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THE WIDTHS OF SINGLE PARTICLE STATES OF ANISOTROPIC STRONGLY CORRELATED ELECTRON SYSTEMS IN SOLIDS

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The damping $\gamma(\varepsilon)$ of electron states in crystals is investigated beyond the phase transition point related to a rearrangement of the Fermi surface. The attention is focused on the alteration of the standard Landau theory due to the emergence of a flat portion in spectrum $\xi(\mathbf{p})$ of the single particle excitations triggered by the rearrangement. In the limit $\varepsilon \to 0$, the width $\gamma(\varepsilon)$ of the states in the region of the Brillouin zone, where the dispersion of $\xi(\mathbf{p})$ has an ordinary order of value is found to depend on ε almost linearly in contrast to the Fermi-liquid-theory result $\gamma(\varepsilon) \sim \varepsilon^2$.

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In a number of articles [1 – 5] within the Landau – Migdal quasiparticle pattern, a new phase transition in strongly correlated Fermi systems, called fermion condensation, has been described. Its salient feature is in the appearance of the fermion condensate (FC), a group of degenerate single particle states whose energy at the temperature $T = 0$ coincides with the chemical potential μ. Owing to the degeneracy, quasiparticle occupation numbers $n(\mathbf{p}, T = 0)$ are no longer 1 or 0. They are determined by a variational condition

$$\delta E_0(n)/\delta n(\mathbf{p}) = \mu, \qquad \mathbf{p} \in \Omega, \tag{1}$$

where $E_0(n(\mathbf{p}))$ stands for the ground state energy, while Ω denotes the FC region whose boundaries are determined by eq. (1) itself. According to Landau, the l.h.s. of (1) is nothing but the quasiparticle energy $\varepsilon(\mathbf{p})$ and therefore this equation implies smearing the Fermi surface – its metamorphosis into a volume in three-dimensional systems or into a surface in two-dimensional ones. The flattening of the single particle electron spectra $\xi(\mathbf{p})$ has been experimentally observed in many quasi-two-dimensional superconductors (see e.g. [6, 7] and refs. therein). Observable single particle peaks have much larger widths γ than those calculated employing Fermi liquid (FL) theory (see e.g. [8, 9]) which is presumably triggered by the flattening of $\xi(\mathbf{p})$ near the Fermi surface.

Our purpose is to calculate the imaginary part of the mass operator $\Sigma(\mathbf{p}, \varepsilon)$ in the presence of a flat portion in the electron spectrum $\xi(\mathbf{p})$ due to the fermion condensation. In the following we consider a crystal with a cubic or square lattice assuming the FC to be positioned in the vicinity of the van Hove points [4, 5], the FC density $\rho_c \equiv \eta\rho$ being rather small. Dealing with the momentum dependence of $\mathrm{Im}\,\Sigma(\mathbf{p}, \varepsilon)$ two regions can be distinguished: i) the FC region Ω which is made up of "patches" adjacent to the van Hove points, and ii) the normal region where the dispersion has an ordinary order of value, i.e. $|\nabla\xi(\mathbf{p})| = p_F/M^*$ where M^* is the effective mass of the ordinary quasiparticle, independent of T. We neglect a possible variations of relevant quantities inside these regions and use the subscript c for the FC region and the subscript n, for the "normal" one: $\mathrm{Im}\,\Sigma(\mathbf{p} \notin \Omega) \equiv \mathrm{Im}\,\Sigma_n$.

Without loss of generality, we restrict ourselves to the decay of a particle, hence implying $\varepsilon > 0$ in all formulas. We focus our attention on the process of the generation of an additional particle-hole (p-h) pair and neglect contributions from collective modes, e.g. spin fluctuations, important for understanding of many properties of high-T_c superconductors [7]. In doing so, we assume that T exceeds the critical temperature T_c of the superfluid phase transition at which the p-p correlations make no difference and such contributions may be omitted. In this article, we deal with energies ε exceeding a characteristic FC energy $\xi_{FC} = \langle \xi(\mathbf{p}\in\Omega) \rangle$ which, in its turn, exceeds T. With these restrictions, the damping $\gamma(\mathbf{p}, \varepsilon > 0)$ is evaluated with the help of the formula [10]

