

Excitation of rotation collective modes in a vortex lattice of clean type-II superconductors

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In a superclean limit the Magnus force on Abrikosov vortices is stronger than friction. Due to this, nondissipative force vortex segments rotate around pinning centers. Waves of such rotations under certain conditions are only weakly damped (not overdamped as is usually the case) and lead to resonances in the ac response. (The excitation of such waves by applied ac field near the surface is considered. Surface impedance, ac resistivity, and magnetic permeability are calculated using elasticity theory of the vortex lattice. [S0163-1829(99)00442-7])

I. INTRODUCTION

Abrikosov vortex dynamics in type-II superconductors under magnetic field is usually thought to be overdamped. Due to large vortex viscosity the displacement waves in vortex lattice do not propagate. In high- T_c superconductors the situation under certain conditions might be different. The dissipation during the vortex motion is at least to large extent due to excitation of quasiparticles inside the vortex core. At small temperatures this process is frozen and instead of usual Bardeen-Stephen friction force $\eta\mathbf{v}$ one has only a nondissipative Magnus force $\eta'\hat{\mathbf{z}}\times\mathbf{v}$ perpendicular to the vortex velocity, where z is the direction of external magnetic field. As evidence to the increasing role of the Magnus force is the famous Hall anomaly.¹ In a series of direct experiments² it was shown that in YBCO single crystals at low temperatures the Hall angle $\tan(\theta_H)\equiv\eta'/\eta$ diverges as T^{-1} and clearly exceeds 1 below 4 K reaching 2.5 at 3 K. This regime was termed by authors of Ref. 2 "superclean limit." Theoretically such a behavior was predicted in Ref. 3. In such a superclean regime vortex dynamics might be nonoverdamped and, for example, displacement waves in the vortex lattice can propagate. This type of phenomenon was used recently⁴ to explain the magnetoabsorption in BSCCO,⁵ although alternative explanations based on the Josephson plasma oscillations exist.⁶

In this paper we consider dynamics of vortices in "superclean" superconductors under applied ac field. Linear response to applied field (microwave impedance, low frequency complex resistivity, and permeability) and local field profiles are calculated. We find that in the superclean limit the system is not overdamped and point out to several resonance effects. Excitation of the nonoverdamped waves by applied ac field modify in an essential way the theory of the linear response developed by Brandt⁷ and Coffey and Clem.⁸

The basic physics of the vortex response in the superclean limit is very simple. Let us first consider the very small inductions case $B/\Phi_0 < \lambda^2$, where λ is the magnetic penetra-

tion depth and $\Phi_0\equiv hc/2e$. In this case we can neglect exponentially small interactions between vortices and consider single vortex dynamics. Assuming that the vortex is pinned (Fig. 1) we describe it by equation of motion for displacement \mathbf{u} :

$$m\ddot{\mathbf{u}} + \eta\dot{\mathbf{u}} + \eta'\mathbf{z}\times\dot{\mathbf{u}} + \alpha\mathbf{u} = 0. \tag{1}$$

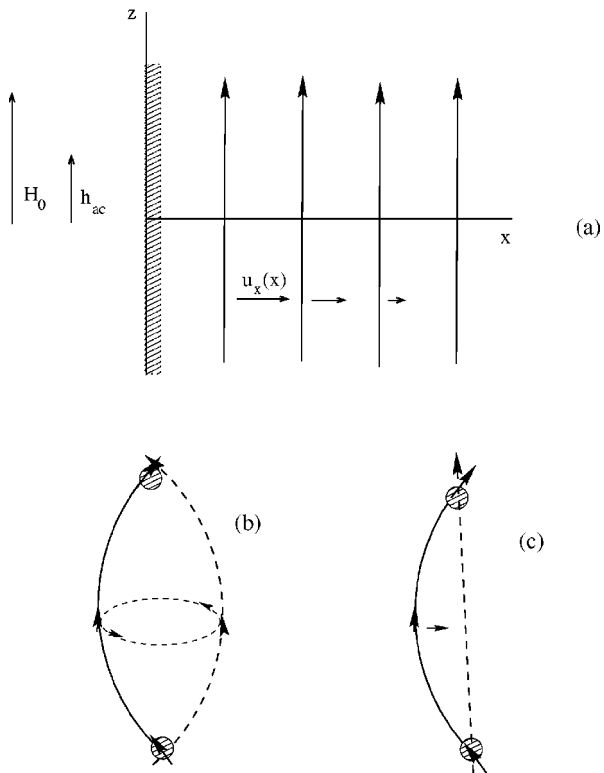


FIG. 1. (a) Positions and displacements of vortices caused by external Lorentz force. (b) Displacement of a vortex segment under influence of ac field in the superclean limit. (c) Displacement of a vortex segment under influence of ac field in the conventional overdamped case.

