} \frac{\varkappa d}{W_0} \operatorname{Re}\left\{ -T_q \left[V_{1q}\left(1 + \frac{\Delta_i}{E_i}\right) + V_{2q}\left(1 - \frac{\Delta_i}{E_i}\right)\right]\right\},$$
$$\sigma_i^- \sim \frac{\pi a^2}{2} \frac{\varkappa d}{W_0} \operatorname{Re}\left\{ -T_q \left[V_{1q}\left(1 - \frac{\Delta_i}{E_i}\right) + V_{2q}\left(1 + \frac{\Delta_i}{E_i}\right)\right]\right\}. \tag{12}$$

It can be seen from (11) that elastic interaction with TLS leads only to a renormalization of the quasiparticle current, without introducing an additional nonlinear dependence in the current-voltage characteristic of the junction This renormalization (which has the meaning of the renormalization of the tunnel resistance), however, depends on the TLS occupation numbers. The fluctuations of these numbers leads therefore to I_{ne}^1 fluctuations that are analogous to those considered in Ref. 3 for a function of normal metals.

As for the contribution of inelastic processes, it is transformed into

$$I_{ni}{}^1 \approx j_0 \sum_i \sigma_i \left(\frac{\Delta_{0i}}{E_i}\right)^2 \int_{\Delta l}^\infty d\varepsilon \nu_l(\varepsilon) \Phi(\varepsilon, V, E_i)$$
$$\times \left\{ \int_{\Delta l}^\infty d\varepsilon \nu_l(\varepsilon) [F_0(\varepsilon - eV) - F_0(\varepsilon + eV)] \right\}^{-1},$$
$$\Phi = [F_0(\varepsilon + E_i - eV) - F_0(\varepsilon + E_i + eV)]$$
$$\times [n_-{}^i(1 - F_0(\varepsilon)) + F_0(\varepsilon) n_+{}^i]$$
$$+ [F_0(\varepsilon + E_i - eV) - F_0(\varepsilon - E_i + eV)][(1 - F_0(\varepsilon)) n_+{}^i$$
$$+ F_0(\varepsilon) n_-{}^i] + [F_0(E_i - \varepsilon - eV) - F_0(E_i - \varepsilon + eV)]$$
$$\times [n_+{}^i(1 - F_0(\varepsilon)) + n_-{}^i F_0(\varepsilon)], \tag{13}$$

or, at $T = 0$ and $V > 0$, taking the possibility of a nonequilibrium population of the TLS into account,

$$\Phi = F_0(\varepsilon + E_i - eV) n_-{}^i + [F_0(\varepsilon - E_i - eV) - F_0(\varepsilon - E_i + eV)$$
$$+ F_0(E_i - \varepsilon - eV) - F_0(E_i - \varepsilon + eV)] n_+{}^i.$$

The effective cross section σ_i for inlastic interaction of the tunneling electron with a TLS is given by

$$\sigma_i \sim \pi a^2 (\varkappa d/W_0)^2 |V_{1q} - V_{2q}|^2. \tag{14}$$

It can be seen that the inelastic processes lead, generally speaking, to an additional nonlinearity of the junction current–voltage characteristic. In particular, at $T = 0$, $n_+ = 0$, and $V > 0$,

$$\frac{dI^1_{n_i}}{dV} \backsim \sum_i \nu_l(eV - E_i)\left(\frac{\Delta_{0i}}{E_i}\right)^2 \approx \int_0^{eV - \Delta_l} dE \bar{P} \nu_l(eV - E_i). \tag{15}$$

We have introduced here the TLS distribution function in the parameters E_i and $\rho \equiv \Delta_{0i}/E_i$ (Ref. 2)

$$P(E, \rho) = \bar{P}(1 - \rho)^{-1/2} \rho^{-1}. \tag{16}$$

In the absence of TLS, on the other hand, $dI_0/dV \propto \nu(eV)$. Since, however, the TLS density of states P is practically independent of the energy E, it follows that

$$\frac{dI_{ni}{}^1}{dV} \sim \bar{P} \int_{\Delta_l}^{eV} \nu_l(x)\, dx,$$

and since the ratio I^1_{ni}/I_0 at typical values of the parameters $\bar{P}$ and σ_i is of the order of 10^{-4}, it is apparently quite difficult to observe such a smooth increment against the background of the nonlinear $I_0(V)$ dependence, at any rate for large junctions. The situation is more favorable for small junctions, where the discrete character of the TLS can become significant. The discreteness can lead in turn to the appearance of a small-scale structure in the dependence of dI_1/dV on V, which was possibly observed in some experiments.[11]