$$\gamma(\mathbf{p}, \varepsilon > 0) = \frac{1}{2} \operatorname{Tr} \iint d\tau_2 d\tau_3 \int_0^{\varepsilon} \frac{d\omega}{2\pi} \int_0^{\omega} \frac{d\varepsilon_3}{2\pi} \left| \Gamma(\mathbf{p}, \varepsilon, \mathbf{p}_1, \varepsilon_1, \mathbf{p}_2, \varepsilon_2, \mathbf{p}_3, \varepsilon_3) \right|^2 \times$$

$$\times \operatorname{Im} G(\mathbf{p}_1, \varepsilon_1) \operatorname{Im} G(\mathbf{p}_2, \varepsilon_2) \operatorname{Im} G(\mathbf{p}_3, \varepsilon_3). \tag{2}$$

Here $\mathbf{p}_3 = \mathbf{q} + \mathbf{p}_2 \equiv \mathbf{p} - \mathbf{p}_1 + \mathbf{p}_2$; $\varepsilon_3 = \omega + \varepsilon_2 \equiv \varepsilon - \varepsilon_1 + \varepsilon_2$ and $d\tau = d^n p/(2\pi)^n$ where n is the dimensionality of the lattice. Since we are only interested in clarifying the energy dependence of quantities at issue we shall omit numerical factors in future calculations.

If FL theory were applicable to the systems with the FC, then the ratio $r(\varepsilon) \equiv |\gamma(\varepsilon)|/\varepsilon$ should be small and $\operatorname{Im} G(\mathbf{p}, \varepsilon)$ could be replaced by $-\pi\,\delta(\varepsilon - \xi(\mathbf{p}))$. However, as we shall see, the damping of single particle states in systems with the FC is drastically enhanced compared to predictions of FL theory. This motivates us to take into account an irregular part σ of $\operatorname{Re}\Sigma$ and write down $\operatorname{Im} G$ in the form

$$\operatorname{Im} G(\mathbf{p}, \varepsilon) = \frac{\gamma(\mathbf{p}, \varepsilon)}{(\varepsilon - \sigma(\mathbf{p}, \varepsilon) - \xi(\mathbf{p}))^2 + \gamma^2(\mathbf{p}, \varepsilon)}. \tag{3}$$

The connection between the quantities σ and γ stems from the usual dispersion relation for the mass operator which implies $\sigma(\varepsilon)/\gamma(\varepsilon) \sim \ln \varepsilon$ at $\varepsilon \to 0$ provided $r_c \sim 1$. For example, dealing with marginal Fermi liquid discussed in [11] one has $\gamma(\varepsilon \to 0) \sim \varepsilon, \sigma(\varepsilon \to 0) \sim$ $\sim \varepsilon \ln |\varepsilon|$.

Since the momentum dependence of quantities in the FC region is neglected the density $N_c(\varepsilon)$ of the FC states becomes

$$N_c(\varepsilon) = -\int_{\Omega} \operatorname{Im} G_c(\mathbf{p}, \varepsilon)\, d\tau = \eta\rho\, \frac{|\gamma_c(\varepsilon)|}{d_c^2(\varepsilon)}, \tag{4}$$

where $d_c^2(\varepsilon) = (\varepsilon - \sigma_c(\varepsilon) - \xi_{FC})^2 + \gamma_c^2(\varepsilon)$. In the weak damping limit, eq. (4) is reduced to $N_c(\varepsilon) = \eta\rho\,\delta(\varepsilon)$. In the strong damping limit $r_c > 1$, the contribution of the FC region $\varepsilon \simeq \xi_{FC}$, is of no importance and then d_c^2 in eq. (4) can be replaced by $d^2(\varepsilon) = (\varepsilon - \sigma_c(\varepsilon))^2 + \gamma_c^2(\varepsilon)$. To avoid unjustified complications we restrict ourselves to the momentum-independent zero Landau harmonic in the Legendre expansion of the interaction amplitude $\Gamma(\cos\chi)$, where χ is the angle between the momenta of incoming particles. In this approximation, we replace $\Gamma_{n_1 n_2}(\mathbf{p}_i, \varepsilon_i) \to \Gamma_{n_1 n_2}(\omega)\,(1 - \sigma_1\sigma_2)$, where the index n_1 (n_2) indicates the number of the FC states in the initial (final) state.

