wave packets.

Owing to the growth of $|\tau| \sim E/m\Gamma$ with increasing energy in (3) at high energies (E - energy of LW particles), we can expect the appearance of macroscopic causality-violation effects. However, if $\vec{u} \approx \vec{v}$ then, as follows from (3), R = vt and no noticeable non-causal effects appear. Interest attaches therefore to the case when the velocities u and v do not coincide. At high energies this can occur only if particle 1 has zero mass.

Let us consider, for example, the case when

$$\vec{u} = -\vec{v}, \quad |\vec{v}| = c = 1$$

corresponding to the decay of an LW particle with emission of a photon or a neutrino in a direction opposite to that of the motion of the LW particle. It then follows from (3) that

$$r = t + 2|\tau|.$$

Thus, the effective velocity ($v_{eff} \approx 1 + 2|\tau|/t$) of the propagation of a signal between two events, namely the collision and the instant of registration, may turn out to be noticeably larger than the speed of light at high energies. For example, if Γ = 50 meV, m = 10 GeV, and the energy is $E \sim 10^{23}$ eV we have $|\tau| \sim E/m\Gamma \sim$ cm, i.e., $|\tau|/t \sim 1$ at $t \sim 1$ cm.

The wavelength of the first particle, at the decay kinematics under consideration, equals $\lambda \sim E/m^2$, i.e., $\lambda/|\tau| \sim \Gamma/m \ll 1$. This means that particle 1 can in principle be localized with an accuracy sufficient for the observation of non-causal effects.

It seems to us that the possibility, in principle, of appearance of macroscopic causality violation at high energies is more readily an argument against the theory developed by Lee and Wick.

In conclusion, the authors thank B.L. Ioffe for many useful discussions.

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FLUCTUATION CONDUCTIVITY OF A TUNNEL CONTACT AT TEMPERATURES ABOVE CRITICAL

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It is known that at temperatures above T_c the superconducting-ordering parameter Δ does not vanish identically, but fluctuates with a certain (small) amplitude) whose value increases as T_c is approached. As shown in [1], this leads to additional conductivity of the fluctuation pairs, equal in the case of thin film to

$$\sigma' = \frac{e^2}{16\hbar d}\frac{T_c}{T - T_c}.$$

We consider in this paper an analogous effect for a tunnel contact between two superconductors. In the case when the voltage across the barrier V does not vanish and $T > T_c$, the fluctuation conductivity of the contact oscillates as a function of the time with the Josephson frequency $\omega = 2eV/\hbar$. (A more detailed exposition of the theory of this effect will be published in the collection " Fizika kondensirovannogo sostoyaniya" (Physics of the Condensed State), No. 4, Physico-technical Institute of the Ukrainian Academy of Sciences, Khar'kov.)

We consider the contact between two identical superconductors with total resistance R. The transverse dimensions of the system are assumed small compared with the coherence length

$$\xi(T) \sim v_0/(T_c|T_c - T|)^{1/2},$$

and accordingly the ordering parameters Δ_1 and Δ_2 of the superconductors can be regarded as independent of the coordinates: $\Delta_1 = \Delta_1(t)$, $\Delta_2 = \Delta_2(t)$. The fluctuations of $\Delta_i(t)$ will be considered in the self-consistent field approximation, meaning that the final results are valid outside the immediate vicinity of T_c (far from the "critical region").

Using the scheme of Ambegaokar and Baratoff [2], let us calculate the superconducting current in the contact I(t), assuming $\Delta_i(t)$ to be arbitrary functions of the time. This yields

$$I(t) = \frac{1}{eR}|\Delta_1(t)\Delta_2(t)| I_m \int_{-\infty}^{t} e^{i\chi(t)} e^{i\chi(t')} \ln\mathrm{cth}\,\frac{\pi(t-t')}{2\beta}dt', \tag{1}$$

where

$$\beta = 1/T_c, \quad \chi(t) = e/\hbar \int_{t(0)}^{t} V(t')dt'.$$

In the derivation of (1) we took into account the fact that near T_c the characteristic time of variation of Δ is large in comparison with $\hbar/T_c$. For a slow variation of the potential V(t) we get from (1)

$$I(t) = \frac{\pi}{4eRT_c}|\Delta_1(t)\Delta_2(t)| I_m [e^{2i\chi(t)} F(\frac{eV(t)}{2\pi T_c})], \tag{2}$$

where the function F(z) has the following meaning:

$$F(z) = \frac{2}{\pi^2 iz}[\psi(\frac{1}{2} + iz) - \psi(\frac{1}{2})] \tag{3}$$

($\psi(z)$ is the logarithmic derivative of the Gamma function). At z = 0 we have F(0) = 1.