Realization of a nonequlibrium TLS population can lead, in our opinion, to intresting experimental possibilities. Such a nonequilibrium population can be reached by a number of methods, which were considered in detail for a tunnel junction in Ref. 3. The first method involves the action of the tunneling electrons that excite the TLS. The corresponding contribution δn_+ at $T = 0$ is of the form

$$\delta n_+{}^i \sim [E_i |K|^2/\mu^2 \max(\hbar/\tau_e, \hbar/\tau_{ph})] \frac{eV - E_i}{E_i} \theta(eV - E_i - \Delta_l), \tag{17a}$$

where τ_{ph} and τ_e are the TLS relaxation times governed by the phonons and bank electrons, respectively.

Another possible disequilibrium source is the electric dipole moment μ_e of the TLS and is due to the "polarization" of the electric field applied to the junction, i.e., to the transition of the TLS into states corresponding to an energy minimum when account is taken of the interaction with the field. When the electric field is turned off, the corresponding distribution is in a disequilibrium whose measure is

$$\delta n_+^i \sim \begin{cases} \left(\frac{\mu_{ei} V}{dT}\frac{\Delta_i}{E_i}\right)^2 2\,\mathrm{th}\left(\frac{E_i}{T}\right)\mathrm{ch}^{-2}\left(\frac{E_i}{T}\right), & \frac{\mu_e V}{dT} \ll 1. \\ \max\left(\frac{\mu_{ei}}{dT}, 1\right), & \frac{\mu_{ei} V}{dT} \gg 1. \end{cases} \quad (17b)$$

Finally, a disequilibrium can be produced by the "pumping" of the TLS by the microwave field. In this case the pumping is resonant; the populations become equalized at very low amplitudes of the microwave field potential, $V_m = V_m^0(\mu_e/d\hbar)(\tau_1\tau_2)^{1/2}$, where τ_1 and τ_2 are respectively the relaxation times of the populations and of the transverse "spin" components of the TLS.[12] At $\mu_e \sim 10^{-18}\ \mathrm{erg}^{1/2}\cdot\mathrm{cm}^{1/2}$, $\tau_1\tau_2 \sim 10^{-14}\ \mathrm{sec}^2$, and $d \sim 3\cdot 10^{-7}$ cm we have $V_m^0 \sim 10^{-6}$ V. The range of the energies corresponding to this disequilibrium is estimated at $|E_i - \hbar\omega| \lesssim (\hbar/\tau_2)(V_m/V_m^0)$.

The possibilities indicated are due to two circumstances. The first is that the TLS relaxation times are quite long (a typical value of τ_1 at $T\sim 1$ K and at $\rho \sim 1$ is 10^{-6} sec), so that after the pump is turned off the disequilibrium of the TLS is preserved much longer than the disequilibrium of the electron or phonon system. The second is the near-zero conductivity of the junctions at $eV < \Delta_l$ $(T = 0)$.

Thus, if a sufficiently strong electric field is initially applied to the junction and produces disequilibrium in a TLS system with energies $E_i \sim \Delta_l$ via the mechanism (17a) or (17b), and the field is decreased abruptly at $t = 0$ to values $eV < \Delta_l$, the junction conductivity at $t \gg t_0 = (R_{\mathrm{ext}} C)$ (R_{ext} is the external-circuit resistance and C is the junction capacitance), the junction conductivity will be due only to inelastic processes involving TLS (we assume $\tau_1 \gg t_0$). The corresponding relaxation current, taking (13) and (16) into acount, is then of the order of

$$I^1(t) \sim GV\sigma\bar{P}d\int_0^\infty dE\,\delta n_+(E)\int_0^1 d\rho\,(1-\rho)^{-1/2}e^{-t\rho/\tau_0} \sim GV\sigma\bar{P}d(\tau_0/t) \times \int_0^\infty dE\,\delta n_+(E), \quad (18)$$

where $\tau_0 \equiv \tau_1(\rho = 1)$. We took it into account that $\tau_1 \propto \rho^{-1}$; G is the tunnel conductivity of the junction in the normal state. At $\bar{P} \sim 10^{-33}\ \mathrm{erq}^{-1}\cdot\mathrm{cm}^{-3}$, $\sigma \sim 10^{-15}\ \mathrm{cm}^2$, $d \sim 3\cdot 10^{-7}$ cm, and $\int_0^\infty dE\delta n_+ \sim 3.10^{-16}$ erg we have

$$I^1(t) \sim 10^{-4} GV\tau_0/t.$$

Besides the nonexponentially relaxing current, the system can also be subject to a nonexponential relaxation of the junction voltage (due to relaxation of the TLS electric dipole moment[3]). Observation of the voltage relaxation is facilitated (compared with a normal-metal junction) by the junction conduction at $eV < \Delta_e$.