The damping γ is decomposed into a sum

$$\gamma(\mathbf{p}, \varepsilon) = \gamma_0(\mathbf{p}, \varepsilon) + \gamma_1(\mathbf{p}, \varepsilon) + \gamma_2(\mathbf{p}, \varepsilon), \tag{5}$$

where the term $\gamma_k (k = 0, 1, 2)$ accumulates contributions containing k final FC Green functions. It should be indicated that for the square or cubic lattice the decay of the FC particle into a final state with no less than two FC states unambiguously results in the appearance all three FC final states.

Since we ignore collective modes reserving the analysis of their role for a future article, overwhelming contributions to γ_0 come from the region of the momentum transfer $q \sim p_F$ quite far from the van Hove points. There the functions Γ_{00}, Γ_{01} are known to vary slowly enough and so do the p-h propagator $A_n(q, \omega)$ the dimensionless imaginary component of which is generally written as

$$b(\mathbf{q}, \omega) = \frac{1}{N_n(0)} \int\int_0^{\omega} \mathrm{Im}\, G(\mathbf{p}, \varepsilon)\, \mathrm{Im}\, G(\mathbf{p}-\mathbf{q}, \varepsilon-\omega)\, d\tau\, \frac{d\varepsilon}{\pi}, \tag{6}$$

where we introduced the density of states $N_n(0) \sim \rho M^*/\varepsilon_F^0$ as a normalization factor, ε_F^0 is the Fermi energy of noninteracting gas with the density ρ. All of them may be replaced by averaged values:

$$b_n(\omega) \sim -\omega\, N_n(0)/\rho, \qquad |A_n(\omega) N_n(0)| \sim 1, \qquad N_n^2(0)\, |\Gamma_{00}|^2 \sim N_n^2(0)\, |\Gamma_{01}|^2 \sim 1. \tag{7}$$

Upon substituting these values into eq. (2) the integrals cease to depend on $\mathbf{p}$, and we arrive at the usual result

$$\gamma_0(\varepsilon) = -\gamma_0\, (\varepsilon^2/\varepsilon_F^0), \tag{8}$$

where γ_0 is a dimensionless positive constant.

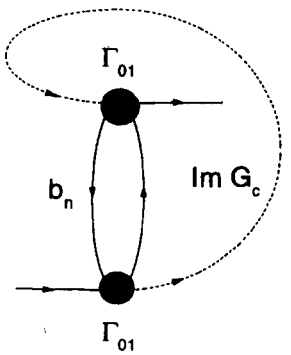

Fig.1. The graphical representation for the contribution γ_{1n} to the total damping

When evaluating the term $\gamma_{1n}(\varepsilon)$, we combine two ordinary Green functions entering eq. (2) into a p-h propagator A_n (see Fig.1) so that

$$\gamma_{1n}(\varepsilon) \sim N_n(0) \int_0^{\varepsilon} N_c(\varepsilon-\omega)\, |\Gamma_{01}|^2\, b_n(\omega)\, d\omega \sim \frac{1}{\rho} \int_0^{\varepsilon} N_c(\varepsilon_1)(\varepsilon_1 - \varepsilon)\, d\varepsilon_1. \tag{9}$$

The magnitude of this integral crucially depends on the value of $r_c(\varepsilon_1) = |\gamma_c(\varepsilon_1)|/\varepsilon_1$. If this ratio is relatively small then the integral (9) receives its overwhelming contributions (of order of 1) from the FC region where $\varepsilon_1 \simeq \xi_{FC}$ and $N_c(\varepsilon) \sim \delta(\varepsilon)$. Then after simple integration we are led to formulas inherent in marginal Fermi liquids [11]

$$\gamma_{1n}(\varepsilon) = -\gamma_{1n} \eta\, \varepsilon, \qquad \sigma_{1n}(\varepsilon) = \frac{2\gamma_{1n}\eta}{\pi}\, \varepsilon \ln \frac{|\varepsilon|}{\varepsilon_L}, \tag{10}$$

where γ_{1n} is a positive numerical constant of order of 1. We see that the magnitude of the linear in ε term is proportional to $\eta < 1$ and therefore it is quite small.