Here α is the Labousch parameter describing restoring pinning force in the x - y plane. As was noted many times⁹ the Magnus force is mathematically identical to ‘‘Lorentz’’ force on a charged particle when Φ_0 is considered as a ‘‘charge.’’ When the friction coefficient η and the dynamical vortex mass m (Ref. 10) are small one obtains (clockwise and counterclockwise) circular motion around the pinning centers [see Fig. 1(b)] with frequency

$$\omega_M \equiv \frac{\alpha}{\eta'}. \quad (2)$$

The order of magnitude is the same as the depinning frequency $\Omega_{\text{depin}} = \alpha/\eta$ which is of order 10-100 GHz.¹³ The friction causes damping of the rotations [see Fig. 1(c)]. To excite this mode resonantly one can apply the external Lorentz force due to ac current in a direction y perpendicular to the dc magnetic field and parallel to the samples surface, see Fig. 1(a). We assume $\mathbf{F}_{\text{ext}}(\mathbf{r}, t) = \hat{\mathbf{x}} F_{\text{ext}}(x) e^{i\omega t}$, namely, that the force is independent of y and z . In this case there is no variations in the z direction and the problem becomes two dimensional. When ω approaches ω_M one obtains resonance in the amplitude of vortex vibrations.

The system of coupled charged oscillators in magnetic field has been considered in various contexts in physics, e.g., in connection with ‘‘magnetophonons’’ in Wigner lattice formed out of electron gas.¹¹ This is very similar to the situation of interacting fluxons in the superclean limit for fields $H_{c1} \ll H \ll H_{c2}$. In this case the repulsion between vortices leading to formation of vortex lattice is logarithmic. The eigenfrequencies $\pm \omega_M$ become bands (the positive frequency is not necessarily the same as the negative when shear is present in the lattice, see below). This medium can be effectively described using the elasticity theory in harmonic approximation. The complex dispersion relation for these waves for arbitrary values of η, η', α , and B within London approximation are presented in Sec. II. These modes lead to number of experimentally accessible effects. In Sec. III we discuss the linear response to ac current flowing at the surface layer of the single crystal in direction perpendicular to that of magnetic field. We calculate the penetration depth λ_{ac} and related to it microwave impedance $Z(\omega)$, low frequency complex resistivity, and magnetic permeability. Compared to standard results^{7,8} we obtain additional large contribution to ac response near the resonant frequency due to excitation of weakly damped waves. In Sec. IV the modi-

fications due to anisotropy (important for applications to high- T_c materials) are made. In Sec. V we conclude by discussing possible destruction of vortex lattice by excitation of such wave and the corresponding shift of the melting line.

II. DISPERSION RELATION FOR WAVES IN VORTEX LATTICE

Neglecting vortex mass in Eq. (1) for single vortex dynamics one obtains the following periodic solution: $u_i(t) = e^{i\Omega_{\pm} t} u_i$, where

$$\Omega_{\pm} = \alpha \frac{i\eta \pm \eta'}{\eta^2 + \eta'^2}. \quad (3)$$

Contribution of interactions between vortices to the vortex dynamics can be taken into account within the harmonic approximation

$$\begin{aligned} \eta \dot{u}_i(\mathbf{R}^a) + \eta' \epsilon_{ij} \dot{u}_j(\mathbf{R}^a) + \alpha u_i(\mathbf{R}^a) \\ + \sum_b \Phi_{ij}(\mathbf{R}^a - \mathbf{R}^b) u_j(\mathbf{R}^b) = 0. \end{aligned} \quad (4)$$

