Dynamic contribution to the fishtail effect in a twin-free DyBa$_2$Cu$_3$O$_{7-\delta}$ single crystal

A.J.J. van Dalen a, *, M.R. Koblishka a,1, R. Griessen a, M. Jirsa b, G. Ravi Kumar c

a Free University, Faculty of Physics and Astronomy, De Boelelaan 1081, 1081 HV Amsterdam, The Netherlands
b Institute of Physics, Academy of Sciences of the Czech Republic, Na Slovance 2, CZ-18040, Praha 8, Czech Republic
c Solid State Physics Division, Bhabha Atomic Research Centre, Bombay-400085, India

Received 19 May 1995

Abstract

Superconducting current densities $j_c$ and dynamic relaxation rates $Q = d \ln j_c / d \ln(dB_c/dr)$, where $dB_c/dr$ is the sweep rate of the external magnetic field $B_c$, were measured as a function of temperature ($5 \text{ K} < T < 65 \text{ K}$) in magnetic fields up to 7 T on a twin-free DyBa$_2$Cu$_3$O$_{7-\delta}$ single crystal by means of a high-sensitivity capacitance torque magnetometer. Above 15 K, we observe a "fishtail" effect, i.e. a pronounced minimum in the $j_c(B_c)$ curve at fields around $B_c = 1 \text{ T}$. The relaxation rate $Q$ shows an anomalous increase at low fields which is correlated to the minimum in the $j_c(B_c)$ curve. Both the $j_c$ versus $B_c$ and $Q$ versus $B_c$ data are used as input parameters into the generalized inversion scheme developed by Schnack et al. [Phys. Rev. B 48 (1993) 13176] to calculate the true critical current density $j_c$ which is by definition independent of relaxation effects. Interestingly, the $j_c(B_c, T)$ curves derived in this way do not show a minimum. This points clearly to a dynamic contribution to the fishtail effect. The true critical current density $j_c(B_c, T)$ decreases weakly with increasing $B_c$ over the entire measured temperature and field range, as expected for single-vortex pinning. This indicates that the observed fishtail effect is not caused by a crossover from single-vortex pinning to pinning of flux bundles. The temperature dependence of $j_c$ is in good agreement with the predictions of a model based on single-vortex pinning caused by spatial fluctuations in the charge-carrier mean free path.

1. Introduction

The fishtail or peak effect in the field dependence of the superconducting current density $j_c(B_c)$ is observed in many bulk and single crystalline samples of high-$T_c$ superconductors [1], in (RE)Ba$_2$Cu$_3$O$_{7-\delta}$ (RE = rare earth), (La$_{1-x}$Sr$_x$)$_3$CuO$_4$, Tl$_2$Ba$_2$CaCu$_2$O$_{8+}$, as well as in Bi$_2$Sr$_2$CaCu$_2$O$_{8+}$, where the feature is often called "arrowhead". Presently, the origin of the fishtail behavior is controversially discussed in the literature. Explanations for this effect include static approaches (where the fishtail effect is present independently of the time scale of the measurement) with granularity on different length scales appearing in the samples at large external magnetic fields [2–5], special impurity phases acting as additional pinning centers [6] or dynamic approaches, where the fishtail shape is caused by creep in the vortex lattice [7,8] or...
a crossover in the pinning regimes, e.g. from single-vortex pinning to a pinning of vortex bundles [9]. A comparison of data from many different samples showing the fishtail effect [10] demonstrates clearly that the fishtail is caused by an interaction of the flux-line lattice with defect structure and not by a specific defect structure itself. In the latter case the presence and the strength of the fishtail effect would vary strongly from sample to sample due to their widely varying pinning properties, which is in contradiction with the experimental observations.

Analysing the flux-creep equation for high-temperature superconductors, Schnack et al. [11] showed that for a correct interpretation of experimental results it is essential to take creep effects into account by which all measured current densities \( j_c (T, B_c) \) (often improperly called critical current) are strongly affected. Therefore, for investigation of the pinning mechanisms we need to calculate the true critical current density, \( j_c (T, B_c) \) which is by definition independent of relaxation effects, i.e. \( U(j_c(T, B_c), T, B_c) = 0 \) with \( U \) denoting the current density, temperature and magnetic field dependent activation energy. Since it is not influenced by relaxation effects, the true critical current density \( j_c \) is only related to the pinning properties of the sample. In this work the so called generalized inversion scheme (GIS) developed by Schnack et al. [12,13] is used to determine \( j_c (T, B_c) \).