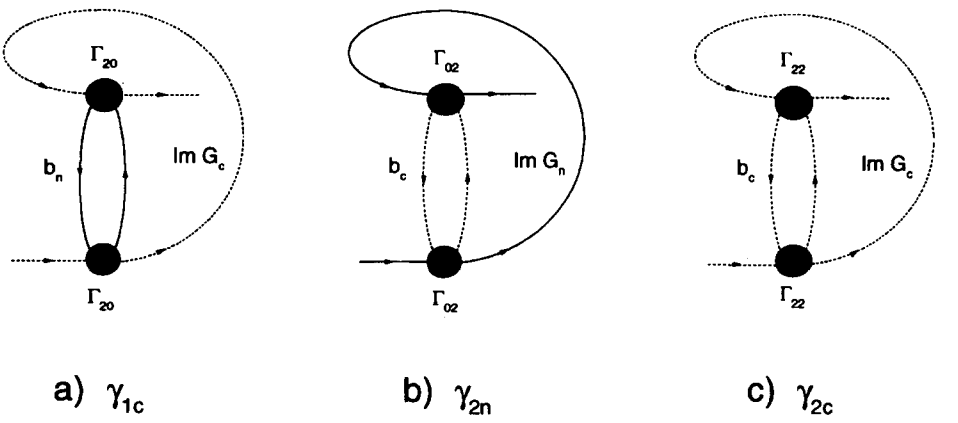

Fig.2. The same for γ_{1c} (a), γ_{2n} (b) and γ_{2c} (c)

In the opposite case of the strong damping with $r_c > 1$, the contribution of γ_{1n} to the damping γ is minor compared to that of γ_{2n}. When evaluating the latter quantity two final FC states come into play (see Fig.2b) and γ_{2n} is given by

$$\gamma_{2n}(\mathbf{p},\varepsilon) \sim -N_n(0) \int\limits_{\Omega'} \int\limits_0^{\varepsilon} \mathrm{Im}\, G_n(\mathbf{p}_1, \varepsilon-\omega)\, |\Gamma_{02}(\omega)|^2\, b_c(\omega)\, d\tau_1\, d\omega. \tag{11}$$

Here we introduced a special notation Ω' for the momentum integration region in (11), which is determined by the requirement that the momentum $\mathbf{p}_1$ is appropriate for arising the pair with momenta $\mathbf{p}_2$ and $(\mathbf{p} - \mathbf{p}_1 + \mathbf{p}_2)$, both in the FC region. We omit the insignificant term σ_n in $\mathrm{Im}\, G_n$ and assume the η value to be sufficient for satisfying inequalities $\varepsilon_{max} > |\gamma_n(\varepsilon)|$, where ε_{max} is the upper energy limit admissible for the normal states in the region Ω', that yields

$$\int\limits_{\Omega'} \mathrm{Im}\, G_n(\mathbf{p}-\mathbf{q}, \varepsilon)\, d\tau_{\mathbf{q}} \sim N_n(0) \int\limits_0^{\varepsilon_{max}} \frac{\gamma_n(\varepsilon)\, d\xi}{(\varepsilon-\xi)^2 + \gamma_n^2(\varepsilon)} \sim -N_n(0).$$

As a result, one finds

$$\gamma_{2n}(\varepsilon) \sim N_n^2(0) \int\limits_0^{\varepsilon} |\Gamma_{02}(\omega)|^2\, b_c(\omega)\, d\omega. \tag{12}$$

The interaction amplitude Γ_{02} is extracted from a set of algebraic equations

$$\begin{aligned} \Gamma_{02}(\omega) &= \mathcal{F}_{02} + \mathcal{F}_{02} A_c(\omega) \Gamma_{22}(\omega) + \mathcal{F}_{00} A_n(\omega) \Gamma_{02}(\omega), \\ \Gamma_{22}(\omega) &= \mathcal{F}_{22} + \mathcal{F}_{22} A_c(\omega) \Gamma_{22}(\omega) + \mathcal{F}_{20} A_n(\omega) \Gamma_{02}(\omega). \end{aligned} \tag{13}$$

This set is easily solved to arrive at

$$\Gamma_{02}(\omega) = \mathcal{F}_{02}/D(\omega), \qquad \Gamma_{22}(\omega) = [\mathcal{F}_{22}\Big(1 - \mathcal{F}_{00} A_n(\omega)\Big) + \mathcal{F}_{20}^2 A_n(\omega)]/D(\omega), \tag{14}$$

where $D(\omega) = \Big(1 - \mathcal{F}_{00}A_n(\omega)\Big)\Big(1 - \mathcal{F}_{22}A_c(\omega)\Big) - \mathcal{F}_{02}^2 A_n(\omega)A_c(\omega)$.