3. LOW-FREQUENCY FLUCTUATIONS IN JOSEPHSON STRUCTURES

We have shown above that the presence of TLS in the tunnel insulator renormalizes the expression for the Josephson current. The normalization depends, generally speaking, on the occupation numbers of the TLS levels [Eq. (9)]. In other words, the expression for the critical Josephson current is found to depend on the TLS states, so that transitions in the TLS are accompanied by fluctuations of this quantity, i.e., at a fixed phase difference $\chi_l - \chi_r \equiv \chi$, by fluctuations of the Josephson current I_s. We are most interested in low-frequency fluctuations connected with slowly relaxing TLS having $\rho \ll 1$. We take into account in (9) only the contribution of elastic processes, and rewrite this contrbution in the form

$$I_{se}^1 \sim \frac{I_{c0}}{A}\sin\chi\sum_i(\sigma_i^+ n_+^i + \sigma_i^- n_-^i), \quad (19)$$

where I_{c0} is the critical current in the absence of TLS. For the occupation-number fluctuation correlator we can obtain the expression

$$(\delta n_-^i, \delta n_-^j)_\omega = \frac{2}{\tau_1}\frac{n_0(1-n_0)}{\omega^2+\tau_1^{-2}}\delta_{ij}, \quad n_0 = F_0(-E_i). \quad (20)$$

Using (16) and (20) we obtain directly an expression for the correlator of the Josephson-current fluctuations:

$$\frac{(\delta I_s, \delta I_s)_\omega}{\sin^2\chi} \equiv (\delta I_c, \delta I_c)_\omega = I_{c0}\mathcal{F}(\omega), \quad (21)$$

$$\mathcal{F}(\omega) \approx \frac{d}{A}\overline{(\sigma^+-\sigma^-)^2}\frac{\bar{P}T}{\omega}, \quad (21a)$$

where the bar denotes configuration averaging. The spectrum of the fluctuations considered obeys thus the $1/f$ law (cf. Refs. 3 and 4).[2)]

The Josephson effect is known also to occur in systems of other type, such as point and bridge *S-C-S* junctions. We shall show that the presence of TLS in the vicinity of the weak link, just as in the case of an *S-I-S* junction, leads to low-frequency Josephson-current noise. To this end, we recognize first that the critical current of the structures in question is directly expressed in terms of their normal resistance R_N (see, e.g., the review[6]). The simplest situation is that of a point junction at $T \sim T_c$ (when the Ginzburg-Landau equation can be used); in this case[13]

$$I_s = I_c\sin\chi, \quad I_c = \frac{\pi}{4}\frac{\Delta_l\Delta_r}{eTR_N}. \quad (22)$$

Far from T_c, the dependence of I_s on χ is more complicated,[14] but the connection with R_N remains the same

$$I_s = \psi(\chi)I_c, \quad I_c \propto R_N^{-1}. \quad (23)$$

It can be shown[15] that if TLS are present in the vicinity of a point junction (or bridge) its resistance increases by an amount proportional to the occupation numbers of the TLS. The increment is estimated at

$$R_N^1 \sim R_N\frac{l}{r_0^3}\sum_i(\sigma_i^+ n_+^i + \sigma_i^- n_-^i), \quad (24)$$

where r_0 is the junction dimension, l is the mean free path of the electrons, while σ_i^+ and σ_i^- are the cross sections for electron scattering by a TLS located on an upper and lower level, respectively. With allowance for (22)–(24) we obtain

$$\frac{(\delta I_s, \delta I_s)_\omega}{\psi^2(\chi)} \equiv (\delta I_c, \delta I_c) = I_{c0}\mathcal{F}(\omega), \quad (25)$$

$$\mathscr{F}(\omega)=\left(\frac{l}{r_0^3}\right)^2\sum_i\frac{(\sigma_i^+-\sigma_i^-)^2\tau_{1i}n_{0i}(1-n_{0i})}{\omega^2\tau_{1i}^2+1}. \qquad (25a)$$

We have retained in this expression the summation over i, since in the case of point junctions the allowance for the discreteness of the set of TLS can be quite significant. We note that the time τ_1 in (25) is determined by the interaction of the TLS with the electrons.