The paper is organized as follows: The experimental method is described in Section 2. Section 3 presents measurements of the superconducting current density \( j_s \) and the dynamic relaxation rate \( Q \) as a function of \( B_c \), \( dB_c/dt \) and \( T \). In Section 4 the true critical current density \( j_c \) is calculated from the temperature dependence of the induced current density \( j_1 \), and of the dynamic relaxation rate \( Q \) by means of the GIS procedure. In Section 5, we discuss the temperature dependence of \( j_c \) and demonstrate that it is for all magnetic fields \( 0.5 \ T \leq B_c \leq 6.5 \ T \) in good agreement with a model of single vortices pinned by spatial fluctuations in the charge-carrier mean free path. In Section 6, the field dependence of \( j_c \) is analyzed, and we show that \( j_c (B_c) \) does not exhibit a minimum in contrast to \( j_1 (B_c) \). The relatively weak decrease of \( j_c (B_c) \) with field is consistent with the predictions of a single-vortex pinning model.

2. Experimental methods

For the measurements of currents and dynamic relaxations, we selected a twin-free Dy-Ba\(_2\)Cu\(_4\)O\(_{7-\delta}\) (DyBCO) single crystal. DyBCO single crystals were grown using a slow-cooling method with the same preparation route as for YBa\(_2\)Cu\(_3\)O\(_{7-\delta}\) single crystals [14]. These as-grown DyBCO single crystals proved to be twin free in previous experiments using polarizarion microscopy and magneto-optic observations of flux distributions which are very sensitive to the presence of twin planes [15]. The dimensions of the crystal \((a \times b \times c) = 0.97 \times 1.23 \times 0.015 \ mm^3\). Torque hysteresis loops are measured using a capacitance torque magnetometer [16] consisting of two capacitance plates connected by a pair of crossed phosphorous bronze springs. The external magnetic field \( B_z \) ranges from zero to 7 T and is oriented at an angle \( \Theta \) with respect to the \( c \)-axis of the crystal. A change of an external magnetic field induces a non-equilibrium superconducting current density \( j_s \) flowing in the \((a, b)\) planes. The corresponding magnetic moment is given by [17]

\[
m = \frac{1}{3} a^2 d (b - a/3) j_s,
\]

where \( d \) is the thickness of the sample and \( a \) and \( b \) are the sample dimensions with \( b \geq a \). The magnitude of the torque \( \tau = m \times B_c \) is simply equal to \( \tau = m \sin \Theta \) because the magnetic moment is oriented parallel to the \( c \)-axis. In this work the angle \( \Theta \) is fixed at 5°. The dynamic relaxation rate \( Q \) is obtained from measurements of \( j_s \) at various sweep rates \( dB_c/dt \) of the magnetic field. As shown by Jirsa et al. [18] and Pfütz et al. [19] the dynamic relaxation rate \( Q \) is closely related to the conventional relaxation rate \( R = d \ln M/dt \) where \( M \) is the magnetization of the sample. This observation is supported by numerical calculations by Schnack et al. [11].

3. Experimental results

In Fig. 1 the field dependence of the superconducting current density \( j_s \) measured at a sweep rate of 40 mT/s is shown for several temperatures between 5 K and 65 K. Above \( T = 10 \) K \( j_s \) exhibits a fishtail effect, i.e. there is a minimum in \( j_s \) at low
fields (around 1 T at 25 K) and a maximum (around 3 T at 25 K). This is in agreement with results from Werner et al. [10] who found that the fishtail effect starts to develop at \( T = 10 \) K. Both the minimum and the maximum shift towards lower fields when \( T \) is increased.

To explore whether or not this behavior is related to flux-creep relaxation effects we determined the relaxation rate \( Q \) from measurements of \( j_i \) at various sweep rates. A typical example of such a measurement is shown in Fig. 2(a) at \( T = 20 \) K. The magnetic field was ramped up and down with a sweep rate of \( dB_x/dt = 40, 30, 20, 15 \) and 10 mT/s, respectively.