Let us consider for simplicity the case of voltages small compared with T_c. As seen from (2), the superconducting current of the contact is a periodic function of the time with frequency $\omega = 2eV/\hbar$ and with a randomly varying amplitude, the mean value of which is $I_m = \pi|\overline{\Delta_1}||\overline{\Delta_2}|/4eRT_c$. The quantity $|\Delta|$ is defined as $\mathrm{Tr}\, e^{-\beta F}|\Delta|)/\mathrm{Tr}\, e^{-\beta F}$, where

$$F = \frac{\hbar^2 C}{2m\xi^2(T)}\int|\Delta|^2 dV = \frac{\hbar^2}{2m\xi^2(T)}CSd|\Delta|^2.$$

Here d is the thickness of the film ($d \ll \xi(T)$), and $C = 7\zeta(3)N/8(\pi T_c)^2$ (see [3]). As a result we obtain for the average amplitude of the Josephson current the estimate

$$I_m \approx \frac{m\xi^2(T)}{2eR\hbar^2 CSd} \sim \frac{T_c}{eR}\frac{mT_c}{\hbar^2 Nd}\frac{\xi_0^2}{S}\frac{T_c}{T - T_c} \tag{4}$$

(S is the area of the contact surface, $S \ll \xi^2(T)$). As expected, I_m increases when T_c is approached.

To determine the spectral composition of the superconducting component of the current, let us calculate the correlator $K(t, t') = \overline{I(t)I(t')}$. At the same time, we take into account the fact that both $\Delta(t)$ and the contact voltage $V(t) = V_0 + v(t)$ fluctuate. The fluctuations of v will be assumed to be "white noise," and then we have in accord with Nyquist's theorem $\overline{v(t)v(t')} = 2RT\delta(t - t')$. Using the time-dependent Ginzburg-Landau equation in the form given in [4], we obtain for the correlation function the expression $\overline{|\Delta(t)||\Delta(t')|} = |\Delta|^2 e^{-\Gamma|t - t'|}$, where $\Gamma = 8/\pi\hbar(T - T_c)$.[1)] It is now easy to calculate $K(t - t')$, and we obtain for the Fourier component of this function the formula

$$K(\omega) = \frac{\pi^2}{4} I_m^2 \frac{\Gamma^*}{(\omega - \omega_0)^2 + 4\Gamma^{*2}}, \quad \omega_0 >> \Gamma^*, \tag{5}$$

where

$$\Gamma^* = \Gamma + \frac{2e^2}{\hbar^2} RT_c .$$

The spectral density of the radiation power at the frequency ω_0 is estimated at

$$\frac{\Delta W}{\Delta\omega} = \frac{\pi^2}{16} \frac{RI_m^2}{\Gamma^*}. \tag{6}$$

The quantity Γ^* can be represented in the form

$$\Gamma^* = \frac{8}{\pi\hbar}(T - T_c^*),$$

where the "renormalized" critical temperature is

$$T_c^* = T_c - \frac{\pi e^2}{4\hbar} RT_c .$$

We see from (5) that the radiation power has a maximum at a value of ω equal to the frequency of the Josephson oscillations $\omega_0 = 24V_0/\hbar$. Unlike the Josephson effect, however, this radiation has now a rather broad spectrum determined by the value of Γ^*. We present some numerical estimates. At $T_c = 10°K$, $T - T_c = 10^{-3}$ °K, $d \sim 10^{-7}$ cm, $S \sim \xi_0^2 \sim 10^{-8}$ cm and R = 1 ohm we obtain for the quantity $\Delta W/\Delta\omega$ at the maximum of the spectral line the estimate $\Delta W/\Delta\omega \sim 10^{-16}$ W/Hz. The minimum value of Γ^* at the indicated values of the parameters is $\Gamma^* \sim 10^9$ sec^{-1}. Assuming a bandwidth $\Delta\omega \sim 10^8$ Hz, we obtain for the total radiation power the estimate $W \sim 10^8$ W. The power radiated to the outside space is (much) lower than this value, and is determined by the ratio of the impedances of the strip line and the vacuum [6].