Omitting contributions of $\mathrm{Re}\, A_c(\omega)$ since they are insignificant [12] and replacing $A_n(\omega)$ by $-N_n(0)$, as before, one finds

$$|N_n(0)\Gamma_{02}(\omega)|^2 \sim \frac{h_{20}^2}{h_{00}^2 + b_c^2(\omega)}, \qquad |N_n(0)\Gamma_{22}(\omega)|^2 \sim \frac{1}{h_{00}^2 + b_c^2(\omega)}, \tag{15}$$

where notations $f_{ab} = \mathcal{F}_{ab}N_n(0)$ and $h_{20} = f_{20}/[f_{22}(1 + f_{00}) - f_{20}^2]$, $h_{00} = (1 + f_{00})/[(f_{22}(1 + f_{00}) - f_{20}^2]$ are introduced. With these results, eq. (12) becomes

$$\gamma_{2n}(\varepsilon) \sim h_{20}^2 \int_0^\varepsilon \frac{b_c(\omega)\, d\omega}{h_{00}^2 + b_c^2(\omega)}. \tag{16}$$

To take the next step one needs to evaluate the propagator $b_c(\omega)$ that, in its turn, requires the calculation of the damping γ_c of the FC particles which is made up of three terms (5) where $\gamma_0(\varepsilon)$ is given by eq. (8) while $\gamma_{1c}(\varepsilon)$ and $\gamma_{2c}(\varepsilon)$, depicted in Fig.2a, Fig.2c, are written as follows

$$\gamma_{1c}(\varepsilon) \sim -N_n(0) \int_\Omega \int_0^\varepsilon \mathrm{Im}\, G_c(\varepsilon-\omega)\, |\Gamma_{20}(\omega)|^2\, b_n(\omega)\, d\tau_1\, d\omega \sim -\frac{h_{20}^2}{\rho} \int_0^\varepsilon \frac{\omega\, N_c(\varepsilon-\omega)\, d\omega}{h_{00}^2 + b_c^2(\omega)},$$

$$\gamma_{2c}(\varepsilon) \sim -N_n(0) \int_\Omega \int_0^\varepsilon \mathrm{Im}\, G_c(\varepsilon-\omega)\, |\Gamma_{22}(\omega)|^2\, b_c(\omega)\, d\tau_1\, d\omega \sim \int_0^\varepsilon \frac{N_c(\varepsilon-\omega)\, b_c(\omega)\, d\omega}{N_n(0)\left[h_{00}^2 + b_c^2(\omega)\right]}. \tag{17}$$

Summing up various contributions to γ_c we are finally led to a nonlinear integral equation

$$\gamma_c(\varepsilon) = -\gamma_0 \frac{\varepsilon^2}{\varepsilon_F^0} + \int_0^\varepsilon \frac{\gamma_c(\varepsilon-\omega)\left[C_1 h_{20}^2 \eta\, \omega - C_2 \delta_F b_c(\omega)\right]}{d_c^2(\varepsilon-\omega)\left[h_{00}^2 + b_c^2(\omega)\right]}\, d\omega, \tag{18}$$

where C_1, C_2 are positive numerical constants of order of 1, $\delta_F = \eta\rho/N_n(0) = \eta\varepsilon_F^0 M/M^*$. The propagator $b_c(\omega)$ defined by eq. (6) reads

$$b_c(\omega) \sim \frac{1}{N_n(0)} \int_\Omega \int_0^\omega \mathrm{Im}\, G_c(\mathbf{p}, \varepsilon)\, \mathrm{Im}\, G_c(\mathbf{p}-\mathbf{q}, \varepsilon-\omega)\, d\tau\, d\varepsilon \sim \delta_F \int_0^\omega \frac{\gamma_c(\varepsilon)\, \gamma_c(\varepsilon-\omega)\, d\varepsilon}{d_c^2(\varepsilon)\, d_c^2(\varepsilon-\omega)}. \tag{19}$$