Here Φ_{ij} is the dynamical matrix and \mathbf{R}^a are locations of vortices usually arranged in the lattice and ϵ_{ij} is the totally antisymmetric tensor. Since we are using elasticity theory the detailed nature of the vortex matter is not very important as long as correct elastic moduli are used and most of the considerations are valid in vortex liquid or glass. We will consider only external forces homogeneous in y and z directions, therefore the only nonzero component of momentum is $k_x \equiv k$. When external force is absent displacement vector for frequency Ω satisfies

$$\begin{pmatrix} (i\Omega \eta + \alpha) + c_{11}(k)k^2 & i\Omega \eta' \\ -i\Omega \eta' & (i\Omega \eta + \alpha) + c_{66}(k)k^2 \end{pmatrix} \begin{pmatrix} u_x(k) \\ u_y(k) \end{pmatrix} \\ \equiv A_{ij} u_j = 0, \quad (5)$$

where c_{11} and c_{66} are (possibly dispersive) elastic moduli of the vortex matter. In the London limit¹²

$$c_{11}(k) = \frac{B^2}{4\pi} \frac{1}{1 + \lambda^2 k^2} \equiv \frac{c_{11}}{1 + \lambda^2 k^2}, \quad c_{66} = \frac{B\Phi_0}{4(4\pi\lambda)^2}. \quad (6)$$

The eigenfrequencies Eq. (3) now become branches:

$$\Omega_{\pm}(k) = \frac{i\eta[2\alpha + k^2 \bar{c}]}{2(\eta^2 + \eta'^2)} \pm \frac{\sqrt{-\eta^2[2\alpha + k^2 \bar{c}]^2 + 4(\eta^2 + \eta'^2)[\alpha^2 + \alpha k^2 \bar{c} + k^4 c_{11}(k) c_{66}]}}{2(\eta^2 + \eta'^2)}, \quad (7)$$

where $\bar{c} \equiv c_{11}(k) + c_{66}$. In the superclean limit $\eta = 0$ one has

$$\Omega_{\pm}(k) = \pm \frac{\sqrt{\alpha^2 + \alpha k^2 \bar{c} + k^4 c_{11}(k) c_{66}}}{\eta'}. \quad (8)$$

The condition that there is a nonzero real part of $\Omega_{\pm}(k)$ (waves) is

$$\tan^2 \theta_H > \frac{k^4 (c_{11}(k) - c_{66})^2}{4(\alpha^2 + \alpha k^2 \bar{c} + k^4 c_{11}(k) c_{66})}. \quad (9)$$

Assuming that moduli have no dispersion, namely, $k\lambda < 1$ one obtains the following condition:

$$k^4[4c_{11}c_{66}\tan^2\theta_H - (c_{11} - c_{66})^2] + k^2[4\alpha\tan^2\theta_H\bar{c}] + 4\alpha^2\tan^2\theta_H > 0. \quad (10)$$

It is obviously satisfied for all k if $\tan\theta_H > (1/2)(c_{11} - c_{66})/\sqrt{c_{11}c_{66}}$. When this inequality does not hold only modes with

$$k^2 < \frac{2\alpha\tan\theta_H[\tan\theta_H(c_{11} + c_{66}) + (c_{11} - c_{66})/\cos\theta_H]}{(c_{11} - c_{66})^2 - 4\tan^2\theta_H c_{11}c_{66}} \equiv k_m^2 \quad (11)$$

have a nonzero real part of $\Omega_{\pm}(k)$. For $B \gg H_{c1}$ one has $c_{11} \gg c_{66}$ and the conditions Eqs. (10) and (11) simplify into $\tan\theta_H > \frac{1}{2}\sqrt{c_{11}/c_{66}} = \sqrt{\ln\kappa(B/H_{c1})}$ and $k^2 < 2\alpha\tan\theta_H(1 + \sin\theta_H)/c_{11}\cos\theta_H \equiv k_m^2$.

