i.e., it is smaller than the contribution of the 2P branch cut (the latter is $\sim 1/\xi$). Thus, although enhanced diagrams do lead to a broadening of the peaks and to a decrease of the valleys between them, they cannot, generally speaking, smooth out the peaks completely.

It is of interest to determine the ratio of the values of σ_n at the peak "$a\xi$" to the values of σ_n at the peak "$2a\xi$." It can be easily estimated by starting from the ratio of the contribution of the (PP) branch cut to the contribution of P to σ_{tot}; we obtain

$$\frac{\sigma(2a\xi)}{\sigma(a\xi)} \approx \frac{1}{\sqrt{2}} \frac{c^2\sigma_{tot}}{16\pi(2a'\xi + 2b)}\left(\frac{N'}{N}\right)^2 \approx \frac{1}{10} c^2\left(\frac{N'}{N}\right)^2, \qquad (6)$$

where c determines the deviation of the quantity N (the vertex of the emission of two P by a hadron) from diagonal ($c^2 \approx 1.3 - 1.8$; see [5]). The ratio $(N'/N)^2$ enters in (6), since when the reggeons are cut, generally speaking, a change takes place in the value of the vertex N. The factor $1/\sqrt{2}$ in (6) takes into account the broadening of the "$2a\xi$" peak.

Observation of oscillations in σ_n in experiment would be of very great interest[1]. The presently available data [6] on the distributions of σ_n contain excessively large errors, and it is therefore desirable to refine them.

In conclusion we thank V.N. Gribov, I.D. Mandzhavidze, S.G. Matinyan, and K.A. Ter-Martirosyan for useful discussions.

[1] D. Amati, S. Fubini, and A. Stangellini, Nuovo Cim. 26, 896 (1962).
[2] A.H. Mueller, Phys. Rev. 4D, 150 (1971).
[3] V.N. Gribov, Zh. Eksp. Teor. Fiz. 53, 654 (1967) [Sov. Phys.-JETP 26, 414 (1968)].
[4] V.N. Gribov and A.A. Migdal, Yad. Fiz. 8, 1002 (1968) [Sov. J. Nucl. Phys. 8, 583 (1969)].
[5] A.B. Kaidalov, ibid. 13, 401 (1971) [13, 226 (1971)].
[6] L.W. Jones, A.E. Bussian, G.D. DeMeester, et al., Phys. Rev. Lett. 25, 1679 (1970).

QUASIFIELD IN SUPERCONDUCTORS AND INCREASE OF T_c UNDER NON-EQUILIBRIUM CONDITIONS

A.G. Aronov and V.L. Gurevich
A.F. Ioffe Physico-technical Institute, USSR Academy of Sciences
Submitted 25 March 1972
ZhETF Pis. Red. 15, No. 9, 564 - 567 (5 May 1972)

The superconducting transition is a second-order phase transition. In this sense it is analogous, for example, to the transition in magnetically ordered systems. It is known that when an external magnetic field is applied to a ferromagnet, the phase transition vanishes, and the ordering parameter (the magnetization) remains different from zero at all temperatures.

---

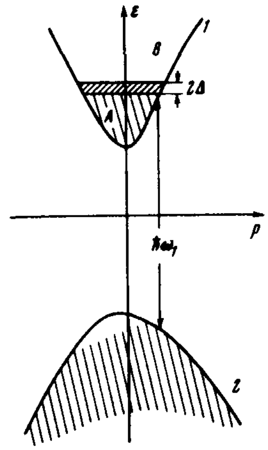
[1] We note that the data of [6] at $E_{lab} \approx 424$ GeV point to a possibility of oscillations of σ_n (peaks at $n \approx 6$, 10, and 14).

Is there an analog of this external field for superconductors? We shall show that such a "quasi-field" can exist under non-equilibrium conditions.

We consider a semiconductor or a semimetal with the band scheme shown in the figure. Under the influence of light with a definite spectral composition (see the figure), the electrons go from the valence band 2 into region A of the conduction band 1, which lies below a superconducting gap of width 2Δ. This frees places near the top of the valence band 2.