It is easily verified that the FL solution $\gamma_c(\varepsilon) \sim \varepsilon^2$, (the first term in eq. (18)) fails in the low energy limit $\varepsilon \to 0$. Indeed, if it were valid then the ratio $\gamma_c(\varepsilon)/d_c^2(\varepsilon)$ could be replaced by $-\pi\delta(\varepsilon)$ while the propagator b_c, by a constant. Upon inserting these results into the integral (18) and simple integration we would obtain a constant term which does not vanish at $\varepsilon \to 0$ as it must be since $\gamma_c(\varepsilon)$ is supposed to be proportional to ε^2.

To find a correct solution at $\varepsilon \to 0$ we employ the scaling approach assuming

$$\gamma_c(\varepsilon \to 0) = -\gamma_c\, \varepsilon^\nu\, |\ln \varepsilon|^\beta, \tag{20}$$

where γ_c is a positive constant while ν and β are critical indices calculated from the requirement of the cancellation between leading terms on the r.h.s. of eq. (18) and $\gamma_c(\varepsilon)$ itself. We shall see that at least one solution has the critical index $\nu < 1$. In this case, the ratio $r_c(\varepsilon \to 0) = |\gamma_c(\varepsilon)|/\varepsilon$ diverges and therefore one can retain in the denominators of eq. (19) only the term $\sigma_c(\varepsilon) \sim \gamma_c\, \varepsilon^\nu\, |\ln \varepsilon|^{\beta+1}$ together with the analogous term $\sigma_c(\varepsilon - \omega)$.

Upon substituting them into eq. (19) we are led to

$$b_c(\omega \to 0) = -\frac{\delta_F}{\gamma_c^2} \int_0^{\omega} \frac{d\varepsilon}{\varepsilon^{\nu}\,(\omega-\varepsilon)^{\nu}\,|\ln\varepsilon|^{\beta+2}|\ln(\omega-\varepsilon)|^{\beta+2}} \simeq \frac{\delta_F}{\gamma_c^2}\,\omega^{1-2\nu}\left|\ln\omega\right|^{-2\beta-4}. \quad (21)$$

Inserting this term into eq. (18) one finds that the term with the factor C_2 prevails on the r.h.s. of eq. (18). The requirement of its cancellation with $\gamma_c(\varepsilon)$ itself yields

$$\nu = 1/2, \qquad \beta = -3/2, \qquad \gamma_c^2 \simeq \delta_F/h_{00}. \quad (22)$$

Knowing $\gamma_c(\omega)$ and $b_c(\omega)$, the damping $\gamma_{2n}(\varepsilon)$ is straightforwardly evaluated from eq. (16). One obtains $\gamma_{2n}(\varepsilon) \sim \varepsilon/\ln\varepsilon$ with the numerical factor proportional to the ratio h_{20}^2/h_{00}. In principle, this ratio may be rather small and then the FL term proportional to ε^2 becomes significant, otherwise $\gamma_n(\varepsilon \to 0)$ is practically a linear in ε function with the prefactor of order of 1. Collecting all results one finally obtains

$$\gamma_c(\varepsilon \to 0) \sim -\sqrt{\frac{\delta_F}{h_{00}}\,\frac{\varepsilon}{|\ln\varepsilon|^3}}, \quad b_c(\omega \to 0) \sim \frac{h_{00}}{\ln\omega}, \quad \gamma_n(\varepsilon \to 0) \sim -\frac{h_{20}^2}{h_{00}}\,\frac{\varepsilon}{|\ln\varepsilon|} + O(\varepsilon^2). \quad (23)$$

It can be verified that this is the single solution of the problem. Thus the damping of normal excitations at low energies shows up marginal behavior with the factor of order of 1 while the damping of the FC states is enormously enhanced. This means that dealing with the damping of single particle states beyond the FC phase transition point FL theory fails even in the limit $\varepsilon \to 0$.

