INFLUENCE OF THE AC CURRENT ON THE NONLINEAR DC RESISTIVE RESPONSE IN A TILTED WASHBOARD PINNING POTENTIAL

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The influence of the *ac* current on the longitudinal and transverse (with respect to the *dc c*urrent direction) even and odd (with respect to the magnetic field reversal) nonlinear anisotropic magnetoresistivities $\rho^{\pm}_{\|,\perp}$ of a type-II superconductor is considered. The current and frequency dependence of the number and position of Shapiro-like steps in the current-voltage characteristic is calculated and analyzed for the transverse geometry at low temperatures. A simple physical picture of the vortex motion in a washboard periodic pinning potential, tilted due to the presence of the dimensionless *dc* driving force, is formulated.

INTRODUCTION

It is well known that the mixed-state resistive properties of type-II superconductors are determined by the dynamics of vortices which again are strongly influenced by the distribution of pinning sites [1]. In the simplest case this distribution is assumed to be periodic in one dimension and can be described by using a washboard-like planar pinning potential (PPP). The temperature-dependent *dc* current uniaxial pinning anisotropy, provoked by such a washboard PPP has been extensively studied recently, both theoretically [2-4] and experimentally [5-9].

As the pinning force in a PPP is directed perpendicular to the washboard channels of the PPP [1], the vortices generally tend to move along these channels. Such a *guided* motion of vortices in the presence of the Hall effect produces an anisotropic transport behaviour for which the even (+) and odd (–) (with respect to the magnetic field reversal), longitudinal ($\|$) and transverse ($\perp$) (with respect to the current density vector **j**) nonlinear *dc* magnetoresistivities $\rho^{\pm}_{\|,\perp}$ depend substantially on the angle α between **j** and the direction of the PPP channels ("guiding direction").

The *dc*-current nonlinear guiding problem was exactly solved recently for the washboard PPP within the framework of a two-dimensional (2*D*) single-vortex stochastic model of anisotropic pinning based on the Fokker-Planck equation and rather simple formulas were derived for the *dc* magnetoresistivities $\rho^{\pm}_{\|,\perp}$ [2,3]. Thereupon the (*dc*+*ac*)-current nonlinear guiding problem was exactly solved in the framework of the Langevin equation [10] in terms of a matrix continued fraction [11] without recourse to the Fokker-Planck approach.

As a result, two groups of new findings were obtained in Ref. [10]. First, for the 2*D dc*–problem, previously solved in Refs. [2, 3], the influence of an *ac* current on the anisotropic $\rho^{dc\pm}_{\|,\perp}$ *dc*-response was calculated and analyzed. Second, for an *ac* current at a frequency ω under *dc* bias the 2*D* nonlinear time-dependent stationary $\rho^{ac\pm}_{\|,\perp}$ *ac*-response as function of frequency ω in terms of the nonlinear impedance tensor $\hat{Z}$ and the nonlinear *ac* response at ω-harmonics were studied.

In the present paper I discuss the influence of the *ac* current on the overall shape and appearance of Shapiro-like steps on the *dc* current-voltage characteristics (CVC) pointed out in Ref. [10] and propose a simple physical picture of the vortex motion in a washboard periodic pinning potential, tilted due to the presence of the dimensionless *dc* driving force.

MAIN RESULTS

The Langevin equation for a vortex moving with velocity $\mathbf{v}$ in a magnetic field $\mathbf{B} = \mathbf{n}B$ ($B \equiv |\mathbf{B}|$, $\mathbf{n} = n\mathbf{z}$, $\mathbf{z}$ is the unit vector in the z-direction and $n = \pm 1$) has the form

$$\eta\mathbf{v} + n\alpha_H \mathbf{v} \times \mathbf{z} = \mathbf{F}_L + \mathbf{F}_p + \mathbf{F}_{th}, \qquad (1)$$

