

Magnetic relaxation in high-temperature superconductors

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We review experimental studies of the time decay of the nonequilibrium magnetization in high-temperature superconductors, a phenomenon known as magnetic relaxation. This effect has its origin in motion of flux lines out of their pinning sites due to thermal activation or quantum tunneling. The combination of relatively weak flux pinning and high temperatures leads to rich properties that are unconventional in the context of low temperature superconductivity and that have been the subject to intense studies. The results are assessed from a purely experimental perspective and discussed in the context of present phenomenological theories. [S0034-6861(96)00403-5]

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I. INTRODUCTION

The discovery of high-temperature superconductivity in copper-oxide-based materials produced a tremendous euphoria in both scientific and technological communities. Some of this euphoria diminished as researchers found that the persistent supercurrents in these materials decay in time. The difficulty was not ordinary resistivity, but rather a dissipation occasioned by the thermally activated motion of magnetic flux inside the superconductor. This magnetic relaxation has now become a fascinating and widely studied phenomenon on its own right. Meanwhile, technological applications have gone forward successfully anyway, adapting and in some respects even benefiting from the magnetic relaxation.

The relaxation is typically observed in measurements of the magnetic dipole moment of high-temperature superconductors using, for example, a vibrating sample magnetometer or a magnetometer based on a SQUID (superconducting quantum interference device). The magnitude of the dipole moment is observed to decrease with time, in an approximately logarithmic manner.

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Large magnetic relaxation rates have been observed in all the known families of high-temperature superconducting materials, including LaBaSrCaCuO, YBaCuO, BiSrCaCuO, TlBaCaCuO, HgBaCaCuO, and NdCeBaCuO, and in all the various material forms, including ceramic powders, single crystals, films, tapes, and fibers.

Magnetic relaxation in superconductors was first studied in low-temperature superconductors, but the effect was typically so small that it took especially sensitive experimental techniques to detect it. A model explaining this effect was first described by Anderson and Kim (1962, 1964). They introduced the basic concept of thermal activation of magnetic flux lines (see Sec. II.A) out of pinning sites, which proceeds at a rate proportional to $\exp(-U/kt)$, where U represents an activation energy, k is Boltzmann's constant, and T is the absolute temperature. This process leads to a redistribution of flux lines, hence of the current loops associated with the flux lines, thus causing a change in the magnetic moment with time. The Anderson-Kim model predicted a magnetic moment changing logarithmically with time, as has been observed in experiment (Kim *et al.*, 1962). Slow physical changes due to exponential activation processes are generally referred to as "creep," and so the Anderson-Kim mechanism became referred to as "magnetic flux creep."

The original magnetic relaxation data reported by Müller *et al.* (1987) on the first high-temperature superconducting material (LaBa)₂CuO₄ were interpreted in terms of a very different model. They recognized that the coupling between grains in their ceramic powder sample was weak, and that a random array of such weak links gives rise to a frustrated glassy state that could relax in a logarithmic way. This phenomenon may well contribute to the relaxation of ceramic-powder compacts of high-temperature superconducting material, but it is surely only part of the story. This became evident with the discovery by Worthington *et al.* (1987) and by Yeshurun and Malozemoff (1988) of comparably strong relaxation in bulk single crystals of YBa₂Cu₃O₇ (YBCO), as is illustrated in Fig. 1. In this figure, the moment per unit volume, or magnetization M , is measured with a SQUID magnetometer after the sample is first cooled in zero applied field to the measurement temperature, after which a field of 0.6 kOe is applied. The moment is negative, or diamagnetic. The data are normalized to M_0 , the first value measured, and sets of data are shown at different temperatures. The rate of magnetic relaxation is large, with a 20% change at 50 K over two decades in time, and the normalized rate further increases with increasing temperature. It was natural to apply the Anderson-Kim model to the results obtained in these crystals. The large size of the effect was explained in terms of relatively small activation energies of the order 10–100 meV. Because of its size, the effect in high-temperature superconductors has been termed rather ominously "giant flux creep."

However, it was soon found that the basic Anderson-Kim theory failed to explain some important details of

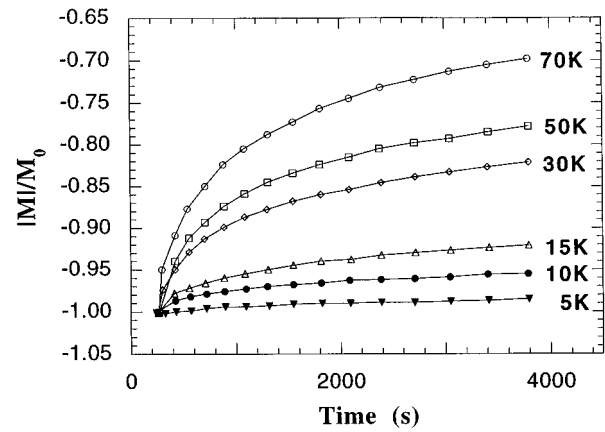


FIG. 1. Demonstration of the size of the relaxation phenomenon, showing the decay of the normalized magnetization as a function of time for a YBCO crystal. The field $H=0.6$ kOe, parallel to the c axis, was applied after cooling the sample in zero field. (After Yeshurun *et al.*, 1988a.)

the magnetic relaxation effect in high-temperature superconductors. For example, the functional form of the time relaxation was found to be only approximately logarithmic. More careful studies over many orders of magnitude in time revealed deviations from logarithmic-time behavior, as illustrated in Fig. 2 for an YBa₂Cu₃O₇ crystal. This figure, from the work of Thompson, Sun, and Holtzberg (1991), shows a semilogarithmic plot of the diamagnetic moment per unit volume versus time after application of a 1 T field. Deviations from the expected straight line are evident. These deviations are discussed in detail in Sec. III.B.3.

Noticeable disagreements with the basic Anderson-Kim model were also observed in the temperature dependence of the magnetic relaxation, described in detail in Sec. III.B.1. These led to a rapid development of new

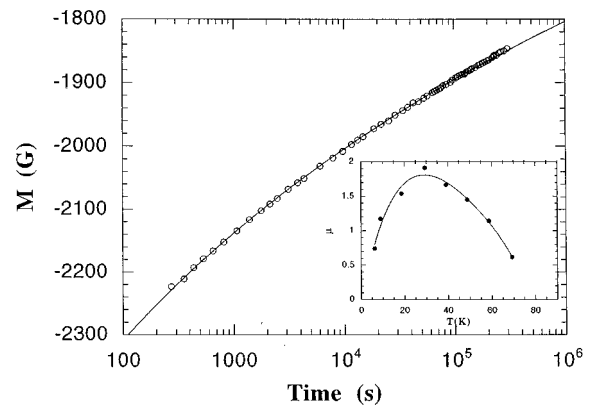


FIG. 2. Semilogarithmic plot of the magnetization M vs time for a single YBCO crystal at $T=30$ K, with a field $H=1$ T applied parallel to the c axis. Deviations from a logarithmic decay are apparent. The solid line shows a fit to the interpolation formula, Eq. (4.21). Inset: The μ values extracted from the fit to Eq. (4.21). (After Thompson, Sun, and Holtzberg, 1991.)

phenomenological models that built on the early work of Beasley *et al.* (1969) to consider complex current-dependent activation energies. These models are reviewed in Sec. IV. Their ability to explain certain important aspects of the experimental data is one of the significant successes of this field. The microscopic basis of these models involves the collective interaction of flux lines. Hence these theories are often referred to in the context of “collective flux creep.” These theories have many parameters and, with only a few exceptions, comparison with experiment is still rather incomplete.

Magnetic relaxation at ultralow temperatures is unexpected on the basis of thermal activation of flux lines. However, a residual, temperature-independent relaxation was measured at milliKelvin temperatures, which has led to the novel concept of relaxation based on a quantum tunneling of vortices (Mota *et al.*, 1988, 1991; Simanek, 1989; Hamzic *et al.*, 1990; Blatter *et al.*, 1991; Fruchter, Hamzic, *et al.*, 1991; Fruchter, Malozemoff, *et al.*, 1991; Ivlev *et al.*, 1991). Relaxation at ultralow temperatures is briefly discussed in Sec. III.B.1 and in the introduction to Sec. IV.

Thermally activated flux motion is in evidence in more than just direct measurements of the magnetic dipole moment. It also causes the broadening of the magneto-resistive curves (Graybeal and Beasley, 1986; Tinkham, 1988a; Palstra *et al.*, 1988; Malozemoff *et al.*, 1989; Iye *et al.*, 1990) and it determines the shape of the voltage-current (V - I) or electric-field/current-density (E - J) curves, causing them to show a smooth, power-law behavior (Koch *et al.*, 1989). The intimate relationship between flux motion and V - I curves was recognized by many authors (Griessen, 1991a; Ries *et al.*, 1992; Sandvold and Rossel, 1992; van der Beek, Nieuwenhuys, *et al.*, 1992; Konczykowski *et al.*, 1993; Caplin *et al.*, 1994). Its origin lies in the fact that flux motion is associated with a Lorenz force $(1/c)\mathbf{J}\times\mathbf{B}$ that can be viewed as driving the flux lines from their pinning sites, and with an electric field $\mathbf{E}=(1/c)\mathbf{B}\times\mathbf{v}$, where \mathbf{v} is the average velocity of the flux lines in the direction of the Lorenz force (see Sec. II.C). For an unpinned but damped flux-line lattice, v is proportional to the current density J , and one obtains a linear relationship between V and I (Bardeen and Stephen, 1965). When pinning is important, the average velocity associated with the thermally activated jumps of flux lines is $v=v_0\exp[-U(J)/kT]$, where the prefactor v_0 may also be a function of J (see Sec. IV.A)—so E is also exponentially dependent on $U(J)$. Thus flux creep is generally associated with a highly nonlinear V - I relationship, dictated by the specific dependence of U on J and by the exponential dependence of E on U .

An example of this can be found in the standard flux creep model of Anderson and Kim, which assumes a linear dependence of the effective energy barrier on the current density, namely $U=U_0(1-J/J_c)$, giving rise to an exponential V - I curve. The logarithmic model of Zeldov *et al.* (1989, 1990) assumes $U=U_0\ln(J_c/J)$, resulting in a power-law V - I dependence. Other models will also be discussed in Sec. IV. It should be noted that the

V - I (or E - J) curves can be generated from magnetization measurements using the relations $M\propto J$ and $E\propto dM/dt$ (see Sec. III.B.6), and vice versa.

Flux creep manifests itself in many other ways. It causes the irreversible part of the magnetic torque to be time dependent (Giovannella *et al.*, 1987; Hergt *et al.*, 1993) and influences the ac magnetic response (Nikolo and Goldfarb 1989; Fleisher *et al.*, 1993; Sun *et al.*, 1993), its harmonics (Shaulov and Dorman, 1988; Shaulov *et al.*, 1990; van der Beek *et al.* 1994), and its frequency dependence (Malozemoff, Worthington, Yeshurun, *et al.*, 1988; Wolfus *et al.*, 1994; Prozorov *et al.*, 1995a). It couples to acoustic and optical signals (Zeldov *et al.*, 1989) and causes electrical and magnetic noise (Clem 1981; Placais and Simon, 1989; Johnson *et al.*, 1990; Jung *et al.*, 1991). In the interests of narrowing the focus of this review to a reasonable level, we concentrate here on experiments that directly measure the magnetic moment, with a variety of dc magnetometers, susceptometers, Hall probes, and so forth.

Magnetic relaxation in high-temperature superconductors is important for several reasons. From a fundamental perspective, the sheer size of the effect, along with its complex dependence on temperature, field, and other parameters, call for explanation and have led to significant new levels of theoretical understanding, including the elaboration of a cooperative theory of thermal activation and development of the novel concept of quantum vortex creep. This understanding has also contributed to the broader understanding of the magnetic phase diagram and pinning mechanisms and to an improved determination of the thermodynamic properties of high-temperature superconductors.

From an applied point of view, the importance of magnetic relaxation lies in the fact that it modifies the current-voltage characteristics of high-temperature superconductors, determines the temperature and time dependence of the current density and dictates limits to the stability of high-temperature-superconductor devices such as persistent-mode magnets used in levitation or in magnetic resonance imaging. These properties are central to the successful commercialization of high-temperature-superconductor technology. Although magnetic relaxation has set limits on the possible applications of this technology, it has by no means ruled them out: once understood, the phenomenon can be dealt with, circumvented, and even exploited.

This review summarizes the many experiments on magnetic relaxation in high-temperature superconductors and relates them to a variety of phenomenological theories. It extends earlier reviews by Malozemoff (1989) and by Hagen and Griessen (1990) on this topic. It is complementary to the largely theoretical review by Blatter *et al.* (1994) and to the broader review of McHenry and Sutton (1994), which discusses the entire magnetic phase diagram of high-temperature superconductors.

The goal of this review is to acquaint the reader with the main concepts, experimental results and theoretical interpretations in this important field. In the extraordi-

nary excitement generated by the discovery of high-temperature superconductors, an incredible array of publications has been generated, often with more enthusiasm than scientific care. This poses a special problem to the reader attempting to grasp this field for the first time: it is hard to know where to turn to get reliable information. This is the special task of this review: to guide the reader to the most reliable and meaningful experimental studies, as well as to the most plausible theoretical interpretations of magnetic relaxation in high-temperature superconductors.

The review is organized as follows: In the next section, basic concepts of magnetic relaxation are described, including the role of flux lines, the nature of magnetic measurements, the critical-state model, and thermal activation. This section is intended for the nonspecialist; those already working in this field can skip it. Section III gives the central content of the review: an assessment of the most reliable experimental results, including a discussion of the many experimental pitfalls that have plagued this field. Section IV gives a more detailed description of the phenomenological theories and their comparison to experiment. For more discussion of these theories, the reader is referred to the review of Blatter *et al.* (1994). Conclusions are summarized in Sec. V.