A stronger condition that $\text{Re}\Omega/\text{Im}\Omega = \Gamma > 1$ can be satisfied only when the superconductor is sufficiently ‘‘clean’’:

$$\tan\theta_H > \Gamma. \quad (12)$$

Moreover it will be satisfied for all k in the ‘‘superclean’’ limit:

$$\cos\theta_H < \frac{2}{\sqrt{\Gamma^2 + 1}} \frac{\sqrt{c_{11}c_{66}}}{c_{11} + c_{66}}. \quad (13)$$

This condition is very restrictive. More relevant case is when only some of the modes $k < k_m(\Gamma)$ have the imaginary part smaller than the real part. The maximal momentum for a given ration Γ is

$$k_m^2(\Gamma) \equiv \frac{2\alpha\Delta}{c_{11} - c_{66} - \Delta(c_{11} + c_{66})} \simeq \frac{2\alpha\Delta}{c_{11}(1 - \Delta)} \quad (14)$$

where $\Delta^2 \equiv 1 - (\Gamma^2 + 1)\cos^2\theta_H$ [always positive due to Eq. (12)].

Polarization of the waves [which follows from Eq. (5)] is the following:

$$\frac{u_y(k)}{u_x(k)} = -\tan^{-1}\theta_H + i\frac{\alpha + c_{11}k^2}{\Omega_{\pm}(k)\eta'}. \quad (15)$$

The fact that the ratio is imaginary means that vortices move on elliptic trajectories.

III. LINEAR RESPONSE UNDER APPLIED ac FIELD

In this section we consider the pinned vortex system response to surface ac current caused by alternating field $h_{ac}e^{i\omega t}$ in direction parallel to dc field H and to the surface of the superconducting half space, see Fig. 1(a). The linear response for such geometry for the case $\eta' = 0$ was considered by Brandt⁷ and Coffey and Clem⁸ also taking into account pinning, viscosity and creep. Since we are interested mostly in the low temperature regime flux creep can be neglected while the Magnus force term is important (creep can be taken into account in a similar manner as in (Refs. 7,8)). When one performs similar calculation for $\eta' > 0$ new resonant phenomena are readily seen. We impose proper bound-

ary conditions using the ‘‘bulk concept’’ methods of Ref. 7 which allows to refer the problem to an equivalent problem in whole space.

The external force is

$$F_{\text{ext}}(x, t) = \frac{Bh_{ac}}{4\pi\lambda} e^{-|x|/\lambda} e^{i\omega t}, \quad (16)$$

$$F_{\text{ext}}(k, \omega) = \frac{Bh_{ac}}{2\pi(1 + \lambda^2 k^2)}. \quad (17)$$

The displacement in momentum space is obtained from Eq. (5) with external force

$$u_x(k, \omega) = \frac{Bh_{ac}[(i\omega\eta + \alpha) + c_{66}k^2]}{2\pi(1 + \lambda^2 k^2)D(k, \omega)}, \quad (18)$$

$$u_y(k, \omega) = \frac{iBh_{ac}\omega\eta'}{2\pi(1 + \lambda^2 k^2)D(k, \omega)}, \quad (19)$$

where $D(k, \omega)$ is the determinant of matrix A :

$$D(k, \omega) = [i\omega\eta + \alpha + c_{11}(k)^2][i\omega\eta + \alpha + c_{66}k^2] - \omega^2\eta'^2. \quad (20)$$

A. Small shear modulus approximation

Very often both c_{66} and k are ‘‘small.’’ If $c_{66}k^2$ is small compared to $(\omega\eta + \alpha)$ one readily reexpresses the result in the form

$$u_x(k, \omega) = \frac{Bh_{ac}}{2\pi\alpha(\omega)[1 + k^2(c_{11}/\alpha + \lambda^2)]} = \frac{2h_{ac}\lambda_C^2(\omega)}{B[1 + k^2\lambda_{ac}^2(\omega)]}, \quad (21)$$

$$u_y(k, \omega) = \frac{i\omega\eta'}{i\omega\eta + \alpha} u_x(k, \omega), \quad (22)$$

where $\alpha(\omega) \equiv i\omega\eta + \alpha - \omega^2\eta'^2/(i\omega\eta + \alpha)$ and modified Campbell penetration depth $\lambda_C^2(\omega)$:

$$\begin{aligned} \lambda_C^2(\omega) &\equiv \frac{c_{11}}{\alpha(\omega)} = \frac{B^2(i\omega\eta + \alpha)}{4\pi[(i\omega\eta + \alpha)^2 - \omega^2\eta'^2]} \\ &= \frac{B^2(i\omega + \tan\theta_H\omega_M)}{4\pi\eta[(i\omega + \tan\theta_H\omega_M)^2 - \omega^2\tan^2\theta_H]} \\ &= -\frac{B^2(i\omega + \tan\theta_H\omega_M)\cos^2\theta_H}{4\pi\eta[\omega - \Omega_+(0)][\omega - \Omega_+(0)]}. \end{aligned} \quad (23)$$