A closer inspection of Fig. 2(a) reveals, however, that on the low-field side of the minimum \( j_i \) is less dependent on \( B_x \) at low sweep rates (10 mT/s) than at high sweep rates (40 mT/s). This implies that relaxation effects are more pronounced at low fields. This is confirmed by the data in Fig. 2(b), where the dynamic relaxation rate decreases rapidly with increasing field below \( \approx 2 \) T. The field dependence of the dynamic relaxation rate obtained at several temperatures between 5 K and 65 K is shown in Fig. 3.

At \( T = 20 \) K a shallow minimum in the field dependence of the relaxation rate develops, which gets more pronounced and shifts to lower fields for increasing temperature.

In contrast to the data reported by Krusin-Elbaum et al. [9] we do not find that \( j_i(B_x) \) and \( Q(B_x) \) are mirror images. Our data show that the maximum in \( Q(B_x) \) is always shifted to lower fields with respect to the minimum in \( j_i(B_x) \). It suggests that the fishtail originates from two effects:

1. a relatively sharp decrease of the true critical current density \( j_c \) with increasing field and
2. a field-dependent relaxation rate.

**4. Determination of the true critical current density \( j_c \)**

In order to test quantitatively the expectation that the field-dependent relaxation rate contributes to the fishtail effect we determined the true critical current density \( j_c \) by using the GIS [12]. This method calcu-
lates the true critical current density \( j_c \) and the activation energy \( U_c \) from measurements of the induced current density \( j_i \) and the dynamical relaxation rate \( Q \) as a function of temperature at a fixed value of the external field \( B_x \). The GIS is based on the general assumption that the temperature and current dependence of the activation energy can be separated as

\[
U(j_i, T, B_x) = g(T, B_x) f(J),
\]

and on the existence of a power-law relation between \( U_c \) and \( j_c \)

\[
g(T, B_x) \propto U_c(T, B_x) \propto [j_c(T, B_x)]^p G(T).
\]

Here \( J = j_i/j_c \) and \( U_c \) is the characteristic pinning energy. The thermal function \( G(T) \) can be derived [12] from the temperature dependence of the prefactor in the power-law relation between \( j_c \) and \( U_c \) from which it follows that \( G(T) \) depends on the pinning regime and the dimensionality of the sample. For example, for a collectively pinned single vortex in the 3D regime the activation energy is given by \( U_c = j_c \Phi_0 L_c \xi^2 \) and the correlation length is given by \( L_c \propto j_c^{-1/2} \). In the 2D regime \( L_c \) is restricted by the thickness of the layers and thus independent of the current density. These results lead to \( p = \frac{1}{2} \) and \( p = 1 \) for a collectively pinned single vortex in three and two dimensions, respectively. The pinning barrier for a flux bundle is in the collective pinning theory given by \( U_c = j_c B_x V_B/\rho \) where \( V_B \) is the volume of the flux bundle and \( \rho \) is the range of the potential. A scaling approach has been used to obtain the volume of the flux bundle [20], leading to expressions for \( U_c \) and \( j_c \) [21]. Elimination of the relevant length scale leads again to a power-law relation between \( U_c \) and \( j_c \) in the regime of flux bundles [12] with \( p = -\frac{3}{2} \) for small bundles and \( p = -\frac{1}{2} \) for large bundles. For the 2D case one finds similar results with \( p = -\frac{3}{2} \) and \( p = 0 \) for small and large bundles, respectively [12,22]. An overview of the calculations leading to values of \( p \) in various pinning regimes and dimensionalities is given in Appendix B of Ref. [13]. The GIS treats the exponent \( p \) as a free parameter. However, from the previous discussion follows that only the values \( p = -\frac{3}{2}, -\frac{1}{2}, 0, \frac{1}{2} \) and 1 are physically meaningful.

During a sweep of the applied magnetic field the activation energy is given by [12] \( U(j_i, T, B_x) = CkT \), which together with Eq. (2) leads to \( \ln(CkT) = \ln g(T, B_x) + \ln f(J) \). The parameter \( C = \ln(2\nu_0 B_x/\alpha dB/\alpha t) \) in which \( \nu_0 \) is an attempt velocity and \( R \) is the effective radius of the sample, can be determined as follows [13].