The electrons[1)] flow out of the region A of the conduction band via two channels. They go over principally from the bottom of band 1 to the free places of band 2 via direct radiative transitions[2)]. The corresponding time of spontaneous emission will be denoted by τ_{A2}. In addition, by absorbing a phonon with energy larger than 2Δ, they can go over into region B of band 1, which lies above the superconducting gap. The corresponding characteristic time τ_{AB} will be assumed large (e.g., because the width of the gap is larger than the characteristic phonon energy).

The outflow of electrons from region B is due also to two causes. First, these are phonon emission processes, which cause the electrons to return to the region A (the corresponding time τ_{BA} will also be assumed large in comparison with the time τ_B of the electron-phonon collisions inside the region B). Second, these are indirect transitions from the region B immediately to the freed places of band 2, with a characteristic time τ_{B2}.



Band scheme and scheme of transitions under the influence of light. The light contains the frequencies $\omega < \omega_1$. Near the top of band 2 a certain number of empty places is freed as a result of the excess of electrons in band 1 in comparison with the equilibrium value.

Let us examine the stationary state, when the electron distribution does not depend on the time. Since the electron-phonon collisions rapidly establish an equilibrium inside each of the regions A and B, and A $\rightleftarrows$ B transitions are rare, the distribution function in each of the regions is a Fermi function, but the chemical potentials μ_A and μ_B are different. The average value $\bar{\mu} = (\mu_A + \mu_B)/2$ increases with increasing total number of electrons in band 1. It can be shown that the position of the center of the gap in the spectrum of the single-particle excitations coincides precisely with this quantity. Then the distribution functions of the excitations below and above the gap will be

$$F_p^{(A,B)} = \left(\exp\frac{\zeta_p \mp \delta\mu}{T} + 1\right)^{-1}; \quad \zeta_p = \mp\sqrt{\xi_p^2 + \Delta^2} \qquad (1)$$

$\xi_p = (p^2/2m) - \mu$, p is the quasi-momentum, and m is the effective mass. The quantities $\delta\mu$ and $\bar{\mu}$ are determined from the corresponding balance conditions.

1)Actually, it would be more correct to speak here not of electrons but of excitations in the superconductor.

2)Indirect transitions also take place, but in this region they can be neglected in comparison with the direct transitions.

Using Bogolyubov's well-known procedure [1][3], we obtain an equation relating the gap with the distribution function of the excitations

$$1 = \frac{g}{2}\left[\int_{p > p_F} \frac{d^3p}{(2\pi\hbar)^3} \frac{1 - 2F_p}{\sqrt{\xi_p^2 + \Delta^2}} - \int_{p < p_F} \frac{d^3p}{(2\pi\hbar)^3} \frac{1 - 2F_p}{\sqrt{\xi_p^2 + \Delta^2}}\right] \tag{2}$$

or, substituting for F_p the function (1)

$$1 = \frac{g}{2}\int \frac{d^3p}{(2\pi\hbar)^3} \frac{1}{\sqrt{\xi_p^2 + \Delta^2}} \,\mathrm{th}\left[\frac{\sqrt{\xi_p^2 + \Delta^2}}{2T} + \frac{\delta\mu}{2T}\right]. \tag{2a}$$

The integration, as usual, is carried out over the energy region $2\hbar\omega_D$ (ω_D is the Debye-phonon frequency) and g is the effective constant of the electron-electron interaction.

At $\delta\mu \neq 0$ and $\Delta \to 0$, the integral in the right-hand side of (2a) diverges. This means that Eq. (2) has a non-zero solution Δ at any temperature T. Thus, $\delta\mu$ plays the role of the quasi-field.

At $\delta\mu/T \gg 1$, the solution of (2a) in first approximation is of the form

$$\Delta = \Delta_0\left[1 - 2e^{-\delta\mu/T} K_0(\Delta/T)\right], \tag{3}$$

where Δ_0 is the gap at $\delta\mu = T = 0$, and $K_0(x)$ is a Macdonald function. This case was in fact considered by the authors earlier [2].

At $\delta\mu/T \ll 1$ and $\Delta \ll T \ll \hbar\omega_D$ we have

$$\Delta = T\left(\frac{T_c}{T}\right)^{2T/\delta\mu} = T\exp\left(-\frac{2T}{\delta\mu}\ln\frac{T}{T_c}\right), \tag{4}$$

where T_c is the phase transition point at $\delta\mu = 0$. Relation (4) is valid when $T > T_c$.[4]

We make a few concluding remarks.