The integral equation (18) is also greatly facilitated at sufficiently high energies at which characteristic values of $r_c(\varepsilon)$ in the integrals standing on the l.h.s. of eq. (18) is fewer that 1. In this case, $N_c(\varepsilon)$ can be replaced by $\eta\rho\,\delta(\varepsilon)$ and we are left with an algebraic equation which after utilizing the relation (19) for b_c and some algebra takes the form

$$\gamma_c(\varepsilon)\left[1 - C_2\,\frac{\delta_F^2\,d_c^2(\varepsilon)}{\gamma_{22}^2\,d_c^4(\varepsilon) + \delta_F^2\,\gamma_c^2(\varepsilon)}\right] = -\gamma_0\,\frac{\varepsilon^2}{\varepsilon_F^0} - C_1\eta\,\varepsilon\,\frac{\gamma_{20}^2\,d_c^4(\varepsilon)}{\gamma_{22}^2\,d_c^4(\varepsilon) + \delta_F^2\,\gamma_c^2(\varepsilon)}. \quad (24)$$

In the weak damping limit $r_c \ll 1$, the C_2 term in the square brackets can be omitted and then the formula (10) is recovered. However, as we have seen above, the smallness of $r_c(\varepsilon)$ cannot always hold as ε goes down to zero. More precisely, the weak damping limit persists until the terms in the square brackets on the l.h.s. of eq. (24) start to cancel each other when ε goes down to 0. A rough estimate for the "watershed" ε_{cr} separating the weak damping limit from the strong one can be obtained by setting this sum to zero at the point $\gamma_c = 0$. As a result, one obtains

$$\varepsilon_{cr} \simeq \sqrt{C_2}\,(\delta_F/\gamma_{22}). \quad (25)$$

At lower energies, solutions (23) can be applied as the first approximation. At energies $\varepsilon > \varepsilon_{cr}$, the value of $b_c(\varepsilon) \sim \delta_F\gamma_c(\varepsilon)/\varepsilon^2 \sim \eta$ becomes small and, according to eq. (12), so does the derivative $d\gamma_{2n}(\varepsilon)/d\varepsilon$. Thus the energy dependence of the damping $\gamma_n(\varepsilon{>}\varepsilon_{cr})$ given primarily by the term (10) again turns out to be linear, the slope being relatively small. More precisely,

$$|\gamma_n(\varepsilon{>}\varepsilon_{cr})| = \gamma_{2n} + \gamma_1\eta\,\varepsilon + \gamma_0\,(\varepsilon^2/\varepsilon_F^0). \quad (26)$$

Upon substituting the propagator $b_c(\varepsilon)$ given by eq. (23) into eq. (16) and integrating over energy one finds $\gamma_{2n} \sim \delta_F$.

Thus the energy dependence of the damping $\gamma_n(\varepsilon)$ of the normal states in the systems with the FC turns out to be more complicated than in marginal FL: one linear regime with the slope of order of 1 transforms into another linear regime with the fewer slope at energies $\varepsilon \simeq \varepsilon_{cr}$. This result does not contradict to recent precise measurements of the linewidths of normal states in the compound Bi2212 [9].

In conclusion, we have analyzed the energy dependence of the damping $\gamma(\varepsilon)$ at some distance from the Fermi surface exceeding the characteristic energy ξ_{FC} of the FC. We have concentrated on crude features of this phenomenon without specific attention to numerical factors in quantities of our interest. We have demonstrated that dealing with the damping of single particle states in these systems there is no room for Landau theory: even in the limit $\varepsilon \to 0$ the ratio $r(\varepsilon) = |\gamma(\varepsilon)|/\varepsilon$ exceeds the unity. Our results for the damping $\gamma_n(\varepsilon)$, significantly differ from those obtained in theory of nearly antiferromagnetic Fermi liquid [13] and in the model of marginal FL considered in [11].

It becomes clear some ten years ago [11, 14] that the analysis of numerous experimental data on the base of the Luttinger liquid, characterized by a vanishing of the renormalization constant z, is an appropriate guide to understanding of properties of strongly correlated electron systems in solids including transport phenomena mostly dependent on the damping of the single particle degrees of freedom. At the same time, other properties of these systems including the existence of the Fermi surface itself are known to be successfully treated within Landau theory. What we have demonstrated in the article is that the fermion condensation model serves as a "bridge" connecting the two approaches.

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