II. BASIC CONCEPTS

A. Flux exclusion and flux lines in superconductors

To understand the context for magnetic relaxation in high-temperature superconductors, it is necessary to review the magnetic properties of superconductors. This review starts with the Meissner effect: Every bulk superconducting material can exclude externally applied magnetic fields, in a certain range of field strength, from its interior, except for a thin surface layer. This flux exclusion is, along with lossless electrical current flow, one of the defining characteristics of superconductivity.

The most fundamental magnetic measurement, such as that obtained using a vibrating sample magnetometer or a magnetometer based on a SQUID, detects the magnetic dipole moment \mathbf{m} , which is a volume integral of the cross product of the current density \mathbf{J} and the radius vector \mathbf{r} in the sample:

$$\mathbf{m} = \frac{1}{2} \int \mathbf{r} \times \mathbf{J}(\mathbf{r}) d^3\mathbf{r}. \quad (2.1)$$

Higher-order multipoles can also be defined and are sometimes measured, but here we deal predominantly with the dipole moment. Magnetization \mathbf{M} is defined as the dipole moment per unit volume, although, as will be explained below, the magnetization of a nonequilibrium superconductor is not a spatially uniform quantity, as it is for single-domain ferromagnetic materials.

In a superconductor exhibiting the Meissner effect, the magnetic moment is diamagnetic, that is, the current loops are of such a polarity as to create a magnetic field opposing and canceling out applied fields in the interior of the sample. For an ellipsoidally shaped sample up to a

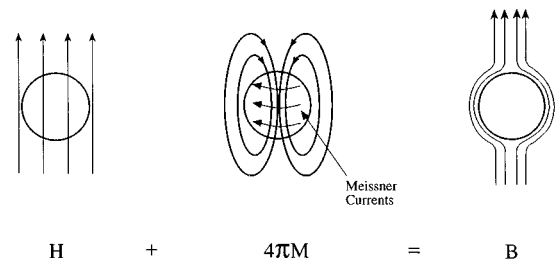


FIG. 3. In a superconductor exhibiting the Meissner effect, the shielding-current loops are of such a polarity as to create a magnetic field opposing and canceling out the applied field H in the interior of the sample. These lossless surface currents, the Meissner supercurrents, create a diamagnetic dipole moment. The magnetization M is defined as the dipole moment per unit volume. The internal field $B = H + 4\pi M$.

certain critical field, this cancellation is complete, except at the surface of the sample. These lossless surface currents are called Meissner supercurrents. This situation corresponds to the exclusion of magnetic flux from the interior, as illustrated in Fig. 3. It is opposite to the behavior of conventional ferromagnetic materials, which concentrate the magnetic flux inside the sample.

Fundamental to our topic is the fact that there are two characteristically different magnetic behaviors corresponding to two classes or types of superconductors. Type-I superconductors exhibit a Meissner effect up to a thermodynamic critical field value H_c at which field first penetrates into the bulk; above this field the material loses all superconductivity. By contrast, type-II superconductors, while exhibiting a Meissner effect up to a first critical field H_{c1} , have a further field range up to a second critical field H_{c2} at which superconductivity is fully quenched. In the intermediate field range, called a “mixed” or Shubnikov phase, magnetic flux can penetrate the bulk of the superconductor, but superconductivity is not fully quenched. (Note that the critical fields H_c and H_{c1} are defined with reference to a sample with zero demagnetization factor. For a sample with a demagnetization factor N , these fields become $H_c/(1-N)$ and $H_{c1}/(1-N)$, respectively.) The conventional magnetic phase diagram of type-II superconductors is shown in Fig. 4(a). High-temperature superconductors are all of type II, and the magnetic relaxation of interest for this review is a property of the mixed phase.

In the mid fifties, along with his theoretical prediction of the mixed phase, Abrikosov (1957) deduced the remarkable way in which magnetic field penetrates the bulk of a type-II superconductor, namely in tubes or cylinders called magnetic flux lines or vortices. Understanding flux lines is fundamental to understanding magnetic relaxation in superconductors. The internal structure of an isolated flux line is illustrated in Fig. 5: At large distances, the magnetic induction field strength \mathbf{B} falls off exponentially with radius from the center of the flux line, with a decay length equal to the London penetration depth λ , typically 1000–2000 Å along the CuO_2 planes at low temperatures. An integral of this field \mathbf{B} over a cross-sectional area of the flux line gives the total

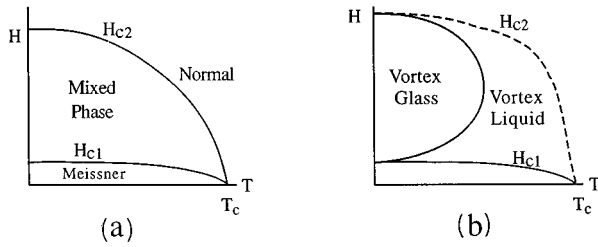


FIG. 4. Magnetic phase diagrams (schematic) for (a) a conventional superconductor, and (b) for high-temperature superconductor. Note that the figure is not drawn to scale, e.g., the reentrant liquid phase just above H_{c1} in (b) is highly exaggerated.

magnetic flux, which Abrikosov predicted would be precisely quantized in multiples of

$$\phi_0 = h/2e = 2 \times 10^{-7} \text{ G cm}^2, \quad (2.2)$$

which is called the magnetic flux quantum (given here in cgs units, with field in Gauss). Lossless supercurrents with density J circulate around the flux line in accordance with Ampere's law:

$$\nabla \times \mathbf{H} = (4\pi/c)\mathbf{J}. \quad (2.3)$$

Here, \mathbf{H} is the magnetic field, which is close to the induction field \mathbf{B} . (The difference is of order H_{c1} , which is small in high-temperature superconductors. Note, however, that this difference may be important at fields close to H_{c1} —see, for example, Vlasko-Vlasov *et al.*, 1994; Indenbom *et al.*, 1995; Zeldov *et al.*, 1995.) The circulating currents explain the alternate term “vortex” often given to flux lines; we will use the words “flux line” and “vortex” interchangeably to refer to the same linear object that involves both flux and circulating currents.

The flux line also has an inner core, of radius approximately equal to a superconducting coherence length ξ ,

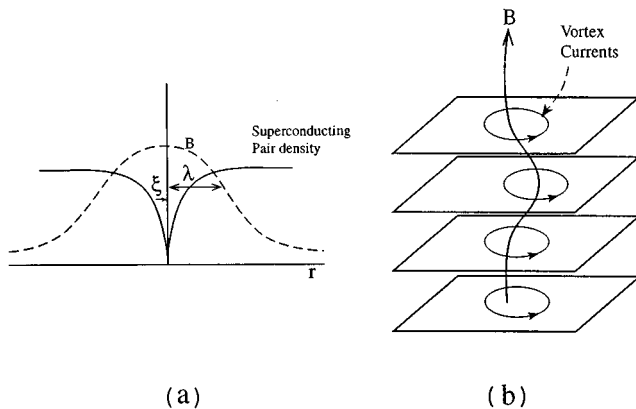


FIG. 5. Flux lines in superconductors: (a) Illustration of the internal structure of an isolated flux line. λ is the Meissner penetration depth and ξ is the coherence length. (b) Illustration of a flux line in a high-temperature superconductor composed of “pancakes” in the Cu-O planes.

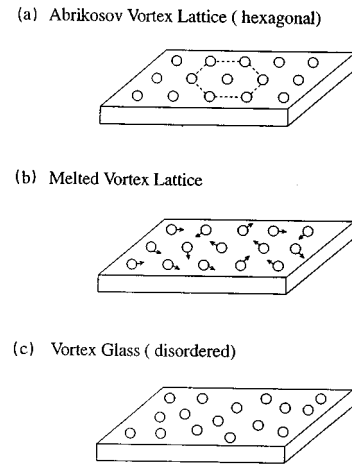


FIG. 6. Schematic description (top view) of (a) Vortices form an hexagonal lattice with a uniform equilibrium density. (b) Melted vortex lattice. (c) Vortex glass.

typically 10–20 Å in high-temperature superconductors at low temperatures, and in this inner core the superconducting order parameter (the superconducting electron-pair density) is suppressed, as shown in Fig. 5(a). The magnetic properties of superconductors, such as the dipole moment of Eq. (2.1), are derived from the superposition of the vortex currents around flux lines and the Meissner currents that circulate at the surface of the sample. A good recent review of basic properties of vortices in high-temperature superconductors was given by McHenry and Sutton (1994).

B. The vortex lattice

In the mixed phase of a perfect (no defects) ellipsoidally shaped isotropic superconductor, the vortices form a hexagonal lattice (see Fig. 6) with a uniform equilibrium density that depends on the external applied field H . In fact, the number density of the vortices times the flux quantum equals the locally averaged magnetic induction field B . Because the vortex density is uniform inside the sample, there is no bulk current flow on a scale much larger than λ . In this case, as in the case of single-domained ferromagnets, which have a uniform internal density of atomic spins, it is useful to define a magnetization or magnetic moment per unit volume, which can be calculated from the net surface current according to Eq. (2.1).

The classic theory of Abrikosov derives the magnetization of an isotropic superconductor as a function of field and temperature, based on the energy terms of a uniform-vortex lattice. This theory has also been extended to treat the highly anisotropic case appropriate to high-temperature superconductors (Kogan, 1981; Balatskii *et al.* 1986; Kogan and Campbell, 1989; Brandt, 1995). Their superconducting CuO_2 planes are only weakly coupled, and for vortices oriented along the c axis normal to the ab plane (the CuO_2 plane), the vortex structure has strong periodic variations along its length.

A suggestive analogy, proposed by Kes *et al.* (1990) and by Clem (1991), is that of a stack of pancakes, with each pancake representing the vortex within each CuO_2 plane, as suggested in Fig. 5(b). The pancakes are only weakly coupled along the vortex axis, and under appropriate excitation can slide past each other. The strength of the coupling determines the degree of anisotropy in the magnetic properties with respect to fields applied along the three principal axes. This anisotropy differs significantly for different families of materials, with YBCO having the lowest anisotropy and BSCCO having among the highest. Another interesting fact is that, for field applied in the ab plane, the vortex current flow pattern is approximately elliptical, with the major axis lying in the ab -plane.

Another aspect of the vortex physics that has emerged in recent years concerns thermodynamic fluctuations. As mentioned above, Abrikosov's theory was a purely classical, mean-field treatment that ignored fluctuations. Treatment of fluctuations in the behavior of vortex lattices started with the work of Kosterlitz and Thouless (1973) on two-dimensional superconducting films. However, the study of the unusual magnetic properties of high-temperature superconductors has spurred a significant extension of the statistical theory to include fluctuations in a three-dimensional anisotropic vortex lattice arising both from temperature and quantum effects.

Individual vortices can bend as a result of their attractive interaction with pinning centers. This attraction competes with the repulsive interaction of the vortices among themselves. Such concepts can be generalized by treating the vortex lattice as an anisotropic elastic medium. The various modes of distortion of such an elastic medium can be excited thermally. Such effects are particularly large in high-temperature superconductors, in part simply because the temperature can be much higher than in low-temperature superconductors, but also because the weak superconducting coupling between the planes allows thermal energy to pull the vortices apart relatively easily. A remarkable insight by Gammel *et al.* (1988) suggested that the vortex lattice melts at a temperature and field below the superconducting transition temperature T_c and the upper critical field H_{c2} . In this context, melting means that the long-range spatial order of the vortex lattice in the ab plane can break down and the vortices can move past each other as in a two dimensional fluid, as suggested schematically in Fig. 6(b). Subsequently, it was recognized that melting in the anisotropic vortex lattice of a high-temperature superconductor is more complicated—two modes of spatial disordering must be considered, even for the relatively simple case of the field being parallel to the c -axis: (1) disorder along the c axis and (2) disorder in the ab plane (Ryu *et al.*, 1992).

These concepts point to the fact that the H - T phase diagram of superconductors must be significantly generalized beyond the conventional one illustrated in Fig. 4(a). An example of the H - T phase diagram of this complex system is shown in Fig. 4(b), but the details depend on many parameters, including the strength of random

defects, to be discussed in the next section. A wide variety of experimental techniques has also uncovered unusual behavior often interpreted as vortex-lattice melting.

Remarkable though these concepts and experimental results may be, they are not the main topic of this review. Rather, we concentrate on other phenomena, equally interesting, which appear well below the melting transition line in the H - T plane. Indeed there is a relationship between the magnetic relaxation at lower temperature and fields and the melting at higher temperatures and fields: the former may be considered a precursor to the latter in the sense that the relaxation arises from the motion of vortices, which becomes increasingly rapid as temperature or field is raised.

C. Defected superconductors: irreversible magnetic properties

The vortex lattice and melting theory described above is complicated enough, but it still misses a major consideration necessary for any understanding and comparison to experiment. This is the fact that all real materials have defects which interact with the flux lines and pin them. As we shall see, the pinning process is a prerequisite for the phenomenon of magnetic relaxation, and it is now understood to have a strong effect on the phase diagram.

Defects in high-temperature superconductors can be of many types. These include the defects of as-grown material, including extended defects such as dislocations, twins, and stacking faults, and local defects, such as inclusions, oxygen nonstoichiometry, or other atomic defects. Other defects can be created by post-processing of the material, for example by mechanical working or by irradiation. Of particular interest are columns of damage created by high-energy ions, because these match the linear geometry of vortices. Often, the damage tracks can be several tens of Angstroms in radius, approximately matching the ab plane coherence length and thus facilitating pinning of the vortex core.