1) It is extremely difficult to satisfy inequalities of the type $\tau_{BA} \gg \tau_B$ in experiments. If there is no strong inequality, then the simple formulas of type (4) are no longer obtained. One might assume, however, that even in this case, in some region of temperatures and intensities, the effect will remain in force, since its existence is formally connected with the divergence of the right-hand side of Eq. (2), which is due to the specific non-equilibrium character of the distribution function in the presence of a gap.

---

[3] Estimates show that at reasonable light intensities the corrections to the excitation spectrum due to the field of the light wave are still small, as is assumed in the derivation of (2).

[4] The question of raising T_c of a superconductor in a strong microwave field was considered earlier by Eliashberg [3].

2) The indicated non-equilibrium character of the distribution function can exist only because of the gap. This means in turn that the considered state is metastable and can be obtained by first illuminating an equilibrium superconductor, and then raising the temperature.

3) The calculation does not take into account the finite lifetime of the Cooper pairs. This account should decrease the value of the ordering parameter Δ, and it is possible that the superconductivity at $T > T_c$ will be gapless.

The authors are grateful to A.I. Larkin, V.S. Maleev, G.E. Pikus, and G.M. Eliashberg for interesting discussions.

[1] L.D. Landau and E.M. Lifshitz, Statisticheskaya fizika (Statistical Physics), Nauka, 1964, p. 300.
[2] A.G. Aronov and V.L. Gurevich, Fiz. Tverd. Tela 14, 1129 (1972) [Sov. Phys.-Solid State 14, No. 3 (1972)].
[3] G.M. Eliashberg, ZhETF Pis. Red. 11, 186 (1970) [JETP Lett. 11, 114 (1970)]; Zh. Eksp. Teor. Fiz. 61, 1254 (1971) [Sov. Phys.-JETP 34, 668 (1972)].

CONTRIBUTION TO ELECTRODYNAMIC THEORY OF VAN DER WAALS FORCES BETWEEN MACROSCOPIC BODIES

Yu.S. Barash and V.L. Ginzburg
P.N. Lebedev Physics Institute, USSR Academy of Sciences
Submitted 20 March 1972
ZhETF Pis. Red. 15, No. 9, 567 - 571 (5 May 1972)

Van der Waals forces between macroscopic bodies (e.g., between two plates - regions 1 and 2) can be calculated by an electrodynamic method using the fluctuation-dissipation theorem. The corresponding theory was initially developed (see [1, 2], Sec. 92) as applied to two half-spaces with parallel boundaries, separated by an empty gap (region 3). A generalization of the calculations to the case of a gap filled with a medium was realized, however, only by using the complicated formalism of the temperature Green's functions. At the same time, the results obtained in [3] for filling of a gap with an arbitrary medium were recently obtained [4, 5] by an incomparably simpler procedure. Namely, the free energy F or the internal energy U of the system (media 1, 2, and 3) are represented in the form of sums of contributions of harmonic oscillators with surface-oscillation (wave) frequencies ω_α, corresponding to the problem in question. For example, it is assumed that

$$U(\ell) = \sum_\alpha \phi(\omega_\alpha, T), \quad \phi(\omega_\alpha, T) = \frac{\hbar\omega_\alpha}{2} \operatorname{cth} \frac{\hbar\omega_\alpha}{2kT}, \tag{1}$$

where ℓ is the width of the gap 3, on which the frequencies ω_α depend, and the index α combines both the discrete variables and the wave vector k in the plane of the gap. Knowing the frequencies $\omega_\alpha(\ell)$, we obtain $U(\ell)$ or $F(\ell)$, and by differentiating with respect to ℓ we subsequently obtain the force $f(\ell)'$, which coincides with that obtained in [1 - 3]. In principle it is easy to generalize the method to the case of more complicated configurations; in the particular case of ideally conducting plates in vacuum, this method reduces to a procedure already long in use [7] (see also [8, 9]). Both in this limiting case, and in the more general case, when all the media 1, 2, and 3 are non-absorbing, the meaning of the expression (1) is obvious, and it is obtained formally by standard quantization of the field in the medium (see, e.g., [10]). The theory of