An important feature of this noise mechanism is, as seen from the expressions obtained, that it is not directly connected with the quasiparticle component of the electron system, so that I_s fluctuates also at $T\ll T_c$ (the ratio $\delta I_s, \delta I_s)/I_c^2$ depends on T only via the TLS occupation numbers), whereas $(\delta I_s, \delta I_s)$ is linear in T.

We examine now how the fluctuations can be observed in experiment. First, they can manifest themselves directly in measurements for which a fixed value of the critical current is essential. It appears, in particular, that the noise can limit the sensitivity of guantum interferometers (SQUID's).

If we connect the junction into a closed superconducting ring, the presence of I_c fluctuations leads ultimately to fluctuations of the magnetic-flux Φ linked with the ring. Using the known relation for a ring with weak coupling

$$\chi+2\pi IL/\Phi_0=2\pi n=\text{const},$$

where I is the current in the ring, L the ring inductance, and Φ_0 is the magnetic-flux quantum, and using also the expression $I=I_c\psi(\chi)$ for the current (where we take into account the possibility that the functions ψ differ from $\sin\chi$ at an arbitrary weak coupling), we easily obtain, assuming $2\pi I_{c0}L\psi'(\chi)\gg\Phi_0$, the relation

$$(\delta\Phi,\delta\Phi)_\omega\approx\left[\frac{\Phi_0}{2\pi}\frac{\psi(\chi)}{\psi'(\chi)}\right]^2\frac{(\delta I_c,\delta I_c)_\omega}{I_{c0}^2}. \qquad (26)$$

We call attention to the following interesting question, which we shall illustrate with the S-I-S junction as the example. A change in the values of Φ, χ, and I_c in a weakly coupled ring is accompanied by a change of its energy E_s which is the sum of the energy of the magnetic flux Φ 1/2 and of the weak-coupling energy[10]:

$$(\hbar I_c/2e)\,(1-\cos\chi);$$

this change can be readily seen to be equal to

$$\delta E_s=(\hbar/2e)\,\delta I_c(1-\cos\chi).$$

Since, at any rate in the limit $T\ll T_c$, the value of I_c is determined by purely elastic processes, the situation with the energy balance is not clear at first glance. It can be shown, however (see the Appendix), that the renormalization of the TLS energy E_i, due to the interaction with the electrons, contains in this case a contribution that depends on χ. The change that must take place in the energy E_i needed for the TLS transition between levels makes up the corresponding weak-coupling energy deficit.

Under conditions of the nonstationary Josephson effect the junction-resistance fluctuations due to the TLS manifest themselves directly as fluctuations of the junction voltage (at $T\sim T_c$ the expression for the static current differs little from the corresponding expression in a normal metal[3,15]). Taking the foregoing into account we obtain the relation

$$(\delta V,\,\delta V)_\omega/V^2\sim\mathscr{F}(\omega) \qquad (27)$$

where $\mathscr{F}(\omega)$ is determined by (25a) for an S-C-S junction and by (21a) for an S-I-S junction), which is valid at any rate at $T\sim T_c$. We note that in this situation it is possible to take account of the TLS disequilibrium due to the interactions of the systems with the electrons. This pertains above all to the case of point S-C-S junctions, when the principal TLS relaxation mechanism is connected with the electron system[15]. The result of this circumstance reduces to replacing the temperature in (25) at $eV>T$ by the energy eV. Obviously, the considered Josephson-junction voltage fluctuations manifest themselves, in particular, as a fluctuational modulation of the Josephson generation frequency ω_0, and the corresponding broadening in frequency is estimated in order of magnitude at

$$\Delta\omega_D\sim\int_{\omega_c} d\omega\,\frac{2e}{\hbar}(\delta V,\delta V)_\omega^{1/2}$$

(where the cutoff freguency ω_c is determined either by the observation time or by the TLS distribution cutoff parameter ρ_c).