Typically, the vortex pinning energy per unit length of a defect in which superconducting pairing is locally destroyed is order of $H_c^2 \xi^2 / 8\pi$, where H_c is the thermodynamic field of the superconductor (equal to $H_{c2} \xi / \sqrt{2} \lambda$). Because of the low value of ξ , this energy is low, corresponding typically to a temperature of order 10–100 K; so thermal effects can evidently be significant. One effect of a random spatial arrangement of pinning centers is to disturb the simple hexagonal arrangement of the ideal Abrikosov vortex lattice. This was first recognized by Larkin (1970) and later by Larkin and Ovchinnikov (1979), who predicted that the lattice would break up into domains with a size corresponding to a finite correlation length for hexagonal order perpendicular to the vortex axis. The higher the pinning strength compared to the vortex interaction energy, the smaller the correlation length. This concept was extended to consider correlation lengths both parallel to the vortex axis and perpendicular to it.

The Larkin-Ovchinnikov state corresponds to a glassy arrangement of vortices because there is no long-range spatial order. However, below a melting temperature, the configuration can be temporally stable and a glassy order parameter can be defined (Fisher *et al.*, 1991). This state has become known as a “vortex glass.” The vortex-glass state is illustrated schematically in Fig. 6(c), and an example of an H - T phase diagram with a vortex-glass phase is shown in Fig. 4(b). The details here depend sensitively on the relative strengths of the vortex interactions, pinning strengths and configurations, and anisotropy of the material.

In the absence of pinning, the vortex array can always reach its thermodynamic equilibrium state, be it a lattice or a fluid. The magnetic moment or magnetization in this case is “reversible,” that is, it does not depend on the magnetic field or temperature history of the sample. The vortex-glass state that arises in the presence of pinning can also be a thermodynamic equilibrium state. However, there is another major physical phenomenon introduced by the presence of pinning, which is essential for our discussion of magnetic relaxation. This is the possibility of nonequilibrium or nonuniform spatial distribution of vortices, which are stabilized at least temporarily by the trapping of vortices in pinning wells. In this case the magnetic properties are “irreversible.” The relaxation of this nonequilibrium state back toward thermodynamic equilibrium (by the hopping or tunneling of vortices out of the pinning wells), is believed to be the basic phenomenon underlying the experimentally observed magnetic relaxation.

The possibility of a nonequilibrium arrangement of vortices is thus of fundamental importance from both a theoretical and practical perspective. In particular, a nonequilibrium gradient in vortex density corresponds to an overall gradient in field which, through Ampere’s law, must correspond to a bulk current. Thus the nonuniform pinning of flux lines inside a mixed-phase superconductor is the basis for bulk superconducting current flow. Although a high Meissner-current density can flow on the surface, it is limited to a layer of thickness λ , which is only a few thousand angstroms, and this is insufficient in most cases for practical application. Thus the bulk-current flow facilitated by pinning underlies bulk superconducting wire technology. It is no wonder that the history of optimizing critical current density in such wires centers on the effort to optimize defects for pinning flux lines.

The combination of the surface Meissner currents and the current vortices around the typically nonuniform spatial distributions of pinned flux lines creates a magnetic dipole moment (and higher multipole moments), which can be measured. In contrast to the simpler case of Meissner flux exclusion discussed above, the magnetic dipole moment can now be either diamagnetic (negative) or paramagnetic (positive), depending on the magnetic history of the sample, as will be explained below.

Further insight into the nonequilibrium behavior of pinned vortices comes from considering the force per unit volume exerted by the vortices on each other in a

density-gradient configuration. This force, derived by Friedel *et al.* (1963), has a form similar to the Lorentz force of electrodynamics, and is given by

$$\mathbf{F} = (1/c)\mathbf{J} \times \mathbf{B}, \quad (2.4)$$

where \mathbf{J} is the spatially averaged supercurrent density and \mathbf{B} is the spatially averaged induction field. Without a pinning force to counter this Lorentz force, the vortices will relax to their uniform equilibrium configuration, and bulk supercurrents will decay to zero. The current density at which the Lorentz force equals the maximum pinning force (in the absence of thermal activation or quantum tunneling) determines the critical current density J_c . The critical current density is a basic concept of superconductors out of equilibrium and is fundamental to bulk superconducting applications.

Thermal excitation or quantum tunneling can cause vortices to escape from their pinning wells and move in the direction of the Lorentz force, thus lowering the current density below J_c . The escape of vortices from pinning wells is the dominant mechanism of magnetic relaxation in superconductors. Thus magnetic relaxation and the reduction of current density are closely linked.

One of the most confusing aspects of the present terminology in the high-temperature-superconductor field is the use of the word “critical” to describe the usual measurements of current density. There is nothing “critical” about this reduced current density, because the voltage rises smoothly, albeit rapidly, with J . It has been necessary, therefore, to select a voltage criterion to define a current density. In this review, we will be careful to distinguish between the true critical current density, a rather inaccessible theoretical construct, and the measured current density which depends on the voltage criterion.

As indicated above, the magnetic moment of any superconductor has two components, one reversible and the other irreversible. The irreversible component, which arises from the nonequilibrium supercurrents, depends on the magnetic and thermal history of the sample. It is this component that undergoes magnetic relaxation. It is also worth emphasizing at this point that a superconductor with an irreversible magnetic moment has a nonuniform distribution of magnetic flux in its interior. Thus the conventional concept of magnetization as magnetic moment per unit volume must be treated with special care: while one can always divide the measured magnetic dipole moment by the sample volume to obtain a “magnetization,” this quantity must not be interpreted as an intrinsic or uniform volume property. In fact, as we shall see in the next section, the magnetization of a nonequilibrium superconductor grows with sample dimension.

D. The critical-state model

To quantify the behavior of pinned superconductors, a relationship must be established between the measured irreversible magnetization M_{irr} and the bulk supercurrent density J . This has been done via the critical-

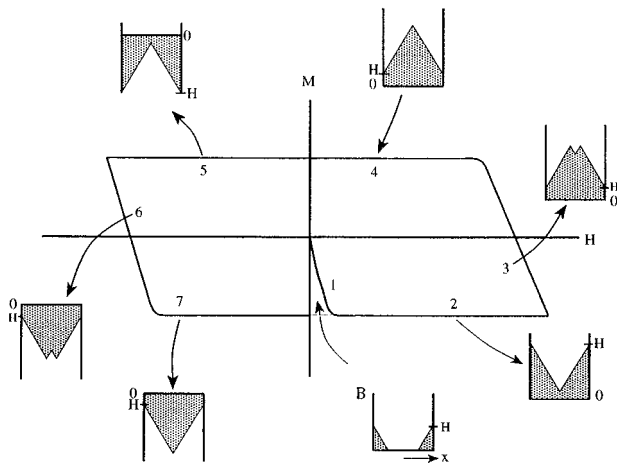


FIG. 7. Illustration of the hysteresis loop for a slab of thickness L , subjected from both sides to an applied field H parallel to its plane, with flux profiles through the slab thickness shown at various points around the loop. The shaded areas represent regions of flux penetration.

state model, first introduced by Bean (1962, 1964). This model is based on two simple assumptions, namely that the supercurrent density is given by a critical current density J_c , and that any changes in the flux distribution are introduced at the sample surface.

In the presence of large magnetic relaxation, the underlying concept of the critical-state model has had to change, but remarkably, the basic structure of the theory has been preserved. Now, instead of a true critical current density, one can use a current density determined by the time scale or electric-field level of the experiment. We proceed here with the conventional description of the theory in terms of J_c .

The specific relationship of M_{irr} and J_c depends on the sample geometry—an extensive review by Campbell and Evetts (1972) summarizes many theoretical results for specific configurations. Here, for illustrative purposes, we describe the simplest case, that of a slab of thickness L , subjected from both sides to an applied field H parallel to its plane. We start from the zero-field state and proceed to consider the hysteresis loop as the field is oscillated through a full cycle, positive and negative, back to zero. We ignore the reversible magnetization and consider J_c to be independent of B . The resulting hysteresis loop is illustrated in Fig. 7, along with flux profiles through the slab thickness at various points around the loop.

Initially, flux penetrates from the surfaces, with a gradient $(4\pi/c)J$ given by Eq. (2.3). The flux fronts, sketched at point 1 in Fig. 7, can readily be calculated to penetrate a distance $cH/4\pi J_c$ into the slab. The magnetization, given by $4\pi M = B - H$ is now

$$4\pi M = (cH^2/4\pi J_c L) - H. \quad (2.5)$$

When the flux fronts reach the center, at $H = 2\pi J_c L/c$, the magnetization reaches its maximum diamagnetic value of

$$M = -J_c L/4c. \quad (2.6)$$

This is a particularly famous result of the Bean critical-state model, showing that the magnetization increases linearly with sample dimension L . Similar results to Eqs. (2.5) and (2.6) can be derived for different points around the hysteresis loop. A remarkable prediction is that the magnetization changes sign on the downward part of the field cycle and becomes paramagnetic (positive). Thus both diamagnetic (flux excluding) and paramagnetic (flux concentrating) behaviors are possible with a non-equilibrium superconductor.

When the field is reduced back to zero after an initial increase to a positive value, flux is trapped inside the sample, giving rise to what is often called “remanent” magnetization M_{rem} . More precisely, this is called “isothermal remanent magnetization” because the process is carried out at constant temperature. A somewhat different kind of remanent magnetization arises when a sample is cooled from above T_c in a magnetic field to a fixed temperature (known as “field-cooling”) at which point field is turned off. This is known as “thermoremanent” magnetization. Expressions for thermoremanent magnetization can be derived in a similar way to those for isothermal remanent magnetization, with some assumptions about the temperature dependence of the parameters.

Magnetic relaxation will occur from any of the non-equilibrium flux states described above. Because of the convenience of making measurements in zero applied field, many studies have used samples with thermoremanent or isothermal remanent magnetization as starting states for magnetic-relaxation studies. However, as will be described in Sec. III, there are severe problems in interpreting these data quantitatively. The most preferred conditions are those on the diamagnetic or paramagnetic plateaus of the hysteresis loop at sufficiently high field to minimize demagnetizing effects.

E. Magnetic relaxation: basic mechanism

Now we are in a position to understand the basic concepts of magnetic relaxation in superconductors. Any process that allows the nonequilibrium configuration of vortices to relax will lead to a redistribution of current loops in the superconductor and hence to a change in the magnetic moment with time. Thus the measured magnetic relaxation can be thought of as being caused by the spontaneous motion of flux lines out of their pinning sites. Such motion usually arises from thermal activation, but it can also arise from quantum tunneling or other external activation, such as mechanical vibrations.

The concept of thermal activation causing hopping of vortices or bundles of vortices out of their pinning-potential wells was first suggested by Anderson (1962) to explain the data of Kim *et al.* (1962) on the relaxation of persistent currents in NbZr tubes. In its simplest form, the idea can be presented as follows. According to the conventional Arrhenius relation, a hopping time t is given in terms of the potential-energy barrier height U , the Boltzmann constant k , and the temperature T :

$$t = t_0 \exp(U/kT). \quad (2.7)$$

The preexponent t_0 (referred to as the “effective” hopping attempt time) can differ from the microscopic attempt time by orders of magnitude, as becomes clear in the more complete theory described in Sec. IV.A. The hopping process is assisted by the driving force $\mathbf{F} = (1/c)\mathbf{J} \times \mathbf{B}$ —see Eq. (2.4). Therefore, U should be a decreasing function of J . In a first approximation, the net barrier is reduced linearly with the current J , according to

$$U = U_0[1 - J/J_{c0}], \quad (2.8)$$

where U_0 is the barrier height in the absence of a driving force, and J_{c0} corresponds to the critical current density (in the absence of thermal activation) required to tilt the barrier to zero in this approximation. Combining Eqs. (2.7) and (2.8) and solving for J , we obtain the classic equation of flux creep

$$J = J_{c0} \left[1 - \frac{kT}{U_0} \ln\left(\frac{t}{t_0}\right) \right]. \quad (2.9)$$

This derivation is highly simplified and not fully correct; a more proper treatment is given in Sec. IV.

Since the magnetization M is proportional to J according to Eq. (2.6), we can immediately see some of the main features of flux creep: namely, the magnetization is expected to decay logarithmically in time and to drop with temperature. The range of validity of Eq. (2.9) will be discussed further in Sec. IV.

Equation (2.9) is often referred to—and we will do so here—as the Anderson-Kim equation for flux creep. Perhaps unfairly to these original authors, who had a broader concept of the general theory, their names are usually used in connection with the version of the theory in which the $U(J)$ dependence is assumed to be linear [Eq. (2.8)]. Later versions of the theory, starting with Beasley *et al.* (1969), considered nonlinear $U(J)$ dependencies. The full impact of these nonlinearities became apparent only recently, after the development of the vortex-glass (Fisher, 1989) and collective-creep (Feigel'man *et al.*, 1989, 1991) theories. The result, in these theories, that the barrier diverges as the current J approaches zero is connected with a novel kind of a phase transition in type-II superconductors (see, for example, Fisher, 1989; Fisher *et al.*, 1991; Huse *et al.*, 1992. See also the discussion of the vortex-glass state in Sec. II.C above).