where $\mathbf{F}_L = n(\Phi_0/c)\, \mathbf{j} \times \mathbf{z}$ is the Lorentz force, Φ_0 is the magnetic flux quantum, c is the speed of light, $\mathbf{j} \equiv \mathbf{j}(t) = \mathbf{j}^{dc} + \mathbf{j}^{ac} \cos \omega t$, where $\mathbf{j}^{dc}$ and $\mathbf{j}^{ac}$ are the *dc* and *ac* current density amplitudes and ω is the angular frequency, $\mathbf{F}_{th}$ is the thermal fluctuation force represented by a Gaussian white noise with zero mean value [2,3], η is the vortex viscosity, α_H is the Hall constant, $\mathbf{F}_p = -\nabla U_p(x)$ is the anisotropic pinning force and one assumes, as in [2,10], an a-periodic pinning potential of the form $U_p(x) = (U_p/2)(1\text{-}\cos kx)$ where $k = 2\pi/a$.

The main quantity of physical interest in our problem is the average electric field, induced by the moving vortex system, which is given by

$$\langle \mathbf{E} \rangle = (n/c)\mathbf{B} \times \langle \mathbf{v} \rangle = n(B/c)(-\langle v_y \rangle \mathbf{x} + \langle v_x \rangle \mathbf{y}), \qquad (2)$$

where $\mathbf{x}$ and $\mathbf{y}$ are the unit vectors in the x and y directions, respectively. As follows from Eq. (1) $\langle v_y \rangle = F_{Ly}/\eta + \delta \langle v_x \rangle$, where $\delta \equiv n\varepsilon, \varepsilon \equiv \alpha_H/\eta$, and so for determination of $\langle \mathbf{E} \rangle$ from Eq. (2) it is sufficient to calculate $\langle v_x \rangle$ from Eq. (1). This calculation gives

$$\langle v_x \rangle (t) = (\Phi_0 j_c / c\eta D)[j^{dc} + j^{ac} \cos \omega t - \langle \sin \mathsf{x} \rangle (t)], \qquad (3)$$

where $D \equiv 1 + \delta^2, j^{dc} \equiv n(j_y^{dc} + \delta j_x^{dc})/j_c, j^{ac} \equiv n(j_y^{ac} + \delta j_x^{ac})/j_c$, with $j_c \equiv cF_p/\Phi_0$, $j_y = j\cos\alpha,\ \ j_x = j\sin\alpha,\ \ \langle \sin \mathsf{x} \rangle (t) = (i/2)[\langle r \rangle (t) - \langle r^{-1} \rangle (t)],\ \ r(t) = \exp[-i\mathsf{x}(t)]$, and $\mathsf{x} = kx$.

Looking for only the stationary *ac* response, which is independent of the initial condition, one may seek all the $\langle r^m \rangle (t)$ in the form [11]

$$\langle r^m \rangle (t) = \sum_{k=-\infty}^{\infty} F_k^m(\omega) e^{ik\omega t}. \qquad (4)$$

On substituting Eq. (4) into Eq. (3) one obtains recurrence equations for the ($j^{ac}, j^{dc}, \Omega, g$)-dependent Fourier amplitudes F_k^m, where $g = (U_p/2T)$, T is the temperature in energy units, $\Omega = \omega \hat{\tau}$ is the dimensionless frequency and $\hat{\tau} \equiv \eta D/kF_p$ is the relaxation time. From the solution of these recurrence equations in terms of matrix continued fraction, as detailed in Refs. [10, 11], one can find the dimensionless average pinning force $\langle F_{px}\rangle(t)$ which is the main anisotropic nonlinear component of the theory, due to the α-dependence on the *ac* and *dc* current input,

$$\left\langle \hat{F}_{px}\right\rangle(t) = -\langle \sin \mathrm{X}\rangle(t) = \sum_{k=0}^{\infty} \mathrm{Im}(\psi_k e^{ik\omega t}) \tag{5}$$

where $\psi_0 \equiv F_0^1(\omega)$ and for $k \geq 1$ one has $\psi_k \equiv F_k^1(\omega) - F_k^{-1}(\omega)$. From Eq. (5) one can decompose the average pinning force $\left\langle \hat{F}_{px}\right\rangle(t)$ into three components

$$\left\langle \hat{F}_{px}\right\rangle(t) = \left\langle \hat{F}_{px}\right\rangle_0^{\omega} + \left\langle \hat{F}_{px}\right\rangle_{t1} + \left\langle \hat{F}_{px}\right\rangle_t^{k>1}. \tag{6}$$