By taking the total derivative of \( \ln(CkT) \) (see above) with respect to \( \ln T \) and \( \ln dB/dt \), one finds by elimination of \( d\ln f/d\ln J \) the following relation for \( C \) (see Ref. [13]):

\[
\frac{d\ln j_c}{dT} = -\frac{d\ln j_c}{dT} + C \frac{Q}{T},
\]

which is only valid at sufficiently low temperatures where the temperature dependence of \( g(B_x, T) \) can be neglected. Since we assume that \( j_c(T, B_x) \) can be written as a product of a field- and a temperature-dependent function, the factor \( d\ln j_c/dT \) is independent of magnetic field. Furthermore, assuming that \( C \) does not depend on the magnetic field [13] it follows from Eq. (4) that \( C \) can be determined from a plot of \( -d\ln j_c/dT \) versus \( Q/T \) for the field values chosen for our experiment. This leads to \( C = 10.6 \pm 1.3 \) in the application of the GIS we used \( C = 11 \).

The measured current densities and relaxation rates are affected by quantum creep for temperatures \( T < 5 \) K. Therefore, these data cannot be used in the
GIS, since the GIS takes only thermally activated vortex motion into account. For the inversion of the measured data, the low-temperature region needs, therefore, to be replaced by extrapolations from the temperature range where quantum creep is negligible. Since it is observed that \( j_s \) decreases exponentially with increasing temperature except at low temperatures, extrapolation of the linear part of the \( \ln j_s \) versus \( T \) curve leads to \( j_s(0) \) which is the current density that would be observed at \( T = 0 \) K in the absence of quantum creep. This current density is equal to the critical current density \( j_c(0) \) because of the absence of thermally activated flux motion at \( T = 0 \) K. The field dependence of \( j_s(0) \) is shown in Fig. 4(a). The weak decrease of \( j_s(0) \) is consistent with a single-vortex pinning model as outlined in Section 6.

The measured temperature dependence of the relaxation rate \( Q(T) \) is also used for the determination of the characteristic pinning energy \( U_c \). One assumes that at low temperatures (but in the absence of quantum creep) and at constant field the \( j_s \) dependence of \( U \) is well described by the interpolation formula of Feigel'man et al. [20]

\[
U(j_s) = \frac{U_c}{\mu} \left[ \left( \frac{j_s}{j_c} \right)^\mu - 1 \right].
\]

Using the fact that during a sweep of the applied magnetic field the pinning energy is given by \( U(j_s, T, B_s) = CkT \) where \( C \) is defined above, one can evaluate the ratio \( j_s/j_c \) with the help of Eq. (5). Together with the definition of the dynamic relaxation rate \( Q \), this leads to [23]

\[
\frac{T}{Q} = \frac{U_c}{\mu} + \mu CT.
\]

In Fig. 5 \( T/Q \) is displayed as a function of temperature for \( B_s = 0.5, 1, 2, 4 \) and \( 6.5 \) T. Similar results are obtained at other magnetic fields. The deviations at high temperature are due to the temperature dependence of \( U_c \) and to thermal fluctuations near the irreversibility line. Extrapolation of the linear part of the curves to \( T = 0 \) K gives the value of the pinning energy \( U_c \) at low temperatures. Repeating this procedure at all the values of the magnetic field considered in this work finally gives the field dependence of \( U_c \) shown in Fig. 4(b). The values of \( j_s(B_s) \) and \( U_c(B_s) \) shown in Figs. 4(a) and (b), respectively, are used in the GIS because they correspond in fact to \( j_s(T, B_s) \) and \( Q(T, B_s) \) at \( T = 0 \) K.