According to Eq. (2.9), the relaxation rate depends on the limiting critical current density J_{c0} , the barrier height U_0 , and the effective attempt time t_0 . To eliminate one or more of these parameters it is convenient to evaluate a normalized relaxation rate $S \equiv (dM_{\text{irr}}/d \ln t)/M_{\text{irr}}$ corresponding to a logarithmic derivative of magnetization versus time. The normalized relaxation rate can be derived directly from Eq. (2.9):

$$S \equiv \frac{1}{M_{\text{irr}}} \frac{dM_{\text{irr}}}{d \ln t} = \frac{d \ln M_{\text{irr}}}{d \ln t} = -\frac{kT}{U_0}. \quad (2.10)$$

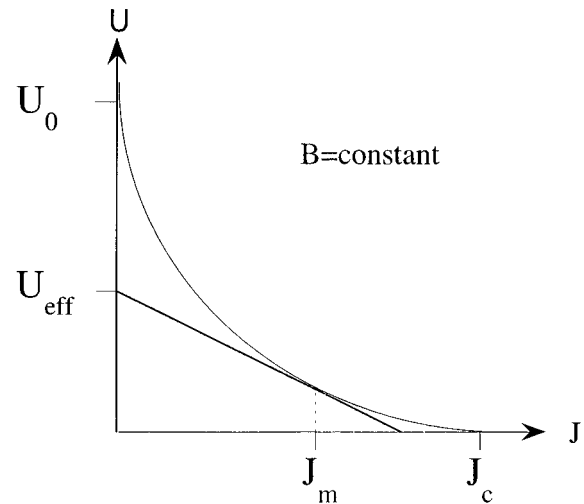


FIG. 8. A schematic illustration of a nonlinear functional form of $U(J)$ at constant magnetic field B . Using the linear approximation [Eq. (2.8)] to determine the pinning barrier yields an apparent value U_{eff} which is smaller than the true pinning potential U_0 . U_{eff} corresponds to the U axis intercept of the straight line tangent to $U(J)$ at the measured current J_m .

Thus in the Anderson-Kim theory, a measurement of the normalized relaxation rate determines the pinning barrier U_0 . This simple result provides one of the major motivations for studying magnetic relaxation in type-II superconductors, namely to determine the energy barrier, and thereby something about the pinning mechanism itself.

However, the nonlinearity in $U(J)$ mentioned above complicates the analysis. To illustrate these complications, we show in Fig. 8, in a schematic way, a nonlinear functional form of $U(J)$. In this figure J_m indicates the experimentally measured current, U_0 is the “true” activation barrier at $J=0$, and U_{eff} is the “effective” activation barrier, which corresponds to the U -axis intercept of the straight line tangent to $U(J)$ at J_m , resulting from the use of Eq. (2.8). This yields an apparent value U_{eff} that is smaller than the “true” pinning potential U_0 at $J=0$. Moreover, this value of U_{eff} may strongly vary with J , depending on the curvature of $U(J)$. This means that different values of U_{eff} will be measured at different times. This problem is particularly severe for high-temperature superconductors because of the large changes in J with time. A straightforward application of Eq. (2.10) is complicated further by the possibility of multiple contributing mechanisms, including distributions of barrier heights (Malozemoff, Worthington, Yandofski, *et al.*, 1988; Griessen, 1990; Gurevich, 1990; Martin and Hebard, 1991; Niel and Evetts, 1991) and surface barriers (Petukov and Chechetkin, 1974; Koshelev, 1991, and 1994; Burlachkov, 1993a).

The sheer size of the magnetic relaxation in high-temperature superconductors made it initially not so clear that thermally activated flux creep was the relevant mechanism; the effect had always been much smaller in the conventional low-temperature superconductors. Early in the development of the high-temperature-

superconductor field, an alternative picture based on a model of disordered Josephson links was proposed (Deutscher and Müller, 1987; Morgenstern *et al.*, 1987; Müller *et al.*, 1987, Aksenov and Sergeenkov, 1988). This model also predicted large, nonexponential magnetic relaxation. The theory was applied not only to ceramic, polycrystalline or granular materials, where there is convincing evidence for Josephson links at the grain boundaries (Chaudhari *et al.*, 1988; Mannhart *et al.*, 1988), but also to crystals. In the latter case there has been much more controversy about the origin, or even the existence, of weak links; twin boundaries and networks of oxygen defects (Däumling *et al.*, 1990) have been proposed. Nevertheless, in crystals magnetic relaxation is now usually interpreted in terms of the flux-creep model.

Returning to the simpler Anderson-Kim flux-creep theory, we can ask why the magnetic relaxation should be particularly large in high-temperature superconductors. As was first suggested by Yeshurun and Malozemoff (1988) and by Dew-Hughes (1988), the large size of the effect can be interpreted in terms of two factors, the temperature and the barrier height. Assuming pinning of the normal vortex line, which has a radius of the order of the superconducting coherence length ξ , the barrier height can be estimated by the condensation energy $(H_c^2/8\pi)\xi^2$ per unit length of the vortex. The small value of the coherence length ξ in high-temperature superconductors leads to a relatively small U_0 in Eq. (2.9). Simultaneously, the high critical temperature of these superconductors increases the numerator in the kT/U_0 factor of this equation. The combined effect increases the importance of the time-dependent term in Eq. (2.9), giving rise to “giant flux creep” (Yeshurun and Malozemoff, 1988).

The basic Anderson-Kim model outlined above explains the large size of the effect. However, as the field has progressed, many new and unanticipated aspects, both theoretical and experimental, have emerged. For example, Eq. (2.10) predicts a linear dependence of the normalized relaxation rate S on temperature, with $S=0$ at $T=0$. Experimentally, however, $S(T)$ neither depends linearly on temperature (Civale *et al.*, 1990), nor does it extrapolate to zero at $T=0$ (Mota *et al.*, 1988; Hamzic *et al.*, 1990). We elaborate on these, and other problems, in Sec. III. The questions raised by the experimental results led to modifications of the basic Anderson-Kim model, as well as to the development of more sophisticated models, to be described in Sec. IV.

III. OVERVIEW OF MAGNETIC-RELAXATION EXPERIMENTS

Magnetic relaxation in high-temperature superconductors has been the subject of numerous experimental studies. Before we summarize the trends that emerge from these studies, we describe, in Sec. III.A, a number of experimental pitfalls, which unfortunately make the conclusions drawn in a large number of articles unreliable. Perceptive screening is required, therefore, to use

fully penetrate the literature, and we attempt such a screening in this review. We hasten to add that we must include some of our own earlier works in this “unreliable” category.

As we already mentioned in the Introduction, magnetic relaxation has been studied in all high-temperature-superconductor families. These include LaSrCaCuO (Mota *et al.*, 1987, 1988), YBaCuO (Hor *et al.*, 1987; Hagen *et al.*, 1988; Tuominen *et al.*, 1988; Wong *et al.*, 1988; Yeshurun and Malozemoff, 1988; Yeshurun *et al.*, 1988a), BiSrCaCuO (Biggs *et al.*, 1989; Lessure *et al.*, 1989; Safar *et al.*, 1989; Yeshurun, Malozemoff, Worthington *et al.*, 1989; Shi, Xu, Umezawa, and Fox, 1990) TlBaCaCuO (McHenry, Maley, Venturini, and Ginley, 1989; Chan and Liou, 1992; Schilling *et al.*, 1992), HgBaCaCuO (Kim and Kim, 1996), the electron-doped superconductor NdCeBaCuO (Fabrega *et al.*, 1992; Yoo and McCallum, 1993), and even the non-copper-based superconductor BaKBiO (McHenry, Maley, Kwei, and Thompson, 1989). All of these studies—as well as hundreds of other publications not listed here—reveal large relaxation rates $dM/d \ln t$. Usually, the YBaCuO (YBCO) samples are somewhat more stable to flux creep than their BiSrCaCuO (BSCCO) and TlBaCaCuO (TBCCO) analogs. Similarly, films are usually more stable than crystals and ceramics of the same family. In addition to the size of the effect, the high-temperature-superconductor families exhibit other common characteristics of magnetic relaxation, summarized in Sec. III.B. In Sec. III.C we describe various experimental approaches to reducing the magnitude of the relaxation effect.

A. Experimental issues in magnetic-relaxation measurements

In this section we review a series of problems that, in retrospect, can be seen to compromise the validity of a large number of early and even present-day studies. This situation is not unusual in any new field, but it is particularly severe in high-temperature superconductivity, where the level of enthusiasm and the rush to publish were extreme in the early years after the discoveries of Bednorz and Müller (1986) and Wu *et al.* (1987).

In essence, the early papers on magnetic relaxation in high-temperature superconductors correctly revealed the impressive size of the effect and identified some valid trends with respect to temperature and field. However, for any quantitative discussion, the reader, especially any theorist seeking to interpret data in detail, is cautioned to look into more recent studies. Experimentalists entering or continuing their work in this field are also encouraged to take note of the experimental issues summarized below.

1. Sample inhomogeneities

The most obvious issue, especially important in the early years, has been the quality of the high-temperature-superconductor samples available for study. The problems in producing homogeneous samples

of such multicomponent compositions and the difficulties with uniform oxygen stoichiometry are well documented (see, for example, Beyers and Shaw, 1989). The magnetic signal, which is an integral over all the circulating current paths in the material, may exhibit peculiar relaxation behavior due to different relaxation rates in the different regions. Clearly, an interpretation of magnetic relaxation data requires a well-defined and regular vortex distribution. A variety of novel spatial-scanning techniques have been developed to map the effect of inhomogeneities on the magnetization, such as scanning micro-Hall probes (Frankel, 1979; Bending *et al.*, 1990; Konczykowski *et al.*, 1990; Lim *et al.*, 1991; Tamegai *et al.*, 1992; Brawner *et al.*, 1993a, 1993b), Hall arrays (Zeldov, Majer, Konczykowski, Larkin, *et al.*, 1994), and magneto-optic imaging (Alers, 1956; Kirchner, 1968; Huebener, 1979; Moser *et al.*, 1989, 1990; Szymczak *et al.*, 1990; Baczewski *et al.*, 1991; Ludescher *et al.*, 1991; Schuster *et al.*, 1991; Dorosinskii *et al.*, 1993; Forkl *et al.*, 1993). These have often revealed unexpected irregularities in the vortex-penetration profile, invalidating simple Bean-model interpretations. It should be noted that scanning of a single Hall probe has the drawback that a flux profile creeps as one scans the sample. Hall-sensor arrays that allow simultaneous electronic scanning of the array elements are definitely preferable. As yet, Hall-probe arrays are limited to sequential measurements along one axis only, whereas the magneto-optic technique provides an instantaneous picture of the whole sample surface. However, the Hall-probe array technique is progressing rapidly and two-dimensional arrays are being developed.

As a crude generalization, the cleanest materials have been small YBCO crystals, although even here a debate still rages about the degree of oxygenation in the interior and the possibility of networks of oxygen defects (Däumling *et al.*, 1990). More recently, a great deal of work has been directed to perfecting BSCCO crystals. Much of this review focuses on these two materials. YBCO films have also reached a high level of quality. TI-based crystals and films of good quality have recently become available, and one can expect much progress in this family of materials in the near future.

Another issue concerns the possibility of multiple mechanisms of pinning, even in a pure material. A dramatic example of this has recently emerged in studies of both YBCO and BSCCO crystals. The problem concerns effects of surface pinning or surface barriers, which become dominant at higher temperatures. We elaborate on this topic in Sec. III.B.4 below. Another good example is the presence in YBCO of twin networks that prevent flux from penetrating or leaving part of the sample.

2. Field inhomogeneities in SQUID magnetometers

A second problem comes from the widespread use of SQUID magnetometers in the study of magnetic relaxation in high-temperature superconductors. While their sensitivity to tiny magnetic signals has made them a

popular replacement for more conventional measurement devices, such as vibrating sample magnetometers, certain precautions should be taken in using these instruments in measurements of hysteretic superconductors (Grover *et al.*, 1991; Suenaga *et al.*, 1992).

SQUID magnetometers are usually operated in a mode wherein the sample is translated multiple times between a set of superconducting pickup coils wrapped around a cylindrical sample chamber. Wired in a gradiometer configuration, these coils are spaced by as much as eight centimeters along the chamber in some widely used commercial instruments. At the same time, a magnetic field, usually generated by a superconducting magnet, bathes this measurement region. The problem arises from the fact that this magnetic field is not perfectly uniform; over the usual range of sample translation, it may vary by a few percent. For measurements of reversible magnetic phenomena like the linear magnetic susceptibility, this is no problem; the resulting signal reflects some average magnetic field experienced by the moving sample. But for hysteretic magnetic materials, including all the irreversible type-II superconductors, translation of the sample up and down through the non-uniform magnetic field causes it to experience a minor loop of its hysteresis cycle. With such multiple loops, the magnetic state may be cycled towards the reversible limit. The effect is analogous to the demagnetization that is commonly performed with ac fields to demagnetize tape recorder heads. Thus instead of a correction of a few percent, suggested by the size of the magnetic-field nonuniformity, a hysteretic signal from a persistent irreversible supercurrent can disappear entirely!

There have been a number of reports in the literature that note a dramatic discrepancy between transport measurements showing a strong critical current and magnetic hysteresis measurements showing reversible behavior and no critical current, under identical conditions of temperature and magnetic field. A valid reason for this discrepancy is that, in magnetic measurements, when the creep is large, the measured irreversible moment may be small even if J_c is large, while in transport measurements the current is directly applied and J_c can be attained. In view of the discussion in the preceding paragraph, we note that the magnetic-field inhomogeneity problem may also contribute to the discrepancy between the magnetic and transport data. Thus it is reasonable to question the magnetic data, especially where authors do not indicate a clear recognition of the above-mentioned problem.