In Eq. (6) $\left\langle \hat{F}_{px}\right\rangle_0^{\omega} \equiv -\langle \sin \mathrm{X}\rangle_0^{\omega} = \mathrm{Im}\,\psi_0$ is the time independent (but frequency dependent) *static* average pinning force; $\left\langle \hat{F}_{px}\right\rangle_{t1} \equiv -\langle \sin \mathrm{X}\rangle_{t1} = \mathrm{Im}(\psi_1 e^{i\omega t})$ is the time-dependent *dynamic* average pinning force with frequency ω of the *ac* current input; $\left\langle \hat{F}_{px}\right\rangle_t^{k>1} \equiv -\langle \sin \mathrm{X}\rangle_t^{k>1} = \mathrm{Im}(\psi_k e^{ik\omega t})$ describes the contribution of the harmonics with $k > 1$ into the dynamic average pinning force.

In order to derive the ω-dependent nonlinear *dc* resistivity and conductivity tensors I first express the time independent part of $\langle E_y\rangle(t)$ and $\langle E_x\rangle(t)$ as (see Eq. (2))

$$\begin{aligned} \langle E_y\rangle_0^{\omega} &= (\rho_f / D)\nu_0^{\omega}(j_y^{dc} + \delta j_x^{dc}), \\ \langle E_x\rangle_0^{\omega} &= (\rho_f / D)[j_x^{dc}(1+\delta^2(1-\nu_0^{\omega})) - \delta j_y^{dc}]. \end{aligned} \tag{7}$$

Here $\rho_f \equiv B\Phi_0/\eta c^2$ is the flux-flow resistivity, $\nu_0^{\omega} \equiv 1 - \langle \sin \mathrm{X}\rangle_0^{\omega} / j^{dc} = 1 + \left\langle \hat{F}_{px}\right\rangle_0^{\omega} / j^{dc}$ where ν_0^{ω} at $j^{ac} = 0$ can be considered as the probability of the vortex hopping over the pinning potential barrier under the influence of the dimensionless generalized moving force $\hat{F}_{Lx}^{dc} = j^{dc}$ in *x* direction [3].

From Eqs. (7) I find the magnetoresistivity tensor for the *dc* nonlinear law $\langle \mathbf{E}_0^{\omega}(\omega)\rangle = \hat{\rho}_0^{\omega}\,\mathbf{j}^{dc}$ as

$$\hat{\rho}_0^{\omega} = \begin{pmatrix} \rho_{xx}^{dc} & \rho_{xy}^{dc} \\ \rho_{yx}^{dc} & \rho_{yy}^{dc} \end{pmatrix} = \frac{\rho_f}{D}\begin{pmatrix} D-\delta^2\nu_0^{\omega} & -\delta\nu_0^{\omega} \\ \delta\nu_0^{\omega} & \nu_0^{\omega} \end{pmatrix}. \tag{8}$$

The *dc* conductivity tensor $\hat{\sigma}_0^{\omega}$, which is the inverse tensor to $\hat{\rho}_0^{\omega}$, has the form

$$\hat{\sigma}_0^{\omega} = \begin{pmatrix} \sigma_{xx}^{dc} & \sigma_{xy}^{dc} \\ \sigma_{yx}^{dc} & \sigma_{yy}^{dc} \end{pmatrix} = \frac{1}{\rho_f}\begin{pmatrix} 1 & \delta \\ -\delta & (D/\nu_0^{\omega})-\delta^2 \end{pmatrix}. \tag{9}$$

As it is seen from Eqs. (8) and (9), the off-diagonal components of the $\hat{\rho}_0^{\omega}$ and $\hat{\sigma}_0^{\omega}$ tensors satisfy the Onsager relation ($\rho_{xy} = -\rho_{yx}$ in the general nonlinear case and $\sigma_{xy} = -\sigma_{yx}$).