The frequency dependent complex ac penetration depth was introduced:

$$\lambda_{ac}^2(\omega) \equiv \lambda^2 + \lambda_C^2(\omega). \quad (24)$$

As is in the usual case^{7,8} $\eta' = 0$ this quantity determines both the surface impedance

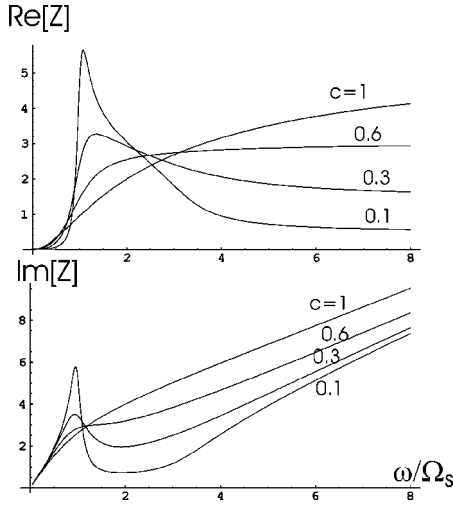


FIG. 2. Frequency dependence of the surface impedance (real and imaginary parts) for $\cos \theta_H = 0.1, 0.3, 0.6, 1$. Magnetic induction $B = 10 H_{c1}$.

$$Z_s(\omega) = \frac{4\pi i}{c^2} \omega \lambda_{ac}(\omega). \quad (25)$$

and the ac resistivity $\rho_{ac}(\omega) \equiv E(x)/J(x) = (4\pi i/c^2)\omega \lambda_{ac}^2(\omega)$. These two quantities exhibit resonance in the clean limit. In Fig. 2 real and imaginary parts of surface impedance for various values of $\cos \theta_H$ and $b \equiv B/H_{c1} = 10$ are shown. When $k^2 c_{66}$ is not negligible but still small compared to α one can correct perturbatively the expression for λ_{ac} :

$$\Delta \lambda_{ac}^2 = \frac{c_{66} \omega^2 \eta'^2}{(\alpha + i\omega\eta)[(\alpha + i\omega\eta)^2 - \omega^2 \eta'^2]}. \quad (26)$$

The dependence of the resonance peak near $c_{66} = 0$ (the melting of the vortex lattice) on c_{66} is quite regular.

B. General case

When the $c_{66} k^2$ term in Eqs. (18)–(20) is not small one can, following Ref. 7, calculate ac electric and magnetic field components caused by vortex displacements. One obtains the following expression for this part of the ac magnetic field:

$$B_1(x, \omega) = \frac{2}{b} \left\{ \frac{r_B(-1)}{[1+k_1^2(\omega)][1+k_2^2(\omega)]} + \frac{r_B(k_1)}{[1+k_1^2(\omega)][k_1^2(\omega)-k_2^2(\omega)]} + \frac{r_B(k_2)}{[1+k_2^2(\omega)][k_1^2(\omega)-k_2^2(\omega)]} \right\}, \quad (27)$$

where poles k_i (zeroes of determinant) and corresponding residua r_i are

$$k_{1,2} = \frac{B \pm \sqrt{B^2 - AC}}{A},$$

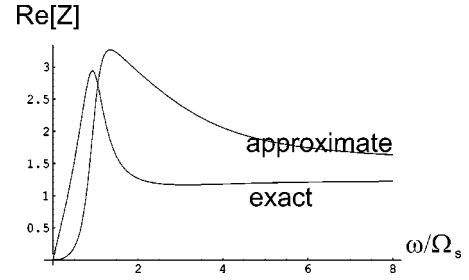


FIG. 3. Comparison of the exact expression for real part of surface impedance (accounting for the effect of shear modulus c_{66}) with an approximate one given by much simpler Eq. (26) for $\cos \theta_H = 0.3$, $B = 10 H_{c1}$.