The usual solution to the field-inhomogeneity problem is to reduce the amplitude of sample translation to only 1–2 cm, in order to stay within a more uniform central section of the superconducting magnet of the SQUID magnetometer. Unfortunately, this cuts down the measurement sensitivity and accuracy. Simple tests as a function of translation amplitude can confirm whether a stable and meaningful measurement has been achieved. It should be noted that there is a significant range of temperatures and fields in which high-temperature superconductors show reversible supercon-

ducting behavior. This is the range between the so-called “irreversibility line” in the field-temperature plane, and the curve describing the upper critical field. In this region, as well as in the region above T_c , full sample translation amplitudes can be restored.

3. Establishment of a fully penetrated flux distribution

Most magnetic relaxation experiments have been interpreted according to the simple assumption that the measured irreversible magnetization M_{irr} is proportional to the persistent current density J . For example, according to the Bean relationship, Eq. (2.6), $M_{\text{irr}} = -JL/4c$ for an infinite slab of thickness L lying parallel to the magnetic field. Here, J is a current that circulates in the plane of the slab, perpendicular to the magnetic field, and in opposite directions on the two sides. However, this relationship is valid only when the field penetrates the whole sample. When a type-II superconductor is first magnetized, flux initially penetrates from the surfaces with oppositely circulating currents on the two surfaces, and until the flux fully penetrates the slab, the relationship between J and M_{irr} is very different; in fact, the inverse- JL relationship $4\pi M_{\text{irr}} = (cH^2/4\pi JL) - H$ holds (H is the applied field). Even greater complexities arise when the field is first increased and then reduced; in a certain range of fields, zones of clockwise and counterclockwise persistent currents coexist. These obviously complicate the interpretation of the magnetic-relaxation data.

In fact, the much-cited early paper by the present authors (Yeshurun and Malozemoff, 1988) on magnetic relaxation in YBCO crystals overlooked this problem, although a note of correction was inserted later in press (Yeshurun *et al.*, 1988a). One solution to this problem is to confirm, most simply via the shape of the magnetic hysteresis loop, that full penetration has occurred before relaxation measurements are initiated. As a rule of thumb, we suggest that full penetration is established for fields that are approximately 1.5 times the field corresponding to the minimum magnetization in the virgin magnetization curve (Wolfus *et al.*, 1989). An alternative solution, used by us among others (Yeshurun *et al.*, 1988a) is to apply the version of the Bean equation that was derived specifically for the partly penetrated case.

4. Complex demagnetizing fields and anisotropy of high-temperature-superconductor crystal platelets

Given the complexities of high-temperature-superconductor powders, the majority of magnetic relaxation studies have by now been done on crystals, with the hope of having a better characterized and more uniform material. Unfortunately, these crystals almost always come in an anisotropic platelet shape and, in addition, their superconducting properties are highly anisotropic. This has led to several problems. On the one hand, measurements with magnetic field parallel to the platelet plane (the principally conducting ab plane of the crystal) have been found to be extremely sensitive to the precise field orientation (Cronmeyer *et al.*, 1990).

Oriental accuracy of tenths of a degree is required even for YBCO, one of the least anisotropic materials. This sensitivity can be seen to arise in part simply from the shape anisotropy of the sample. Current loops circulating perpendicular to the a or b axes in the ac or bc planes circumscribe a small area limited by the thickness of the crystal, while current loops circulating perpendicular to the c axis, in the ab plane, circumscribe a much larger area. Since the magnetic moment is proportional to the product of current and cross-sectional area, the component parallel to the c axis easily dominates with even a slight misorientation. With relaxation rates likely to differ for the different orientations, it becomes very difficult to deconvolute the different contributions. Any study of relaxation with fields in the ab orientation must make explicit mention of the alignment accuracy.

A complementary problem occurs for measurements with magnetic field perpendicular to the platelet plane, which is the usual case reported in the literature. In the low-field range, where the flux densities \mathbf{B} due to circulating currents in the sample are non-negligible compared to the applied magnetic field, the platelet geometry leads to strong demagnetizing fields. These fields can be visualized as “wrapping around” the platelet and so actually lie parallel to much of the platelet surface (for recent references see Brandt and Indenbom, 1993, and Zeldov, Clem *et al.*, 1994). Thus rather than having a clean configuration with field \mathbf{H} and flux \mathbf{B} parallel to the c axis, the field and flux directions have a complex configuration within the sample. In addition, during relaxation the induction at the sample surface changes. This effect may be substantial if the magnetic moment is large. Hence there exist few measurements taken in strictly constant \mathbf{B} .

The field profiles have been calculated, assuming a hypothetical J_c anisotropy (Däumling and Larbalestier, 1989; Conner and Malozemoff, 1991; Conner *et al.*, 1991). However, considering the complexity of these profiles and the intermixing of different anisotropy components, it seems hopeless to try to extract useful quantitative information about the underlying relaxation processes from such a configuration. An experimental approach to bypass these problems is to apply a sufficiently high magnetic field such that the demagnetizing fields are negligible in comparison and the relative changes in \mathbf{B} are small. In this case, the flux lines \mathbf{B} lie essentially along the field direction \mathbf{H} . The size of the field required to overcome the demagnetization problem can usually be determined empirically from an examination of the hysteresis loop, most conveniently from the field range required to reverse the loop and reestablish the high-field hysteretic plateau usually observed in these materials. Of course, this solution limits the range of field and flux density that can be explored. It should be noted here that in thin films, the volume and hence the magnetic moment is very small, so that the demagnetization corrections are less important. Also, in some cases, it can be shown directly that the relative changes in the magnitude of \mathbf{B} are not important (van der Beek, Kes, *et al.*, 1992).

An alternative solution is to synthesize samples with a less platelet-like shape, or even to artificially stack crystals to reduce their demagnetizing fields. Recently this has been achieved, in a tour de force, with multiple BSCCO crystals, and indeed a noticeable change in shape of the hysteresis loop was confirmed (Kishio *et al.*, 1991), as predicted from calculations (Conner and Malozemoff, 1991). Unfortunately, while this approach can give an approximate idea of the hysteresis loop shape without demagnetization distortions, it fails to provide a clean geometry for magnetic relaxation for yet other reasons, which are discussed below.

5. Complexities of the remanent state

Additional problems afflict studies of magnetic relaxation in the remanent state, especially those seeking to establish the precise functional form of the relaxation. In general, the pinning strength—and hence the critical current density—depends on the flux density, and this is particularly true in the low-flux-density limit. Thus in the remanent state, or at low applied magnetic fields H (compared to the maximum flux density B), any magnetic-relaxation measurement integrates the contributions of regions with significant variations in properties. Without some detailed microscopic, spatially resolved probe of magnetic relaxation, it is impossible to deconvolute these contributions.

There is a yet more fundamental problem with remanent-state magnetic measurements. This goes back to the basic equation for flux-line current flow [Eq. (4.2) below]. The point is that the flux-line current is proportional to the hopping rate, but it is also proportional to the density of flux to start with. A high hopping rate is not effective if there is no flux present! And so the low-flux-density region can gate the flux creep. To make matters worse, in the remanent state, the lowest flux density is at the surface, where material properties may differ from the bulk.

The theoretical solution of the flux-creep diffusion equation [Eq. (4.2)] is straightforward if B changes linearly from the edge to the center. This requirement is fulfilled when the variations of B across the sample are relatively small, i.e., $\Delta B/B \ll 1$. Experimentally, this condition can easily be achieved by using sufficiently large external fields. However, it is difficult to assure this condition in the remanent state (Hagen and Griessen, 1990). Nonlinear profiles may cause deviations from time-logarithmic relaxation. Therefore, any efforts to study such deviations in a remanent configuration for the purpose of comparison to more detailed theories must be viewed with skepticism.

Interestingly, understanding of magnetic relaxation in the remanent state are less problematic in flat samples where the demagnetization factor is significant. In a flat sample, the sample contribution to B is composed of components parallel (B_{\parallel}) and perpendicular (B_{\perp}) to the applied field. As a result of demagnetization, B_{\parallel} has a different sign at the edge and at the center, and there exists a contour inside the sample where $B_{\parallel} = H$ exactly. When flux creep is present, and the induction at the

sample center increases, B_{\parallel} at the sample edge decreases. On the contour of $B_{\parallel} = H$, the component B_{\parallel} does not change with time. In the remanent state, this contour corresponds to that where $B_{\parallel} = 0$. Evidently, the relaxation rate there of B_{\parallel} is zero. However, the relaxation process continues by relaxation of the component B_{\perp} . Thus the magnetic moment in the remanent state can relax because of the nonzero demagnetizing factor.

6. Determination of the irreversible component of the magnetization

Magnetic relaxation comes just from the irreversible component of the magnetization M_{irr} in the superconductor, and thus one can determine the rate $dM_{\text{irr}}/d \ln t = dM/d \ln t$ without concern for background contributions to the total M from either the reversible magnetization or from background contributions such as sample holders and substrates. However, we have made clear with respect to Eq. (2.10) that knowing the absolute size of M_{irr} is important for evaluation of the pinning energy. Unfortunately, the correction of data for these contributions is rarely confirmed in the literature.

The irreversible component of the magnetization M_{irr} can be extracted from the hysteresis loop, examples of which are given in Fig. 9. This figure presents three typical hysteresis loops observed for YBCO samples. Figure 9(a) exhibits a symmetric loop (around $M=0$), similar to that predicted by the Bean model. In contrast, the hysteresis loop shown by Figs. 9(b) and 9(c) are asymmetric. The asymmetry in Fig. 9(b) is caused by a superposition of reversible and irreversible contributions. The asymmetry in Fig. 9(c) is caused by the presence of surface barriers, as discussed in Sec. III.B.4. In Figs. 9(a) and 9(b), half the difference between the plateau values at constant applied fields corresponds to M_{irr} , while the average, a slight offset from symmetry around $M=0$, corresponds to M_{rev} . However, this procedure is incorrect for Fig. 9(c), where Bean-Livingston surface barriers, not bulk pinning, dominate the irreversible features of the hysteresis loop. The signal for the presence of such surface barriers is a value of $M \approx 0$ on the field-decreasing branch of the hysteresis loop.

It should be mentioned that, because of the large creep, the measured initial value of the magnetization as well as the shape of the hysteresis loop may be affected by the sweep rate of the external field (Pust, 1990; Schnack *et al.*, 1992; Gurevich and Küpfer, 1993; Ji *et al.*, 1993; Jirsa *et al.*, 1993). In fact, the sweep-rate dependence of M_{irr} and the time dependence of the magnetic moment during relaxation hold the same type of information (Caplin *et al.*, 1995; Perkins *et al.*, 1995), and the two experiments can be cross-correlated to extract both J_c and the barrier height (Lairson *et al.*, 1991; Schnack *et al.*, 1993; Griessen, Wen, *et al.*, 1994; Wen *et al.*, 1995).

An example of a hysteresis loop in which the applied magnetic field has periodically been halted and then steadily increased is shown in Fig. 10. After decaying (the vertical segments) while the field is held constant, M rapidly retains its previous plateau as soon as the field

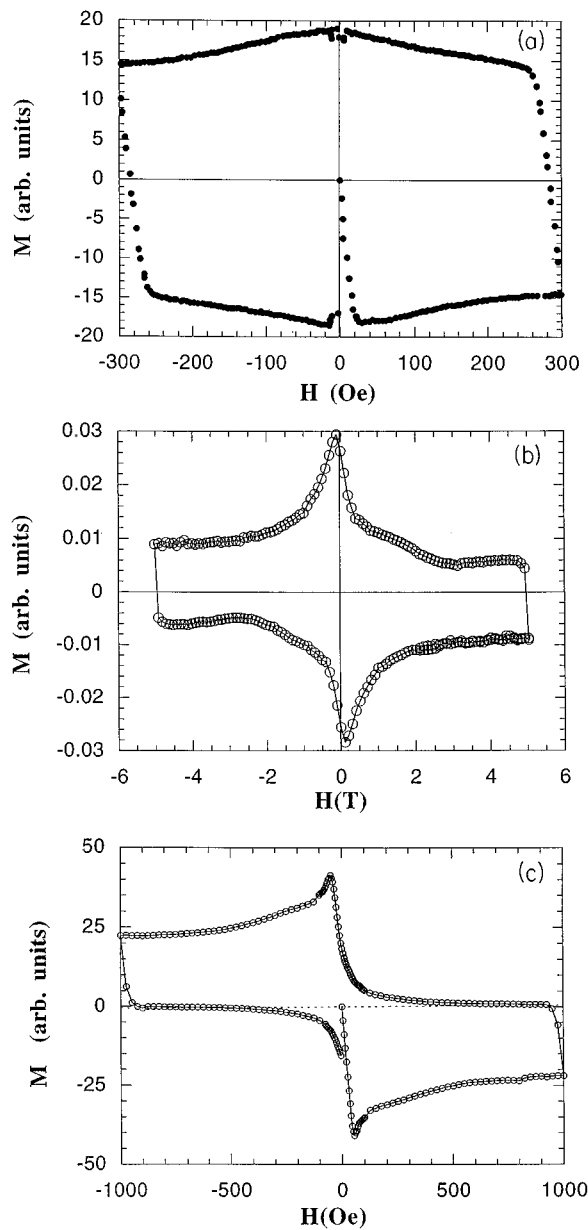


FIG. 9. Typical magnetization curves M vs H , measured in high-temperature superconductors. The irreversible magnetization is half the difference between the upper and the lower branches. (a) $M(H)$ for YBCO film at $T=80$ K. The reversible magnetization in this case is negligible. (After Yeshurun *et al.*, 1993.) (b) $M(H)$ for YBCO crystal at 84 K. Note the “asymmetry” in the values of the lower and upper branches of the magnetization caused by the contribution of the reversible magnetization. The reversible magnetization is the average of the upper and the lower branches. (c) $M(H)$ for a YBCO crystal exhibiting asymmetry due to surface barriers.

increase is resumed, and the plateau level depends on the rate of increase of the field. Unfortunately, many studies have assumed M_{irr} to be equal to the plateau value, ignoring the above mentioned rate dependence and the asymmetry around $M=0$, which arises from the reversible magnetization.