We present a few quantitative estimates. It can be seen from (21a) and (25a) that the fluctuations considered are particularly important for S-C-S junctions (since the number of TLS in the junction region is small by virtue of the small-dimensions). Putting $\bar{P}\sim10^{32}-10^{33}\,\text{erg}^{-1}\cdot\text{cm}^{-3}$, $(\sigma^+-\sigma^-)^2\sim10^{-30}\,\text{cm}^4$, and $r^0\sim3\cdot10^{-6}$ cm, we get $\mathscr{F}(\omega)$ ω of the order of 10^{-8}–10^{-7}. A quantitative comparison of the fluctuations due to the TLS with Nyquist fluctuations at $T\sim T_c$ (or at $I>I_c$) for an S-C-S junction with $R_N\sim(0.1$–$1)\Omega$ shows that at $\omega\lesssim(10^5$–$10^6)$ cm^{-1} the low-frequency fluctuations can exceed the Nyquist fluctuations already at $I\sim(1$–$10)$ μA. This estimate agrees in order of magnitude with the experimental estimate of the $1/f$ noise for point junctions.[6]

4. CERTAIN SPECIAL PHENOMENA IN JOSEPHSON STRUCTURES

We consider first of all phenomena in a superconducing ring closed by a Josephson S-I-S junction with an amorphous film in the case when a nonequilibrium population of the TLS was produced by some method at the instant $t=0$. (Such a nonequilibrium population can be produced, for example, by applying a voltage to the junction and turning it off abruptly at the instant $t=0$; in the system considered here this can be effected, in particular, by a suitable change of the external magnetic field.) Since the relaxation times in the electron (and phonon) system are much shorter than for the TLS, we assume that at $t>0$ all the disequilibrium is due to the indicated nonequilibrium population $\delta n_-(E)$ that relaxes slowly to zero. Turning to Eq. (9) for the Josephson current, we note that in our situation it contains a contribution

$$\sim\sin\chi\frac{I_{c0}}{A}\sum\rho_i\sigma_i\delta n_-^{\,i}(E_i)\,\varphi(\Delta,E_i)$$
$$\sim\sin\chi I_{c0}\sigma\bar{P}d\int_0^\infty dE\,\delta n_-(E)\,\varphi, \qquad (28)$$

where $\varphi(\Delta, E_i)$ is a cetain function such that at $E_i \sim \Delta$ we have $|\varphi| \sim 1$; we disregard it in our order-of-magnitude estimate. The relaxation in the set of TLS will thus be accompanied by a change in the value of I_c, which leads as a result to relaxation of the magnetic flux Φ linked with the ring. For the corresponding variable part of Φ and within the framework of a calculation similar to (18) we obtain the estimate

$$\delta\Phi \sim \frac{\Phi_0}{2\pi} \operatorname{tg} \chi \bar{P}\, d\sigma \left(\int_0^\infty dE \delta n_-(E) \tau_0(E) \right) \frac{1}{t}. \qquad (29)$$

We note that although our calculation was performed for an *S-I-S* junction, a similar behavior can be expected in the case of other Josephson structures.

We propose here also to use the nonstationary Josephson effect to investigate the echo phenomenon, which is among the most interesting features of the behavior of glasses in alternating fields. This effect can serve as one of the most convincing proofs of the existence of TLS. Its gist is that if two (or more) ac electric fields of definite frequency, duration, and amplitude are applied to a glass and are separated by some time interval, a coherent response pulse, or echo signal, is obtained after a definite time lapse.[12] We shall not discuss in detail the mechanism of this phenomenon, which has been described in the literature many times (see, e.g., Ref. 12). We indicate only that in the case of the simplest (two-pulse) echo the first ac pulse with amplitude $\mathscr{E}_0$ and duration Δt must satisfy the relation $2\mu_e \mathscr{E}_0(\Delta t/\hbar) = \pi/2$, while the second (which follows the first after a time T_1) should have a doubled duration or amplitude (compared with the first). The echo signal is produced at a time T_1 after the second pulse.