The experimentally observable longitudinal and transverse *dc* magnetoresistivities $\rho_{\parallel}^{dc} = E_{\parallel}^{dc} / j$ and $\rho_{\perp}^{dc} = E_{\perp}^{dc} / j$ (where j^d is the *dc* current density $(j^d)^2 = (j_x^{dc})^2 + (j_y^{dc})^2$ have the form

$$\begin{aligned} \rho_{\parallel}^{dc} &= (\rho_f / D)[(D - \delta^2 \nu_0^{\omega})\sin^2\alpha + \nu_0^{\omega}\cos^2\alpha], \\ \rho_{\perp}^{dc} &= (\rho_f / D)[\delta \nu_0^{\omega} - D(1-\nu_0^{\omega})\sin\alpha\cos\alpha], \end{aligned} \tag{10}$$

where in order to separate the even and odd components of $\rho_{\parallel,\perp}^{dc}$ one should use $\nu_0^{\omega\pm}(n) = [\nu_0^{\omega}(n) \pm \nu_0^{\omega}(-n)]/2$ which are the even and odd components (relative to the magnetic field inversion) of the function $\nu_0^{\omega}(n)$ (compare with Eqs. (13)-(14) in Ref. [4]).

DISCUSSION

Below I present a graphical analysis of the Ω-dependent *dc* nonlinear response calculated in the transverse ($\alpha = 0$) T-geometry (the current density vector is parallel to the pinning channels, see details in Ref. [3]) at low temperatures ($g > 10$). However, it is instructive to consider first a simple physical picture, as is depicted in Fig. 1, of the vortex motion in a washboard planar pinning potential, *tilted* due the presence of a dimensionless *dc* driving force $0 < \xi^d < \infty$, under the influence of the *effective* dimensionless driving force $\hat{f} = \hat{F}_{px} + \hat{F}_{Lx} = -\sin x + \xi^d$. Here and in the following $\xi^d \equiv j^d/j_c$ and $\xi^a \equiv j^a/j_c$ are the dimensionless *dc* and maximal *ac* current density magnitudes in j_c units, respectively.

If the temperature is zero, the vortex is at rest with $\xi^d = 0$ at the bottom of the potential well of the PPP. When the PPP is gradually lowered by increasing ξ^d, then for $0 < \xi^d < 1$ appears an asymmetry of the left-side and right-side potential barriers for a given potential well, and in this range of ξ^d the effective force $\hat{f}$ changes its sign periodically. With gradual ξ^d -increase there will come a point where $\xi^d = 1$, and for $\xi^d > 1$ the lower right-side potential barrier disappears, the effective motive force $\hat{f}$ becomes positive everywhere along x and the vortex is in the "running" state, i.e. it is periodically changing its velocity with the dimensionless frequency $\omega_i = \sqrt{(\xi^d)^2 - 1}$. So the static CVC of this periodic motion at $\xi^d > 1$ is a result of the time-averaging of the stationary time-dependent solution of the equation of motion $dx/d\tau = \hat{f}$ with $\tau = t/\hat{\tau}$. Eventually, the probability of the vortex overcoming the barriers of the PPP $\nu_0 \equiv \nu_0^{\omega}(\xi^a = 0, \omega = 0)$ at zero temperature is

$$\nu_0 = \begin{cases} 0, & \xi^d < 1, \\ \sqrt{1 - (1/\xi^d)^2} & \xi^d \geq 1, \end{cases} \tag{11}$$

i. e., $\nu_0(\xi^d > 1)$ monotonically tends to unity with increasing ξ^d.