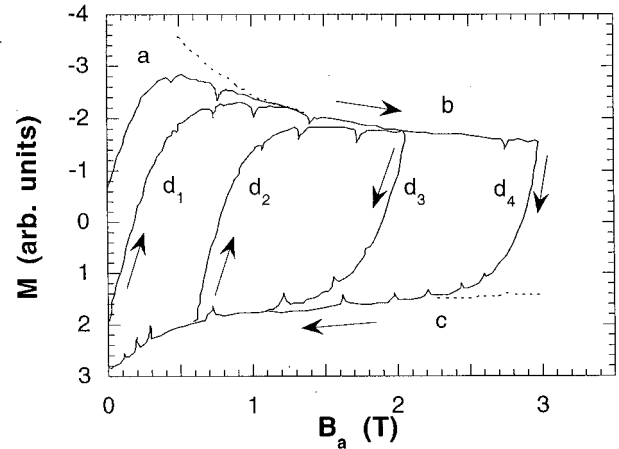


FIG. 10. Hysteretic magnetization curves in a sintered YBCO in which the applied magnetic field has periodically been halted. The vertical segments reflect flux creep. (After Kwasnitza and Widmer, 1991.)

7. Step-by-step procedure for measurements of magnetic relaxation

As a summary of this subsection we propose here one possible step-by-step procedure for measurement of magnetic relaxation that avoids the experimental pitfalls described above and facilitates comparison with theory. As an example we choose to describe measurements of flux exit in a high-temperature-superconductor crystal after a step down in the external field. One may proceed as follows:

(i) Ensure the quality of the sample, field homogeneity, and the orientation of the field relative to the crystal axes.

(ii) Measure the magnetic hysteresis loop at the temperature T for which relaxation will be measured, at a fixed sweep rate up and down. Report the sweep rate along with the hysteresis loop.

(iii) From the virgin curve (starting from $M=0, H=0$) of this loop, determine the field H_m corresponding to the minimum magnetization. The first field for full flux penetration, H^* , is approximately $1.5H_m$.

(iv) Estimate the field above which the ascending and the descending branches of the magnetization coincide. This field is the irreversibility field H_{irr} for the given sweep rate.

(v) Determine the irreversible component $M_{\text{irr}}(H)$ by averaging the ascending and the descending branches of the hysteresis loop.

(vi) Cool the sample again at zero field from above T_c to the temperature T . Apply a field H smaller than H_{irr} (to allow for magnetic relaxation) but large enough (at least $3H^*$, preferably much larger) to ensure full penetration and approximately linear flux profiles.

(vii) Decrease the field by a step ΔH of order $2H^*$ to ensure reversal of the flux profiles. (For measurements in a step-up field, this step is not needed.)

(viii) Measure the magnetization as a function of time. At this step $dM/d \ln t$ may be determined.

(ix) To determine the normalized relaxation rate $d \ln M/d \ln t$, normalize the data in step (viii) by M_{irr} .

(x) Go back to step (vii) if a series of data is desired at different field strengths.

Certainly, the above steps are not the only possible sequence. With insight to the problems mentioned above, one can construct equivalent routes with meaningful results.

B. Summary of experimental results

We turn now to a summary of experimental studies that, bearing in mind the above issues, appear in our judgment to offer reliable results. We start by noting that by far the majority of useful magnetic-relaxation studies have been performed in the context of a hysteresis-loop measurement, that is, starting with a sample cooled at zero field and isothermally varying the field to an appropriate level.

Another configuration of interest is that of field-cooled magnetization which, in early measurements, was thought to give equilibrium (Müller *et al.*, 1987). In fact, a small relaxation was detected in the field-cooled magnetization of YBCO crystals (Yeshurun and Malozemoff, 1988) as well as in other materials (Norling *et al.*, 1988). This had an important implication in comparing glass and flux-pinning theories: glassy theories implied equilibrium under field cooling, while a critical-state model for field cooling (Krusin-Elbaum *et al.*, 1990) predicted a small irreversible magnetization, with its consequent relaxation.

1. Temperature dependence

Perhaps the greatest number of studies focus on the temperature dependence of the relaxation rate. The relaxation, determined in a modest interval of time, is usually approximately logarithmic. Therefore, temperature dependencies are conveniently reported in terms of the relaxation rate $dM/d \ln t$, or in terms of the normalized relaxation rate $(1/M_{\text{irr}})dM/d \ln t \equiv d \ln M_{\text{irr}}/d \ln t$. In the latter case, the normalization is usually performed simply with the initial value of M_{irr} , as long as the relative change in M_{irr} over the time of the measurement is small. As discussed with respect to Eq. (2.10), there are theoretical reasons to focus on the normalized rate $d \ln M_{\text{irr}}/d \ln t$ but, as mentioned earlier, this involves corrections for reversible magnetization and substrate or sample-holder contributions. Thus determinations of the normalized rate are less reliable in general than those of the unnormalized data.

Figure 11 exhibits typical data on $dM/d \ln t$ for an YBCO crystal measured in a constant and fully penetrating magnetic field in the conventional isothermal hysteresis-loop procedure. All such experiments reveal a peak as a function of temperature, which tends to shift to lower temperature with increasing magnetic-field strength (see, for example, Földeàki *et al.*, 1989; Xu *et al.*, 1989, Xu *et al.*, 1991). This peak may, at first, look

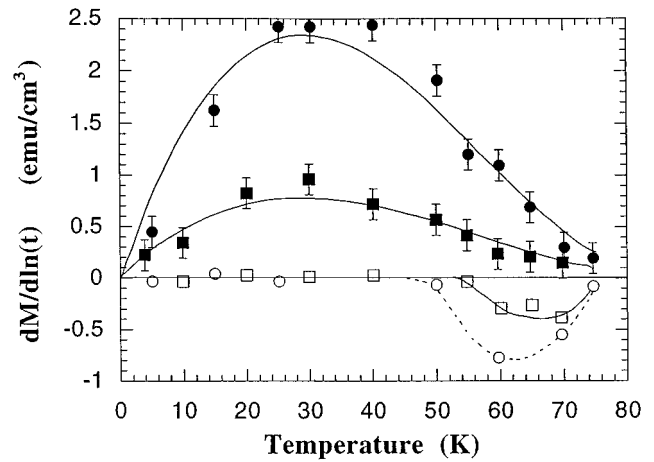


FIG. 11. Temperature dependence of the relaxation rate of field-cooled (open symbols) and zero-field-cooled (solid symbols) magnetization for fields parallel (circles) and perpendicular (squares) to the orthorhombic c axis of YBCO crystal. (After Yeshurun and Malozemoff, 1988.)

surprising because intuitively one would expect higher relaxation rates at higher temperatures. However, one should keep in mind that in the experiment, with a fixed time window, one probes different parts of the relaxation curve at different temperatures. At low temperatures the relaxation is expected to be slow and the relaxation measurements probe the initial process. On the other hand, at high temperatures the decay is fast and for the (fixed) experimental time window only the “tail” of the relaxation process is measured. This explains the apparent slow relaxation rate at high temperatures.

An example of the temperature dependence of the normalized relaxation rate $d \ln M_{\text{irr}}/d \ln t$ is given in Fig. 12, from the work of Civale *et al.* (1990), for an YBCO crystal in an applied field of 1 T parallel to the c axis. (Similar results were obtained by Campbell *et al.*, 1990)

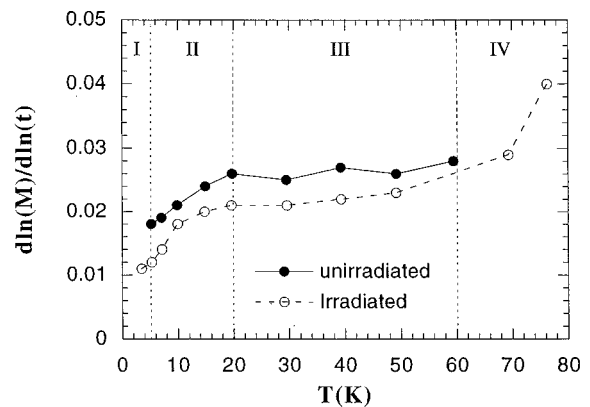


FIG. 12. Temperature dependence of the normalized relaxation rate for a YBCO crystal in an applied field of 1 T parallel to the c axis. The dashed line shows data obtained for the same crystal after irradiation with 3 MeV protons (flux of 10^{14} cm^{-2}). (After Civale *et al.*, 1990.)

and by Fruchter *et al.*, 1990.) The two curves in Fig. 12 represent the results before and after irradiation with 3 MeV protons to introduce additional pinning centers. Results such as these were a major surprise in the field. Although the linear temperature dependence predicted by the Anderson-Kim model [Eq. (2.7)] had never been adequately verified in low-temperature superconductors (Beasley *et al.*, 1969), the basic model had been assumed to be reasonable. According to this model, with increasing temperature, the barrier height could be expected to decrease, as quasiparticle excitations reduced the superconducting condensation energy, and was constrained to go to zero at T_c ; thus the normalized relaxation rate was expected to curve upwards. The results of Fig. 12 show the opposite trend, with the appearance of a plateau in the intermediate temperature range.

It is useful to divide the discussion of these data into four regions, labeled I to IV in the figure. At the lowest temperature (region I, below 2 K), the data appear to extrapolate to a finite intercept. This was also a surprise, since thermal activation was expected to freeze out completely at $T=0$. Finite relaxation at very low temperatures was first detected in a Chevrel-phase low-temperature superconductor (Mitin, 1987). More detailed studies have now been performed on high-temperature superconductors down to the millikelvin range, and the appearance of a low-temperature plateau has been confirmed in a variety of samples (Mota *et al.*, 1988, 1991; Hamzic *et al.*, 1990; Ivlev *et al.*, 1991), including high-quality YBCO (Fruchter, Malozemoff, *et al.*, 1991; Seidler *et al.*, 1993 and 1995) and BSCCO crystals (Fruchter, Hamzic, *et al.*, 1991; Aupke *et al.*, 1993), and a polycrystalline TBCCO (Tejada *et al.*, 1993). Typical data, from the work of Aupke *et al.*, are presented in Fig. 13. The magnetic relaxation at these ultralow temperatures has led to the suggestion of a novel mechanism of flux creep, namely quantum tunneling of vortices, facilitated by the small coherence length.

The findings of Seidler *et al.* show that quantum tunneling is far from a simple process. In their first paper (1993), they found that the tunneling in YBCO is dependent on the magnetic field, disappearing for fields above a certain critical value. Seidler *et al.* (1995) studied the temperature dependence of the tunneling for temperatures below 1 K, finding a noticeable linear temperature contribution to the action for vortex tunneling.

In the range from 4 K to about 20 K (region II), the data of Fig. 12 show a monotonic upward tendency, although the data are not detailed enough to identify a well-defined linear temperature dependence. This region has usually been interpreted, although rather lamely, in terms of Eq. (2.10) derived from the Anderson-Kim model. Barriers in the range 10–50 meV have been inferred.

The plateau in the intermediate temperature range of Fig. 12, from roughly 0.2 to 0.8 of T_c , is labeled region III. It is present in a number of studies on YBCO materials, although other data show only a broad peak. It is remarkable that for all of these YBCO materials, including flux-grown and melt-processed crystals and films, the

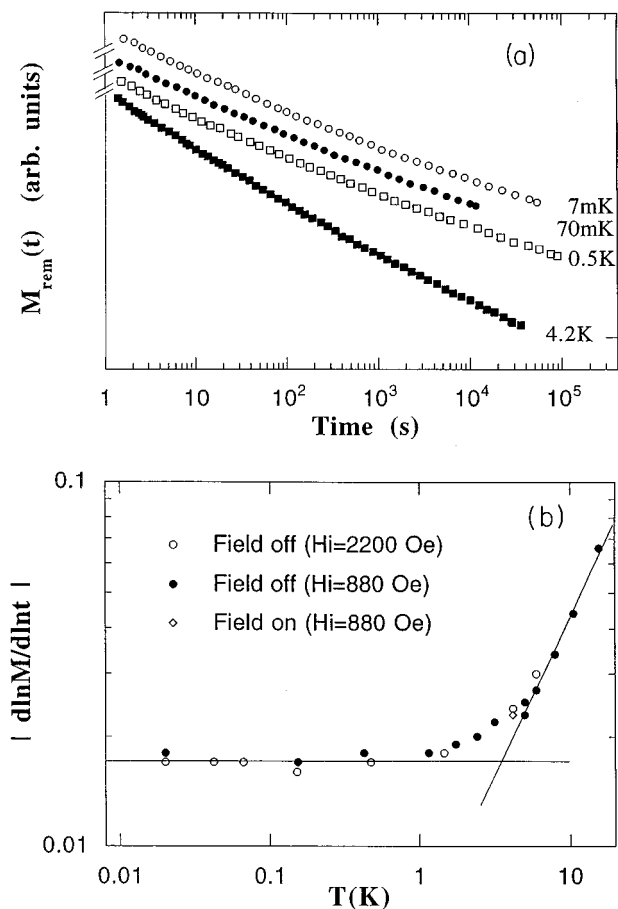


FIG. 13. Demonstration of flux creep at ultralow temperatures. (a) Time dependence of the remanent magnetization in a BSCCO crystal after cycling the crystal in a magnetic field of 880 Oe perpendicular to the ab plane. (b) Normalized decay rates of the remanent magnetization, as functions of temperature, after cycling the same crystal in magnetic fields of 880 and 2200 Oe. (After Aupke *et al.*, 1993.)

plateau or maximum value of $d \ln M_{\text{irr}} / d \ln t$ falls in the rather narrow range between 0.02 and 0.04 for fields of order 1 T applied parallel to the c axis. This has led to the suggestion of an approximate universality of the plateau values (Malozemoff and Fisher, 1990).