$$A \equiv b - \alpha - i\omega \cos \theta_H,$$

$$B \equiv \frac{1}{2} [4\omega^2 - 4\alpha^2 - 8i\omega\alpha \cos \theta_H + (a + i\omega \cos \theta_H)(4b + 1)],$$

$$C \equiv -4(\omega^2 - \alpha^2 - 2i\omega\alpha \cos \theta_H)$$

$$r_B(q) \equiv (q/4 - \alpha - i\omega \cos \theta_H) \exp(iqx). \quad (28)$$

Here we use “natural” units. Unit of length is λ , unit of magnetic field is H_{c1} : $b \equiv B/H_{c1}$. Similarly the electric field part caused by the vortex displacements is

$$E_1(x, \omega) = \frac{2}{b} \left\{ \frac{r_E(-1)}{[1+k_1^2(\omega)][1+k_2^2(\omega)]} + \frac{r_E(k_1)}{[1+k_1^2(\omega)][k_1^2(\omega)-k_2^2(\omega)]} + \frac{r_E(k_2)}{[1+k_2^2(\omega)][k_1^2(\omega)-k_2^2(\omega)]} \right\}, \quad (29)$$

where now the residua are

$$r_E(q) \equiv \frac{\omega(q/4 - \alpha - i\omega \cos \theta_H)}{q} \exp(iqx). \quad (30)$$

This should be superimposed with the Meissner field part. Generally the microwave impedance is given by

$$Z(\omega) = \frac{4\pi E_1(x=0)}{c B_1(x=0)}. \quad (31)$$

As an example we compare the exact expression for the real part with the approximate one given by Eq. (25) in Fig. 3 (for $b = 10, \cos \theta_H = 0.3$). Generally resonance becomes sharper for larger shear moduli.

Similar resonance effects can be seen also in the frequency dependences of ac resistivity and magnetic permeability of a slab of thickness d . The complex ac resistivity is

$$\rho_{ac}(\omega)|_{x=0} = E_1(x=0)/J(x=0), \quad (32)$$

where the current density on the surface is $J(x=0) = (c/4\pi)(\partial/\partial x)B_1(x)|_{x=0}$. The magnetic permeability is calculated from

$$\mu = \langle B_1(x)/h_{ac} \rangle, \quad (33)$$

where $\langle \dots \rangle$ denoted average over the sample.

IV. EFFECT OF ANISOTROPY

The ac experiments are rarely carried out on relatively isotropic low- T_c materials such as Nb. Usually one would like to study highly anisotropic materials such as high- T_c superconductors. The results of the previous Secs. II and III can be easily extended to the case of anisotropic three-dimensional (3D) superconductor. We will consider the case of the uniaxial anisotropic superconductor (such as YBCO) with the c axis perpendicular to the sample surface and vortices parallel to either a or b directions considered equivalent for simplicity, see Fig. 1(a). The anisotropy is characterized by the coefficient $\Gamma \equiv \sqrt{m_c/m_{ab}} > 1$. In this case one should account for anisotropy of the parameters characterizing vortex matter elasticity, pinning force, and viscosity for vortex displacements along and perpendicular to the sample surface (our y and x directions, respectively).

Various properties of the mixed state in anisotropic superconductors were extensively studied in a number of works (see e.g., reviews 12,15 and references therein). For our purposes we make use the appropriate expressions for converting scalar quantities appearing in the isotropic case into generally tensorial quantities. Using the anisotropy coefficient Γ we relate them to an equivalent isotropic superconductor (for which all the quantities will be marked with subscript ‘‘0’’). Different vortex viscosities along the x and y axes are^{15,16}

$$\eta_x = \Gamma \eta_0, \quad \eta_y = \Gamma^{-1} \eta_0. \quad (34)$$

As far as Labusch parameter is concerned we assume that vortex pinning is mainly provided by pointlike pins separated by distances of order L in ab plane. Thus the anisotropy of the Labusch parameter $\alpha_{x,y}$ is determined by the anisotropy of the vortex line tension $P_{x,y}$ for displacements of a vortex along the c axis and along the ab plane.^{12,17} Components of Labusch parameters become