Besides radiation from the contact proper, the investigated effect can become manifest also in the form of current "steps" produced when the contact is irradiated by an external high-frequency power of frequency Ω, and occurring at voltage values $V_n = n\hbar\Omega/2e$ [7]. The heights of such steps are determined by the value of I_m.

It follows thus from our results that certain features of the nonstationary Josephson effect are retained in the normal state at temperatures somewhat higher than T_c. Our estimates show that in principle this effect can become experimentally observable.

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1) We note that above T_c we can regard our system as a gapless superconductor. This means that we can apply to it the considerations developed by Gor'kov and Eliashberg [5], concerning the existence of a time-dependent Ginzburg-Landau equation.

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ELECTRIC PROPERTIES OF A SEMICONDUCTOR IN THE FIELD OF A STRONG ELECTROMAGNETIC WAVE

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1. In [1, 2] we investigated the electric and magnetic properties of a semiconductor in the field of a strong electromagnetic wave $\vec{A}(t) = \vec{A}\cos\omega t$ with a frequency ω exceeding the width Δ of the forbidden band. It was shown that the stationary state of the system is a state of saturation, wherein the strong-field absorption coefficient is equal to zero (in the absence of recombination) (see [3]). In addition, it was shown that in the electron (hole) spectrum there is a gap near the quasimomentum p_0 (see Fig. 1)

$$\lambda_p = \frac{1}{2} e A v_{cv} p_0 = \sqrt{m(\omega - \Delta)}, \quad \hbar = c = 1, \tag{1}$$

where $\vec{v}_{cv}$ is the matrix element of the transition between the conduction and valence bands; the effective masses of the electrons and holes are assumed to be equal.

We consider in this paper the absorption of a weak electromagnetic field $\vec{A}_1(t) = \vec{A}\cos\omega_1 t$ with frequency $\omega_1 \sim \omega$,[1)] and also direct current in a semiconductor in the saturation state.

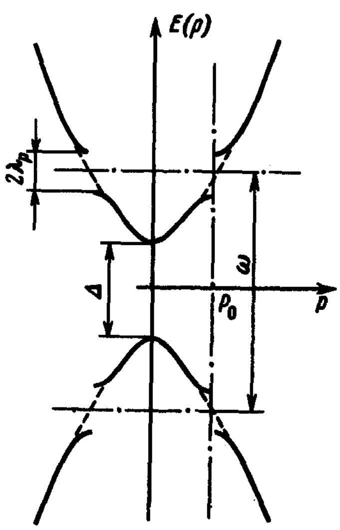

Fig. 1

2. Using standard perturbation-theory methods, we obtain the average Fourier component of the current in the first approximation in the weak field $\vec{A}_1$

$$j_i(\omega_1 q) = K_{ij}(\omega_1 q) A_{1j}(\omega_1 q)$$

where the imaginary part K_{ij}, describing the absorption is given at $Q \to 0$ and $T = 0$ by (see Eq. (26) of [1])

$$K_{ij}(\omega_1) = \frac{\pi e^2}{2} \sum_p v^i_{cv}(p) v^j_{cv}(p) [u_p^4 \delta(2\epsilon_p + \omega - \omega_1) - v_{-p}^4 \delta(2\epsilon_p - \omega + \omega_1)], \tag{2}$$

where

$$u_p^2, v_{-p}^2 = \frac{1}{2}\left(1 \pm \frac{p^2 - p_0^2}{2m\epsilon_p}\right), \quad \epsilon_p = \sqrt{\left(\frac{p^2 - p_0^2}{2m}\right)^2 + \lambda_p^2}.$$

1) A similar method is used in gases, as pointed out to us by I. I. Sobel'man.