A Josephson junction with an amorphous insulator constitutes, in our opinion, a system that permits direct observation of the echo. The exciting pulses can be produced by applying to the junction voltage pulses of suitable duration Δt. The signal echo consists of an alternating field that appears in the barrier after a corresponding time interval and is due to the oscillating dipole moments of the TLS. As a result, a voltage pulse of duration Δt is produced on the junction; its amplitude can be estimated at

$$\tilde{V}_e \sim (4\pi/\varepsilon)\, \mu_e \bar{P}\, (\hbar/\Delta t)\, d.$$

We point out that in estimates of the time interval during which the echo is observable it is necessry to take into consideration, besides the usual relaxation mechanisms,[12] also the interaction of the TLS with the bank electrons.[3]

The echo can apparently be observed with the aid of Josephson bridges deposited on substrates of an amorphous material (insulator). The alternating fields generated in the bridge penetrate into the insulator, and the characteristic scale of the region whose TLS interact effectively with the indicated fields is determined by the dimension (length) r_0 of the bridge. In this case the duration of the first pulse is given by $\sim[4\mu_e \tilde{V}/(r_0\pi\hbar)]^{-1}$, where $\tilde{V}$ is the amplitude of the Josephson voltage, while the amplitude of the echo response is of the order of $\tilde{V}_e \sim (4\pi/\varepsilon)\mu_e \bar{P} \times (\hbar/\Delta t) r_0$. We note, however, that since the electric field produced in the substrate by the bridge is not uniform, the echo signal can be greatly broadened.

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APPENDIX

Renormalization, due to superconducting effects, of the parameters of TLS in the insulating layer of an *S-I-S* Josephson junction

Starting from the Hamiltonian (1), we obtain a system of equations for the components of the TLS "spin" *S*. Proceeding as in Ref. 3 and taking into account only the terms that make the principal contribution to the sought renormalizations, we have

$$\begin{aligned} \frac{d\langle S_x\rangle}{dt} &= \frac{i}{\hbar}\left[iE_i\langle S_z\rangle + \ldots + i\frac{\Delta_i}{E_i} J\langle S_y\rangle \right], \\ \frac{d\langle S_y\rangle}{dt} &= \frac{i}{\hbar}\left[-iE_i\langle S_x\rangle + \ldots - \frac{\Delta_i}{E_i} J\langle S_x\rangle + i\frac{\Delta_{0i}}{E_i} J\langle S_z\rangle \right], \\ \frac{d\langle S_z\rangle}{dt} &= \frac{i}{\hbar}\left[\ldots - i\frac{\Delta_{0i}}{E_i} J\langle S_y\rangle \right], \end{aligned} \qquad (A.1)$$

where

$$J = -\sum_{pp'} \left[\frac{1}{2} + \frac{1}{2}\frac{\Delta_l}{\varepsilon_l}\frac{\Delta_r}{\varepsilon_r}\cos\chi \right] \times \frac{1}{\varepsilon_{pl}+\varepsilon_{pr}} \frac{E_i}{\Delta_i} \left[\left| T+U+\frac{1}{2}\frac{\Delta_i}{E_i}K \right|^2 - \left| T+U-\frac{1}{2}\frac{\Delta_i}{E_i}K \right|^2 \right]. \qquad (A.2)$$

Comparing (A.2) with (9), we call attention to the fact that the J contribution that depends on $\chi_l - \chi_r$ is equal to (we assume for simplicity that $T \ll T_c$)

$$J = -\frac{\hbar}{2e}\cos(\chi_l - \chi_r)\,[I_c^+ - I_c^-]\frac{E_i}{\Delta_i},$$

where $I_c^+ - I_c^-$ is the difference between the critical currents corresponding to the position of the TLS on the upper and lower levels, respectively. It can be seen from (A.1) that the presence of the terms considered is equivalent to a renormalization of the TLS parameters, $\delta\Delta_i = J$ and $\delta E_i = (\Delta_i/E_i)J$, whence it follows directly that the contribution that depends on the phase difference is

$$\delta E' = -\frac{\hbar}{2e}\cos(\chi_l - \chi_r)\,[I_c^+ - I_c^-].$$

[1)]We assume that the barrier molecules have a dipole moment, and there is no correlation between the orientations of these dipoles. It was found that an important role can be played in the situaticn considered by spatial fluctuations of the barrier dipole moment; an alternating electric field interacting with these flucuaticns excites sounds.

[2)]We note that an $1/f$ law is obtained only if the integral with respect to ρ diverges all the way to $\rho = 0$. This implies, first, a corresponding TLS distribution in the Anderson model, and second, infinite dimensions of the junction. In the opposite case the "cutoff" with respect to ρ takes place at frequencies $1/\tau_{\max}$, where $\tau_{\max}$ corresponds to the most slowly relaxing TLS of the given particular junction. Using (16), however, it is easy to show that at $\overline{P}dAT \gg 1$ the value of $\tau_{\max}$ is with overwhelming probability $\tau_0 \rho_c^{-1}$, where ρ_c is the TLS-distribution cutoff parameter

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