$$\alpha_x = \Gamma \alpha_0, \quad \alpha_y = \Gamma^{-1} \alpha_0, \quad (35)$$

where $\alpha_0 = P_0/L^2$. Changes in elastic properties of the vortex lattices due to anisotropy are as follows. While the compression modulus c_{11} remains unchanged in the small k limit (nondispersive part), the shear modulus c_{66} for displacements along the y axis is significantly reduced.^{15,18–20}

$$c_{66}^{\parallel} = \frac{1}{\Gamma^3} c_{66}^{(0)} = \frac{B\Phi_0}{64\pi^2\lambda^2\Gamma^3}. \quad (36)$$

The off diagonal components of viscosity tensor η_{xy} should obey the Onsager principle, thus since it is a phenomenological parameter at the present time, we write as before

$$\eta_{xy} = -\eta_{yx} = \eta'. \quad (37)$$

After these modifications the matrix A_{ij} in Eq. (5) takes the form

$$A_{ij} = \begin{pmatrix} (i\Omega \eta_x + \alpha_x) + c_{11}k^2 & i\Omega \eta' \\ -i\Omega \eta' & (i\Omega \eta_y + \alpha_y) + c_{66}^{\parallel}k^2 \end{pmatrix} = \begin{pmatrix} \Gamma(i\Omega \eta_0 + \alpha_0) + \frac{c_{11}}{\Gamma}k^2 & i\Omega \eta' \\ -i\Omega \eta' & \frac{1}{\Gamma} \left(i\Omega \eta_0 + \alpha_0 + \frac{1}{\Gamma^2} c_{66}^{(0)}k^2 \right) \end{pmatrix}. \quad (38)$$

From Eq. (38) one can see that eigenfrequencies determined by the condition $\det A_{ij} = 0$ are still described by Eq. (7) with simple replacements $c_{11} \rightarrow c_{11}/\Gamma$ and $c_{66} \rightarrow c_{66}^{\parallel}/\Gamma^2$. This means that elasticity now plays a lesser role in governing the vortex oscillations compared to the pinning and viscosity's role which has been enhanced. Moreover the role of shear as compared to compression becomes negligible and the approximation made in Sec. IIIA becomes better. Expression (23) for the Campbell penetration depth derived in Sec. IIIA after replacement $c_{11} \rightarrow c_{11}/\Gamma$ becomes

$$\lambda_C^2 = \frac{c_{11}}{\Gamma \alpha_0(\omega)}, \quad (39)$$

where $\alpha_0(\omega)$ corresponds to isotropic superconductor. The reduction in the Campbell penetration depth leads to corresponding modifications in the ac penetration depth λ_{ac} Eq. (24) in turn influencing the surface impedance and the ac resistivity. The position of the resonance peak however is un-

changed comparatively to equivalent isotropic superconductor (although the polarization of the displacement wave changes).

V. DISCUSSION

In this paper we determined conditions under which a nonoverdamped ‘‘rotation’’ (around pinning centers) waves exist in clean type-II superconductors. There are clear indications that these conditions can be met in untwinned YBCO single crystals.² It is even possible that these conditions can be met in some low- T_c materials such as superclean Nb.¹⁴ Excitation of such waves by applied ac field near the surface is considered. The simplest realistic geometry is the superconducting half space with the dc magnetic field creating vortices parallel to the surface. We considered the direction of the surface ac field parallel to the dc magnetic field. In this case linear response characteristics such as surface imped-

ance, ac resistivity and magnetic permeability were calculated using the elasticity theory of the vortex lattice. The most pronounced effect of the rotation waves is resonance at characteristic frequency of order $\Omega_s = \alpha/\eta'$. It is comparable or larger than the depinning frequency $\Omega_{\text{depin}} = \alpha/\eta$ which is of order 10–100 GHz.¹³

In the region of resonance amplitude of vortex displacements become quite large and nonlinear effects might show up. Qualitatively one expects the following. Vortex lattice can be at least locally destroyed. Therefore Larkin domains will be smaller leading to increase in critical current and the

melting line on the B - T phase diagram shifts to lower values of the magnetic field.

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