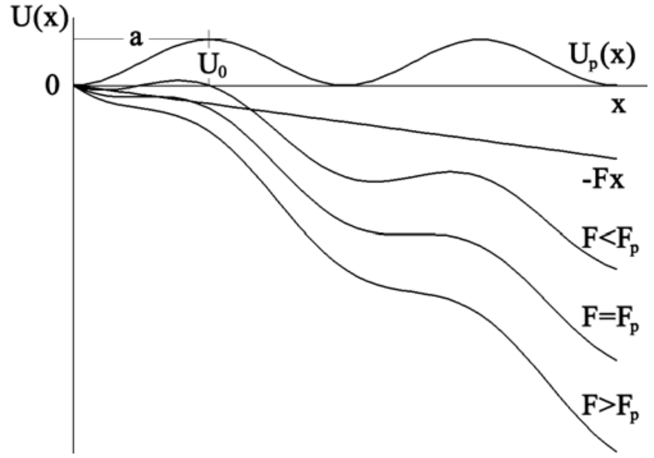

Fig. 1. Modification of the effective pinning potential $U(x) \equiv U_p(x) - Fx$ with gradual increase of the Lorentz force $F = \hat{F}_{Lx}$, where $U_p(x)$ is a periodic pinning potential. If the condition $F < F_p \equiv \hat{F}_{px}$ is satisfied, the initial potential well (with its depth U_0 and width a) is tilted but maintains the average vortex position. When $F = F_p$, the right-side potential barrier disappears. At last, when $F > F_p$, the vortex motion direction coincides with a direction of the moving force F.

If the temperature is nonzero, a diffusion-like mode appears in the vortex motion. At low temperatures ($g \gg 1$) and 0 < ξ^d < 1 the thermoactivated flux-flow (TAFF) regime of the vortex motion occurs by means of the vortex hopping between neighboring potential wells of the PPP. The probability of these hops to occur at low temperatures is proportional to $\exp[-g(1-\xi^d)]$, i. e. strongly increases with T and ξ^d due to the lowering of the right-side potential barriers at their tilting. On the other hand, at ξ^d just above unity (when the running mode is yet weak), the diffusion-like mode can strongly increase the average vortex velocity even at relatively low temperature due to the strong enhancement of the effective diffusion coefficient of an overdamped Brownian particle in a tilted PPP near the critical tilt [12] at $\xi^d = 1$.

Now I consider the influence of a small ($\xi^a \ll 1$) *ac* current density with frequency ω on the CVC in the limit of very small temperatures ($g \gg 1$). In this case the physics of the *dc* response is quite different depending on the ξ^d value with respect to unity. If $\xi^d < 1$, the vortex is mainly localized at the bottom of the potential well where it experiences small ω-oscillation, if one neglects for the moment very rare hops to neighboring wells. The averaging of the vortex motion over the period of oscillations in this case cannot change the CVC which existed in the absence of the *ac*-drive.

If, however, $\xi^d > 1$, the vortex is in the running state with the internal frequency of oscillation $\omega_i = \sqrt{(\xi^d)^2 - 1}$. If $\omega \neq \omega_i$, the CVC is changed only in a second-order perturbation approach in terms of the small parameter $\xi^a \ll 1$ (as it was shown for the analogous resistively shunted Josephson junction problem [13]) because the CVC is not changed in the linear approximation in this case. However, for $\omega = \omega_i$ appears the problem of a synchronization of the running vortex oscillations at the ω_i-frequency