The plateau is baffling in the Anderson-Kim model since Eq. (2.10) implies that the barrier must increase linearly with temperature. Other models have been proposed (Griessen, 1990), involving a distribution of barrier heights, although the approximate universality in the plateau in such a variety of materials is hard to understand.

Another explanation for the plateau emerges from recent theories involving collective vortex creep (Feigel'man *et al.*, 1989, Feigel'man and Vinokur, 1990) or a vortex-glass model (Fisher, 1989), to be discussed in Sec. IV. We note that the plateau value is field dependent (Fruchter, Malozemoff, *et al.*, 1991). For example, in a YBCO crystal, the plateau value is shifted from 0.05 down to 0.015 with field (parallel to the c axis) increased from 0.2 T to 1.7 T. A possible explanation for the field dependence of the plateau value is based on the

collective-creep theory [see Sec. IV.C.1 and, in particular, Eq. (4.23)]. This theory predicts a larger creep parameter μ for collective creep than for single-vortex creep, in agreement with a smaller value of the plateau at high fields. The parameter μ is discussed more fully in Sec. IV.C.1.

Region IV concerns temperatures approaching T_c , where experimental limitations are most severe since M_{irr} becomes small. As a result, it is difficult to conclude whether near T_c the normalized rate increases, decreases, or maintains the same constant value. Perhaps the most accurate data to date are those of Konczykowski, Malozemoff, and Holtzberg (1991), who used a miniature Hall probe for relaxation measurements near T_c . They found in a YBCO crystal that the plateau is maintained up to a fraction of a degree from T_c . We note that region IV, approaching T_c , cannot be explained by the Anderson-Kim model simply because $kT \gg U$. In this regime, the phenomenon of magnetic relaxation is closely related to the dissipation found in transport measurements (Palstra *et al.*, 1988; Tinkham, 1988a; Iye *et al.*, 1990).

In contrast to YBCO, data on BSCCO and TBCCO crystals show no plateau in the intermediate temperature range. Instead, the normalized relaxation rate exhibits a clear maximum in this region. In fact, several groups (Nam, 1989; Zavaritsky and Zavaritsky, 1991; Zhukov *et al.*, 1992) reported, both for BSCCO and TBCCO, a double peak in the temperature dependence of the normalized relaxation rate, with an intervening dip around 30 K. This effect may be related to a transition between two different mechanisms of pinning, e.g., one involving bulk pinning and the other surface barriers.

2. Field dependence

Several groups dealt with the field dependence of the relaxation rate or the normalized relaxation rate (see, for example, Mota *et al.*, 1987; Yeshurun *et al.*, 1988b; Xu and Shi, 1990; Chaddah and Bhagwat, 1991; Moschalkov *et al.*, 1991; Krusin-Elbaum *et al.*, 1992; Yeshurun, Bontemps, *et al.*, 1994). We begin by describing the field dependence of the relaxation rate $dM/d \ln t$. In one of the first studies, Mota *et al.* (1987) reported for SrBaLaCuO ceramics a time-logarithmic decay of the isothermal remanent magnetization. For small fields H , applied and then removed to induce the remanent state, the rate was proportional to H^3 . While data on the remanent state are hard to interpret (see Sec. III.B.5), it can be presumed that the H^3 behavior stems from an only partially penetrated flux state. This behavior is expected to change at higher fields; well above the first field at which full flux penetration occurs, the remanent state should become field independent and the relaxation rate is expected to saturate.

Another explanation for the H^3 behavior stems from the analysis of Burlachkov *et al.* (1992) of the magnetization in the presence of surface barriers (see Sec. III.B.4.). The barriers hinder flux penetration every-

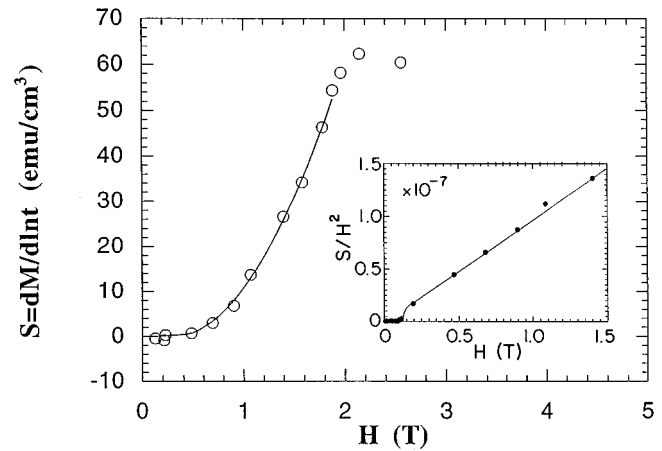


FIG. 14. Relaxation rate of the zero-field-cooled magnetization as a function of field for fields parallel to the orthorhombic c axis of a YBCO crystal at 6 K. The sharp onset of magnetic relaxation at $H = H_{c1}$ is demonstrated in the inset, which shows the relaxation rate, normalized to H^2 , as a function of field. The solid lines are fit to an extended Bean model. (After Yeshurun, Malozemoff, Wolfus, *et al.*, 1989.)

where except at a few point defects, where the barrier is suppressed. As a result, the flux profiles are different from the usual Bean profiles, and the leading term in the irreversible magnetization is proportional to H^3/J_c^2 instead of the usual H^2 behavior. Brandt and Indenbom (1993), in their study of the magnetization of a flat strip in a perpendicular magnetic field, explained the H^3 behavior without invoking surface barriers. They emphasize differences between the perpendicular geometry, for which the demagnetizing fields are important, and the longitudinal geometry, where demagnetizing fields are negligible. In particular, they show that the magnetization of a flat strip of width a in a partial critical state is $M = \pi a^2 H [1 - H^2/3J_c^2]$. In both explanations, in the presence of flux creep the critical current density J_c is replaced by the instantaneous current density J and the normalized creep rate is proportional to H^3 .

In the conventional isothermal hysteresis-loop procedure, where relaxation is measured in the presence of a field, a more complex behavior is observed. Typical data for a YBCO crystal are shown in Fig. 14. For fields above H_{c1} the relaxation rate follows a power-law increase in the low-field limit, slows down at higher fields, and peaks at a temperature-dependent field. The maximum, which is shifted to lower fields as the temperature increases, usually reflects full flux penetration. Similar behavior was reported by Moschalkov *et al.* (1991) for $\text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_x$ single crystals. This behavior can be interpreted (Yeshurun *et al.*, 1988b) on the basis of an extended Bean model (1964) that includes a field-dependent J_c , by introducing the Anderson-Kim logarithmic time dependence of the instantaneous current J , Eq. (2.9). It should be noted that a peak in the field-dependent relaxation rate had already been observed in conventional superconductors (Beasley *et al.*, 1969).

Magnetic relaxation data as a function of field are used for the determination of the field dependence of the activation energy. Typically, the activation energy decreases with increasing field. It is worthwhile noting that one may also utilize the onset of magnetic relaxation as a function of the field for the experimental determination of the lower critical field H_{c1} (Yeshurun *et al.*, 1988b). The logic behind this experiment is simply that at the Meissner state, for fields below H_{c1} , the system is reversible and no relaxation is expected—see the inset of Fig. 14. However, the presence of surface barriers can complicate this interpretation—see Sec. III.B.4.

The limited data on the behavior of the normalized relaxation rate, $d \ln M_{\text{irr}} / d \ln t$, as a function of field, show that it increases with field for weak fields. Several works report a crossover to a slower decay at intermediate fields (Krusin-Elbaum *et al.*, 1992; Yeshurun, Bontemps, *et al.*, 1994). This is interpreted in terms of a crossover from a single-vortex creep to collective creep. The theory of collective creep (Feigel'man *et al.*, 1989; Feigel'man *et al.*, 1991; Blatter *et al.*, 1994) will be discussed briefly in Sec. IV. We note that the crossover to a slower decay at intermediate fields is closely related to the anomalous second peak, or “fishtail,” observed in the hysteresis loops of YBCO, TBCCO, and BSCCO (Däumling *et al.*, 1990; see also Yeshurun, Bontemps, *et al.*, 1994 and references therein). This was suggested by Krusin-Elbaum *et al.* (1992) based on their observation of a strong correlation between the anomalous peak in the magnetization curves and the field dependence of the normalized relaxation rate in YBCO crystals at an intermediate temperature range. They proposed that the apparent anomalous peak is a result of slower relaxation in the field range of the peak. This interpretation requires further investigation in view of recent studies of YBCO crystals (see, for example, Yeshurun, Yacoby, *et al.*, 1994; Gordeev *et al.*, 1994; Küpfer *et al.*, 1994), which revealed no correlation between the peak and the relaxation rate at higher temperatures.

3. Deviations from time-logarithmic relaxation

The basic Anderson-Kim-model predicts that the time dependence of the magnetization $M(t)$ follows an approximately logarithmic law, i.e., $M(t) \propto \ln t$. Indeed, this prediction has been generally confirmed for conventional superconductors (Kim *et al.*, 1963a, 1963b; Beasley *et al.*, 1969). However, in the new superconductors, deviations from the logarithmic decay law were reported even in the initial stage of magnetic relaxation studies in these materials (Yeshurun and Malozemoff, 1988)—see Fig. 2. The description of these deviations is one of the main goals of the collective flux-pinning models to be described in Sec. IV. As emphasized in Sec. III.A, experimental conditions must be particularly well defined to make meaningful conclusions about such deviations, which are usually small. Also, it is essential to take data over a truly large time scale in order to be able to reliably test various models.

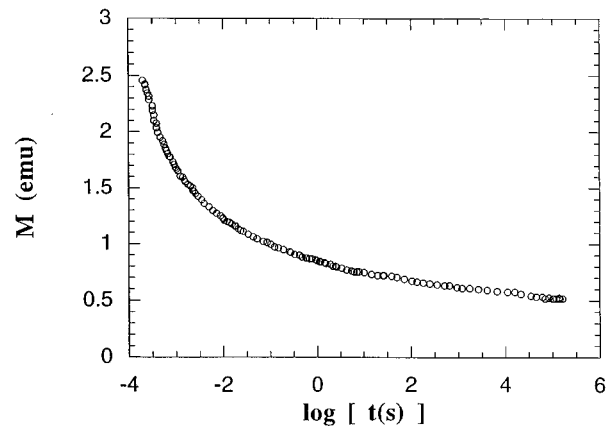


FIG. 15. The decay of the remanent magnetization over 9 orders of magnitude in time for melt-textured YBCO, demonstrating nonlogarithmic decay. (After Gao *et al.*, 1991.)

Gurevich and Küpfer (1993) pointed out that deviations from logarithmic time dependence of the magnetization are expected at the initial stage of the relaxation. This initial nonlogarithmic behavior is due to a transient redistribution of magnetic flux over the sample cross section. The duration of this stage is determined by the sample size, flux-creep rate and the rate of change of the magnetic field. To demonstrate this transient stage, they chose a specific set of magnetic-field changes in their study of grain-oriented YBCO with 5% of Ag. They prepared the system in the following way: at a constant temperature they applied a sufficiently high initial field, ranging from 2 T to 10 T in order to ensure complete flux penetration and a fully critical state in the sample. The field was then reduced with a constant sweep rate to the field at which the magnetic moment was measured as a function of time. Their measurements showed a plateau in the magnetization versus time in the transient regime during time intervals of 1–100 s; these time intervals are inversely proportional to the rate of change of the external field.

For experimental reasons, most time measurements started approximately 1–100 s after the change in field is completed and extended for 3 to 4 decades of measurement time. Several groups, however, explored short-time relaxation features using a variety of techniques, such as pulsed magnetometry (Gao *et al.*, 1991), fast field-switching time (Pozek *et al.*, 1991), and measurement of magnetic hysteresis loops at large sweep rates (Keller *et al.*, 1990; Jirsa *et al.*, 1993). In some cases, exponential time decays were observed, in agreement with the prediction of the thermally assisted flux-flow model (Kes *et al.*, 1989; van der Berg *et al.*, 1989). A particularly dramatic example of nonlogarithmic decays was given by Gao *et al.* (1991), who extended magnetic-relaxation measurements to times as short as 0.1 ms and so succeeded in measuring over nine orders of magnitude in time—see Fig. 15. The figure demonstrates a strong nonlogarithmic relaxation in the remanent mag-

netization of melt-textured YBCO, although the use of the remanent state complicates the interpretation (see Sec. III.A.5).

Systematic long-time relaxation measurements by Thompson, Sun, and Holtzberg (1991), on a proton-irradiated YBCO crystal, have also shown deviations from logarithmic behavior that fit the so-called collective-pinning interpolation formula, $\mathbf{J} \propto [1 + (\mu k \mathbf{T} / U_0) \ln(t/t_{\text{eff}})]^{-1/\mu}$ —see Eq. (4.21), in which μ is a parameter that characterizes the creep rate. A typical example of their data is shown in Fig. 2. The inset of Fig. 2 shows the μ values extracted from the fit of the data to the interpolation formula, for various temperatures. Within the vortex-glass framework (Fisher, 1989), there is no explanation for either the temperature dependence of μ or for values of μ larger than 1. However, as discussed in Sec. IV.C.1, the collective-creep theory provides for larger exponents. Also, the measured temperature dependence of μ agrees qualitatively with the predictions of this theory.