The GIS with various values of \( p \) has been applied to our measurements of \( j_s \) and \( Q \) for magnetic fields ranging from \( B_s = 0.5 \) T to \( B_s = 6.5 \) T in steps of 0.25 T. We found that the choice \( p \leq 0 \) leads to a divergence of \( g(T) = U_s(T)/U_s(0) \) and to \( j_s > j_c \) above a certain temperature. This would indicate that the vortex system is in the flux-flow regime,
For 3D single-vortex pinning due to fluctuations in the transition temperature ($\delta T_c$ pinning) one finds
\begin{equation}
\frac{j_c(t)}{j_c(0)} = (1 - t^2)^{3/6}(1 + t^2)^{5/6}
\end{equation}
\begin{equation}
\frac{U_c(t)}{U_c(0)} = g(t) = (1 - t^2)^{1/3}(1 + t^2)^{5/3}
\end{equation}

where $t = T/T_c$. Note that there are no fitting parameters in Eqs. (7) and (8). The choice $p = \frac{1}{2}$ in the inversion gives a remarkably good agreement between the experimentally determined functions.

Fig. 6. Temperature dependence of the critical current density $j_c$ (a) and the characteristic pinning energy $U_c$ (b) calculated from the measured current densities $j_c$ and relaxation rate $Q$ with the generalized inversion scheme at $B_e = 0.5, 1, 2, 4$ and 6.5 T. The solid lines represent the theoretical predictions based on pinning caused by fluctuations in the mean free path [24]. The dotted lines represent the theoretical predictions based on pinning caused by fluctuations in $T_c$. Note that these lines are not a fit to the data, since they are entirely determined by $t = T/T_c$ [See Eqs. (7) and (8)]. Inversion at other magnetic fields give similar results.

5. Temperature dependence of $j_c$

The calculated values of $j_c(t)/j_c(0)$ and $g(t)$, where $t = T/T_c$, for $B_e = 0.5, 1, 2, 4$ and 6.5 T are shown in Figs. 6(a) and (b), respectively. Calculations for other values of $B_e$ give similar results. In the inversion we have used $p = \frac{1}{2}$, for all values of the magnetic field, which is justified by the arguments given above. For 3D single-vortex pinning due to fluctuations in the charge-carrier mean free path ($\delta l$ pinning) the temperature dependence of $j_c$ and $U_c$ are predicted to be [12,13,24]
\begin{equation}
\frac{j_c(t)}{j_c(0)} = (1 - t^2)^{5/2}(1 + t^2)^{-1/2},
\end{equation}
\begin{equation}
\frac{U_c(t)}{U_c(0)} = g(t) = 1 - t^4.
\end{equation}
\( \frac{j_c(t)}{j_c(0)} \) and \( g(T) \) and the theoretical predictions for \( \delta l \) pinning at fields 0.5 \( T < B_s < 6.5 \) T. This leads to the conclusion that the 3D single-vortex pinning regime extends from \( B_s = 0 \) T up to \( B_s \geq 6.5 \) T at all temperatures between 5 K and 65 K which implies that the fishtail effect is not caused by a crossover of pinning regimes. This contrasts with the results of Krusin-Elbaum et al. [9] who conclude that in a \( \text{YBa}_2\text{Cu}_3\text{O}_7 \) single crystal a crossover from single-vortex pinning to pinning of flux bundles takes place at \( T = 45 \) K independently of the magnetic field. Furthermore, the results displayed in Fig. 6 clearly demonstrate that the pinning is caused by \( \delta l \) pinning and not by \( \delta T_c \) pinning. It is interesting to note that this has also been observed earlier in \( \text{YBa}_2\text{Cu}_3\text{O}_7 \) and \( \text{YBa}_2\text{Cu}_3\text{O}_4 \) films [24].

Another indication that the fishtail effect is not caused by a crossover of pinning regimes as proposed in Ref. [9] is given by the observation of a pronounced minimum in the \( j_c(B_s) \) curves of a \( \text{Pb} \) irradiated \( \text{DyBa}_2\text{Cu}_3\text{O}_{7-\delta} \) single crystal [25]. If indeed a crossover so defined between pinning regimes would be the cause of the fishtail effect, then the introduction of strong pinning centers into a sample showing already the fishtail effect should lead to the disappearance of this effect, since \( j_c \) is increased significantly and the boundary between single-vortex pinning and pinning of bundles should be shifted to higher temperatures and fields. However, torque measurements on a \( \text{Pb} \) irradiated \( \text{DyBa}_2\text{Cu}_3\text{O}_{7-\delta} \) single crystals [25] at two different orientations of the field show a clear fishtail effect at temperatures above \( T = 20 \) K. This indicates that the fishtail effect is not related to a specific arrangement of pinning centers. These conclusions are supported by results of Wen et al. [26] who observed that the position of the minimum of the \( j_c(B_s) \) curves does not scale with the critical current density.