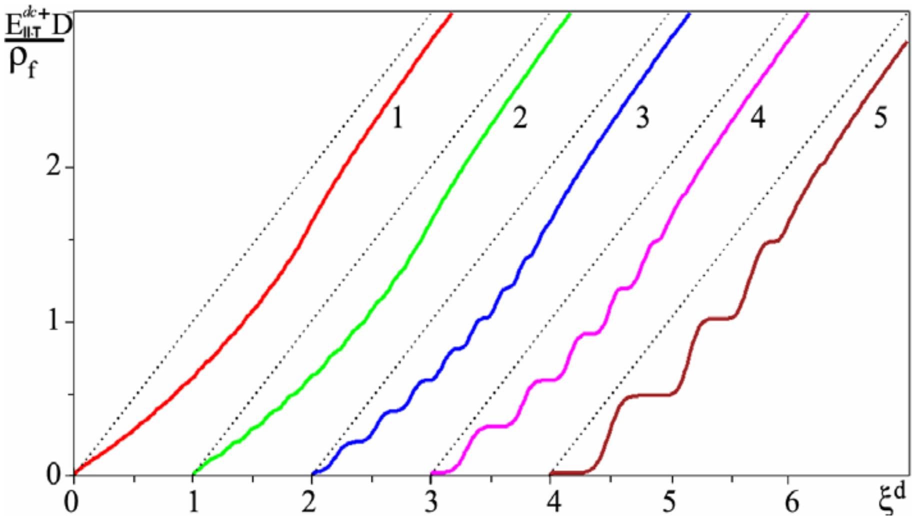

Fig. 2. The longitudinal CVC $E_{\parallel,T}^{dc+}(\xi^d)$ for Ω = 0.05(1), 0.1(2), 0.2(3), 0.3(4), 0.5(5) with ξ^a = 1 and g = 50. Dotted lines correspond to Ohm`s law $E_{\parallel,T}^{dc+} = \xi^d$. The curves 2 – 5 are shifted by 1 along the ξ^d-axis for clarity.

with the external driving frequency ω. As a result, the average over the period of the oscillating vortex velocity is locked in with ω in some interval of the *dc* current density ξ^d even within the frame of a first-order perturbation calculation (see Fig. 2). The width of this first synchronization step (or the so-called "Shapiro step" in the resistively shunted Josephson junction problem) has been found in Ref. [14] and the calculation in the spirit of this reference gives the boundaries of ξ^d where the step occurs as

$$(\xi_\omega^d - \xi^a / 2\xi^d) < \xi^d < (\xi_\omega^d + \xi^a / 2\xi^d). \quad (12)$$

Here ξ^d is the current density which gives $\omega_i = \sqrt{(\xi_\omega^d)^2 - 1} = \omega$, i. e. $\xi_\omega^d = \sqrt{1+\omega^2}$. In this case the size of the first Shapiro-like step on the CVC is ξ^a / ξ_ω^d. In higher approximations (in terms of $(\xi^a)^n$), where n runs through all of the integers) the Shapiro-like steps on the CVC appear at the frequencies $\Omega = n\omega$ and $\Omega_i = n\omega_i$, respectively. The width of the n-th step for $\xi^a \to 0$ is proportional to $(\xi^a)^n$, i. e. strongly decreases with increasing n [15].

In Fig. 3 the longitudinal CVC $E_{\parallel,T}^{dc+}(\xi^d)$ is plotted for various ξ^a showing “Shapiro” steps. The plot in Fig. 3 looks like similar curves discussed earlier [16] for the CVC of the microwave driven resistively shunted Josephson junction model at T = 0 where the overall shape of the CVC and the different behaviour of two types of Shapiro steps in the adiabatic limit was explained. Our graph, in comparison with the curves of Ref. [16], is smoother due to the influence of a finite temperature. The longitudinal CVC $E_{\parallel,T}^{dc+}(\xi^d)$-dependence demonstrates several main features. First, in the presence of the microwave current the *dc* critical current $\xi_c^d(\xi^a)$ is a decreasing function of the *ac* driving force. The physical reason for such a behaviour lies in the replacement of the *dc* critical current by the total *dc* + *ac* critical current. Second, with gradual ξ^a-increase the zero-voltage step reduces to zero and all other steps