6. Magnetic-field dependence of \( j_c \)

Since the GIS is used to obtain \( j_c(T) \) at fixed values of the magnetic field one has to recollect the results from inversions performed with data obtained at different magnetic fields in order to construct the \( j_c(B_s) \) curves. These curves are shown in Fig. 7. In sharp contrast to the measured \( j_c(B_s) \), the true critical current density obtained by means of the GIS does not exhibit a minimum but remains virtually constant or even decreases when the magnetic field increases from 1 T to 6 T. The \( j_c(B_s) \) curves show a plateau for \( 2 T < B_s < 6 \) T for temperatures below \( T = 30 \) K. For higher temperatures the critical current density decreases monotonically with increasing field. The measured induced current density and the calculated critical current density obtained at \( T = 35 \) K and 45 K are redrawn in Fig. 8 for clarity. This figure demonstrates two important features:

1. The induced current density is approximately four times lower than the critical current density, which demonstrates that the usual assumption \( j_s = j_c \) is not valid in this case.

2. In contrast to the induced current density \( j_s \), which shows a pronounced fishtail behaviour, the critical current density \( j_c \) decreases monotonically with increasing field.

The absence of a minimum in \( j_c \) demonstrates clearly the dynamic character of the observed fishtail effect.

From Fig. 4(b) one can see that the activation energy increases slightly with increasing field strength between \( B_s = 1 \) T and \( B_s = 5 \) T, whereas an enhanced increase is observed for fields larger than 5 T. This increase of \( U_c \) with field causes a reduction of the relaxation rate which leads to a smaller differ-
The observed weak field dependence of the critical current density $j_c$ can be compared with predictions of the collective-pinning theory for small and large vortex bundles [20,21,27]. Together with an estimate for the crossover field $B_{cb}$ where the system changes from single-vortex pinning to pinning of small bundles this leads to a determination of the pinning regime.

The crossover field $B_{cb}$ is determined by the condition $L^* = \alpha \xi_0$ where $L^*$ is the single-vortex correlation length along the $c$-axis, $\alpha = (m/M)^{1/2}$ is the mass anisotropy and $\xi_0 = (\Phi_0/B_0)^{1/2}$ is the average distance between vortices. This condition is equivalent to $B_{cb} = \beta \xi_0 J_c/J_0$ with $\beta = 5$, $J_0 = 4B_0/3\sqrt{6} \mu_0 \lambda$ is the depinning current density, $\mu_0 B_{c2} = \Phi_0/2\pi \xi_0^2$ is the upper critical field and $B_c = \Phi_0/2\sqrt{2} \pi \lambda \xi_0$ is the thermodynamic critical field. From the collective-pinning theory follows that the length $L^*$ is given by [28]

$$L^* = \xi^{1/2} \left( \frac{J_0}{J_c} \right)^{1/2}.$$

At low temperatures $j_c = 2 \times 10^{10}$ A/m$^2$ which, together with the numerical values $\lambda = 150$ nm, $\xi = 2$ nm and $\alpha \approx 0.2$ leads to $L^* \approx 4.3$ nm. The corresponding value $B_{cb} \approx 3.7$ T in itself does not exclude the possibility of small-bundle pinning for the high-field data presented in this work. Therefore, a more detailed analysis is required to determine the type of pinning in our case. This is done by comparing the observed field dependence of the critical current density with predictions from collective-pinning theory, where the anisotropy has to be included.

The field dependence of the critical current density in the bundle-pinning regime can be evaluated by noting that the Lorentz force on the vortex bundle is balanced by the pinning force which leads to $J_c = U_c/B_c V_c$ where $V_c$ is the volume of the flux bundle and $U_c$ is the characteristic pinning energy. Comparing this energy to the shear energy $c_{66}(\xi/R_c)^{2} V_c$ and using the expressions for the transverse bundle size $R_c$ derived by Blatter et al. [29] together with $c_{66} \approx \Phi_0 B_c/(8\pi \lambda)^2$ leads to

$$j_c \approx j_{c1} \left( \frac{L^*}{\xi_0 \alpha_0} \right)^2 \exp \left( -2 \frac{c_{66}}{\xi_0 \alpha_0} \right)$$

in the local limit which corresponds to small-bundle pinning. Here $j_{c1}$ is the critical current density in the single-vortex regime and $c_{66}$ is a constant of order unity. Note that the maximum of this function occurs at $\alpha_0 = (3c_{66}/L^*)^{1/3} \xi_0 / \xi$ which corresponds to a field lower than the crossover field $B_{cb}$. Therefore in its domain of validity, Eq. (10) predicts that $j_c$ decreases with increasing magnetic field. In the non-local limit (large-bundle pinning) one finds

$$j_c \approx j_{c2} \left( \frac{\xi_0 \alpha_0 \lambda}{L^*} \right)^4.$$

Eqs. (10) and (11) describe the influence of a magnetic field on the critical current density once $\lambda$ and $L^*$ are known. They both predict a rather strong field dependence of $j_c$ in the small-bundle limit since application of Eq. (10) gives $j_c(6 \text{ T})/j_c(3.7 \text{ T}) \approx$
0.35, where $B_e = 3.7$ T has been chosen as lowest field value, since Eq. (10) is not valid for fields $B_e < B_{cr} \approx 3.7$ T. At high temperatures $j_e$ decreases, leading to an increase of $L_v$ which finally leads to a stronger field dependence of $j_e$. These results are in contradiction with the results displayed in Fig. 7 where it is shown that $j_e$ is only weakly field dependent. Note that a precise knowledge of the anisotropy is not necessary, since only the ratio $L_v/\varepsilon = \xi/\eta/\xi_e$ plays a role in the determination of the critical current density in the small-bundle pinning limit. Application of Eq. (11) for large-bundle pinning leads directly to $j_e \propto B_e^{-3}$ which is in total disagreement with the experimental results. Therefore we conclude that the weak dependence of $j_e$ on magnetic field indicates that the system is indeed in the single-vortex regime.

A similar conclusion was reached for oxygen-deficient YBa$_2$Cu$_3$O$_{7-x}$ films, for which $U_c$ and $j_e$ were determined independently. The work showed that $U_c \propto (j_e)^{1/2}$ exactly as predicted for single-vortex pinning [see also Eq. (3) and the discussion in the text]. This shows [30] that in YBa$_2$Cu$_3$O$_{7-x}$ the vortices are individually pinned for fields up to 7 T and nominal oxygen contents $x$ ranging from $x = 6.55$ to $x = 6.98$.

7. Conclusions

In summary, calculations of the critical current density by means of the GIS reveal that the true critical current density $j_c(B_e)$ does not show a minimum. This demonstrates that the observed minimum in the measured $j_c(B_e)$ curves is caused by a higher relaxation at low fields corresponding to the reduced activation energy. This clearly demonstrates the presence of a dynamic contribution to the so-called “fishtail” effect. The fact that $j_e$ is only weakly dependent on $B_e$ indicates that the system is in the single-vortex limit for fields up to 6.5 T. This conclusion is further supported by the observation that the best agreement between the calculated and theoretical temperature dependence of both the activation energy and the critical current density is achieved for $U_c \propto j_e^p$ with $p = \frac{1}{2}$ for all magnetic-field values used in this work. This indicates that the fishtail effect in DyBa$_2$Cu$_3$O$_{7-\delta}$ is not caused by a crossover from single-vortex pinning to pinning of vortex (super) bundles. Furthermore, it is shown that the pinning is due to spatial fluctuations in the charge-carrier mean free path, as was found earlier by Griessen et al. [24] in several YBa$_2$Cu$_3$O$_7$ and YBa$_2$Cu$_3$O$_x$ films.

Acknowledgements

This work is a part of the research program of the Stichting voor Fundamenteel Onderzoek der Materie which is financially supported by NWO. GRK is grateful to the European community for granting a fellowship under the Indo-EC agreement. One of us (MJ) could participate thanks to the fellowship of the Commission of the EC (No. 2395).

References