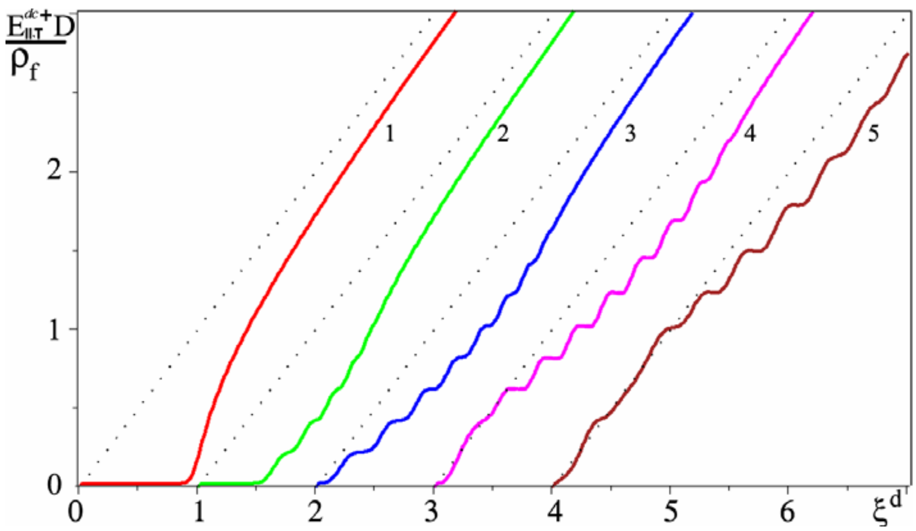

Fig. 3. The longitudinal CVC $E^{dc+}_{\parallel,T}(\xi^d)$ is calculated for ξ^a =0.01(1), 0.5(2), 1(3), 1.5(4), 2(5) with Ω = 0.2, and g = 50. Dotted lines correspond to Ohm`s law $E^{dc+}_{\parallel,T} = \xi^d$. The curves 2 – 5 are shifted by 1 along the ξ^d-axis for clarity.

appear. Such steps are common because they do not oscillate and spread over a *dc*-current range about twice the critical current $2\xi^a_c$. These steps are steps of the first kind and they distort the CVC like a relief bump with a concave shift from the ohmic line. With further ξ^a-increase this relief bump shifts toward higher ξ^d-values. Below this range steps of the second kind appear. These microwave current-induced steps oscillate rapidly and remain close to the ohmic line over the *dc*-current range $\xi^d \le \xi^a - 1$.

To summarize, one can determine three (ξ^d, ξ^a)-ranges where the CVC-behaviour is qualitatively different. In particular, for large *dc* bias current densities ξ^a+1 < ξ^d the CVC asymptotically approaches the ohmic line without microwave induced steps. For an intermediate *dc* current range ξ^a – 1 < ξ^d < ξ^a + 1, the CVC curve deviates from the ohmic line as a series of concave bumps with stable steps. For the lower *dc* current range ξ^d < ξ^a – 1 the steps oscillate with the microwave current along the ohmic line. With gradual Ω-increase the size of the steps increases whereas their number decreases.

CONCLUSION

In the present work I have theoretically examined the influence of the *ac* current on the anisotropic *dc* current-voltage characteristic of a type-II superconductor in the mixed state. A simple physical picture of the vortex motion in a tilted washboard periodic pinning potential has been proposed and elucidates a rich physics arising from the combination of *dc* and *ac* driving, and the low temperature mediated vortex hopping (or running) in a washboard pinning potential. The experimental realization of this model in thin-film geometry [8, 9, 17] opens up the possibility for a variety of experimental studies of directed motion of vortices under (*dc*+*ac*) – driving simply by measuring longitudinal and transverse voltages. Experimental control of the frequency and magnitude of the driving forces, damping, Hall constant, pinning parameters and temperature can be effectively realized.

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