Fits to the interpolation formula have become a standard procedure in analyzing a nonlogarithmic dependence of magnetic relaxation. Sometimes, especially in the case of large μ values, it is difficult to distinguish between the interpolation formula and other models (such as a power law) by examining the time dependence of the magnetization alone. In these cases, the predictions of the interpolation formula were tested by plotting the time dependence of the normalized relaxation rate (Konczykowski *et al.*, 1993). A power law implies a constant S value, whereas the interpolation formula yields a linear increase of $1/S$ with $\ln t$ —see Eq. (4.22). It should be noted that the interpolation formula cannot always be applied, for example, in the presence of surface barriers—see Sec. III.B.4.

In some cases, where the interpolation formula failed, other formulas were used to fit the relaxation data. For example, Liu *et al.* (1992) observed a power-law time relaxation of the magnetization in a $\text{LuBa}_2\text{Cu}_3\text{O}_7$ single crystal. Similarly, Hergt *et al.* (1993) demonstrated that their data for a BSCCO crystal fit a power law in time. A reasonable fit to a power law has been reported for the temperature and field regimes of the anomalous second peak—the “fishtail” (Yeshurun, Bontemps, *et al.*, 1994). A transition from logarithmic flux creep to power-law behavior was reported by Xu *et al.* (1989). They related this crossover to the approach to an irreversibility temperature above which irreversible magnetization vanishes on practical experimental time scales. Safar *et al.* (1989) focused on the low-field behavior of the flux creep in BSCCO ceramics and demonstrated a crossover from a logarithmic dependence at low temperatures to a power-law dependence, at a temperature well below the irreversibility line. Xue *et al.* (1991) found that a different form of the interpolation formula, $M \sim \{\ln[(t+t_2)/t_1]\}^{-1/\mu}$ fits their pulsed data better than a power law (μ , t_1 , and t_2 were set as free fitting parameters). Brawner *et al.* (1993a) performed local measurements of the relaxation of the remanent magnetization in a single crystal of YBCO. They observed that at all

temperatures from 4.2 to 74 K, and independent of the location, the relaxation followed a power-law dependence $(1+t/\tau)^{-1/\sigma}$, as calculated by Vinokur *et al.* (1991) for a pinning barrier that depends logarithmically on J —see Sec. IV.

4. Magnetic relaxation over surface barriers

As has been shown by Bean and Livingston (1964), at the surface of type-II superconductors there appears a potential barrier to vortices entering or leaving the sample. The Bean-Livingston barrier arises from the competition between the repulsion of a vortex from the surface due to its interaction with the shielding currents and the attraction of the vortex to its “mirror image.” The Bean-Livingston barrier should be very pronounced in high-temperature superconductors, as has already been discussed (Kopylov *et al.*, 1990; McElfresh *et al.*, 1990; Konczykowski, Burlachkov, *et al.*, 1991; Burlachkov *et al.*, 1991 and 1994, Burlachkov, 1993a; Mints and Snapiro, 1993). Both surface barriers and bulk pinning determine the irreversible properties of the sample. Their relative contributions depend on the quality of the surface and vary with temperature and field.

Experimental evidence for surface barriers in high-temperature superconductors comes from several sources. The most typical “fingerprint” of surface barriers is found in hysteresis-loop measurements, where the descending branch of the loop hugs $M=0$ (Campbell and Evetts, 1972). Evidence is also found in the decreased width of the hysteresis loops after low-dose electron irradiation (Konczykowski, Burlachkov, *et al.*, 1991). Normally, irradiation is expected to increase pinning and irreversible magnetization by introducing defects. However, if surface barriers dominate the irreversibility, irradiation is expected to reduce the irreversibility by destroying surface perfection, on which surface barriers depend. As an immediate consequence of this effect, the apparent lowest field for flux penetration is reduced dramatically. Anomalous behavior in the temperature dependence of the lowest field for flux penetration, which is a direct consequence of retardation in flux entry due to surface barriers, has been observed at low temperatures (McElfresh *et al.*, 1990) and in the vicinity of T_c (Konczykowski, Burlachkov, *et al.*, 1991).

Of course surface barriers, as a source for irreversible magnetization, affect flux creep. Burlachkov (1993a, 1993b), using a model proposed by Clem (1974), predicted different relaxation rates for flux entry and exit. For flux exit the magnetization $M(t)$ was predicted to depend logarithmically on time, whereas for flux entry $M(\ln t)$ appeared to be a strongly nonlinear function with a downward curvature. Moreover, the initial relaxation rate $dM/d \ln t$ was predicted to be much larger for flux entry than for exit, in contrast to the case of conventional bulk creep. (Note that Beasley *et al.* (1969), also predicted such an asymmetry for bulk creep, but only in the case for which $\Delta B/B$ is large.) Asymmetry in the relaxation rate for flux entry and exit, in the pres-

ence of surface barriers, has been reported for BSCCO (Yeshurun, Bontemps, *et al.*, 1994).

As long as surface barriers dominate the irreversible magnetization, the magnetic relaxation associated with the irreversibility they create can be directly compared with the theoretical predictions. The challenge is to isolate this effect, which can drastically complicate the temperature dependence of the relaxation rate, and to avoid regions where this mechanism overlaps with a bulk mechanism. In cases where bulk pinning and surface pinning are of equal importance, there are reports of crossover phenomena. Chikumoto *et al.* (1992) found that flux creep in BSCCO crystals exhibited two different regimes as a function of time, as well as of temperature and magnetic field. For example, at a constant temperature and field $d \ln M/d \ln t$ exhibited a crossover to a slower decay; this crossover was shifted to shorter times at higher temperatures. They showed evidence that the long-time regime is dominated by surface barriers. Weir *et al.* (1991) reported the opposite behavior in YBCO, namely that surface barriers dominate the dynamics at short times. Burlachkov (1993a) pointed out that in the case of a competition between bulk and surface pinning, the initial relaxation is determined by the smaller pinning energy between U_{bulk} and U_{surface} . Thus, for example, if $U_{\text{bulk}} < U_{\text{surface}}$ then the bulk relaxation should be observed first. But as the persistent current associated with bulk pinning decreases, U_{bulk} increases and, after the crossover time, U_{surface} dominates the relaxation behavior.

We feel that the study of relaxation effects associated with surface barriers is in its infancy. In particular, systematic studies of relaxation rates as a function of field and temperature are still lacking. Also, the predicted asymmetry between flux entry and exit needs to be explored further. This will require overcoming the experimental difficulties posed by the initial almost-zero value of M in measurements of flux exit.

5. Memory effect

Rossel *et al.* (1989) and later Kunchur *et al.* (1990) reported a “memory effect” of magnetic relaxation in crystals of YBCO and BSCCO. In their experiments the sample was left in a quasiequilibrium state at a constant field H_1 and temperature T for a certain waiting time t_w . Then the applied magnetic field was increased to H_2 and the magnetic relaxation was measured. The memory effect was manifested by a “break,” or an inflection point, in the logarithmic time dependence of the magnetic relaxation at a time t_b after the change in the field. Rossel *et al.* reported that $t_b \cong t_w$. They related this observation to a similar memory effect in spin glasses and interpreted the results in terms of a superconducting-glass model (Morgenstern *et al.*, 1987). Kunchur *et al.* (1990) proposed a much simpler, and more natural explanation based on the flux-creep model. According to their model, the change in the relaxation rate was expected from the time evolution (due to relaxation in J) of the complex flux profile produced by the two consecu-

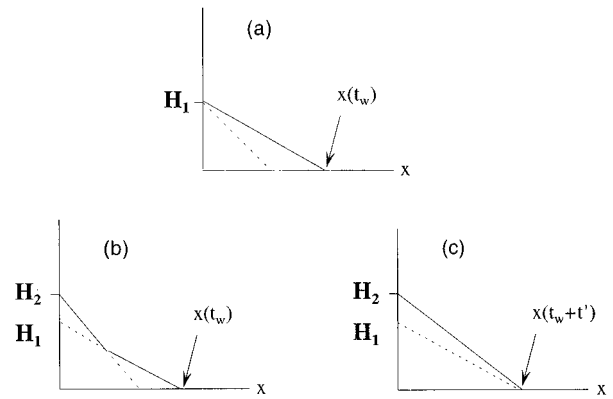


FIG. 16. Flux profile produced by two consecutive field steps. (a) Dotted line: the initial flux-density profile in the sample, for the field H_1 . Solid line: the profile after time t_w . (b) The field is increased to H_2 and the superposition of the initial profile of H_2 over the time-evolved profile of H_1 is sketched by the solid line. Note the “break” in this profile. (c) Thermally activated flux creep causes the slope to decrease with time, and after a time t' , the break disappears and the “memory” of the previous state is erased.

tive field steps, as illustrated in Fig. 16. The dotted line in Fig. 16(a) shows the initial flux-density profile in the sample, for the field H_1 . The solid line in this figure is the profile after time t_w . At this time, the field is increased to H_2 and the superposition of the initial profile of H_2 and the time-evolved profile of H_1 is sketched by the solid line in Fig. 16(b). Note the “break” in this profile. Thermally activated flux creep causes the slope to decrease with time, and after time t' the break disappears and the “memory” of the previous state is erased, as shown in Fig. 16(c). These changes in the profile affect the measured relaxation rate. Thus a crossover in the relaxation behavior is expected at time t' ; beyond this time, the relaxation should exhibit the usual logarithmic behavior for a direct field change from zero to H_2 . Kunchur *et al.* expressed t' in terms of H_1 , H_2 , T , and t_w and demonstrated that, in the general case, $t' \neq t_w$.

6. V - I curves

A relationship between flux motion and V - I , or E - J , curves (Griessen, 1991a; van der Beek, Nieuwenhuys, *et al.*, 1992; Ries *et al.*, 1992; Sandvold and Rossel, 1992; Konczykowski *et al.*, 1993; Caplin *et al.*, 1994) is based on the fact that this motion creates an electric field $\mathbf{E} = (1/c)\mathbf{B} \times \mathbf{v}$, where \mathbf{v} is the average velocity of the flux lines in the direction of the Lorenz force (see Sec. II.C). For an unpinned but damped flux-line lattice, \mathbf{v} is proportional to the current density \mathbf{J} , and one obtains a linear relationship between V and I (Bardeen and Stephen, 1965). When pinning is important, the average velocity associated with the thermally activated jumps of flux lines is $v = v_0 \exp(-U(J)/kT)$, where the prefactor v_0 may also be a function of J (see Sec. IV.A); so E is also exponentially dependent on $U(J)$. Thus flux creep is generally associated with a highly nonlinear V - I rela-

tionship, dictated by the specific dependence of U on J and by the exponential dependence of E on U .

For example, the standard flux-creep model of Anderson and Kim assumes a linear dependence of the effective energy barrier on the current density, $U = U_0(1 - J/J_c)$, giving rise to an exponential V - I curve. The logarithmic model of Zeldov *et al.* (1989, 1990) assumes $U = U_0 \ln(Jc_0/J)$, resulting in a power-law V - I dependence. The vortex-glass and the collective-creep models (see Sec. IV.C.1) result in nonexponential E - J characteristics of the type $E \propto \exp\{-\text{const}/J^\mu\}$, where μ depends on magnetic field, temperature and current (see Blatter *et al.*, 1994).

V - I (or E - J) curves can be generated from magnetization measurements (Ries *et al.*, 1992; Sandvold and Rossel, 1992; Konczykowski *et al.*, 1993; Küpfer *et al.*, 1994). Because dM/dt is proportional to the electric field at the sample surface and M is proportional to the current, plots of dM/dt versus M represent the E - J dependence. This contactless technique allows J - E measurements in ranges inaccessible in conventional transport methods (Konczykowski *et al.*, 1993; Küpfer *et al.*, 1993).

A power-law V - I relationship observed in transport measurements of superconductors has been classically attributed to nonuniformity in J_c , and models have been derived (Warnes and Larbalestier, 1986; Plummer and Evetts, 1987) on this basis to explain the power law in low-temperature superconductors. In high-temperature superconductors, however, the power-law V - I relation can come from flux creep (Griessen, 1991a; Sun *et al.*, 1991), as explained above. Sometimes, both effects can be observed, for example in recent studies on strained and unstrained wires (Adamopoulos and Evetts, 1993; Malozemoff *et al.*, 1992).

7. Experimental determination of $U(J)$

Beasley *et al.* (1969) were the first to point out that the linear dependence of the activation energy U on the persistent current J , Eq. (2.5), is only a first approximation. In general, $U(J)$ is a nonlinear function that may diverge at $J=0$, as discussed in detail in Sec. IV.C. Maley *et al.* (1990) proposed a technique, based on the analysis of flux-creep measurements, for an experimental determination of $U(J)$. From the rate equation for thermally activated motion of flux (Beasley *et al.*, 1969) they showed that $U = AT - kT \ln|dM/dt|$, where A is a time-independent constant. Both $kT \ln|dM/dt|$ and $M_{\text{irr}} \propto J$ can be experimentally determined. Thus one may generate plots of $kT \ln|dM/dt|$ vs M_{irr} at different temperatures which, up to an additive constant AT , present U vs J . Typical results are presented in Fig. 17(a) for polycrystalline YBCO. The principal effect of increasing temperature is to produce monotonically decreasing values of M_{irr} . As M_{irr} decreases, the isotherms in the figure become progressively steeper. Indeed, selection of a single constant A multiplied by T for each temperature causes all of the data to fall on the same smooth $U(J)$ curve, as demonstrated in Fig. 17(b). From these data

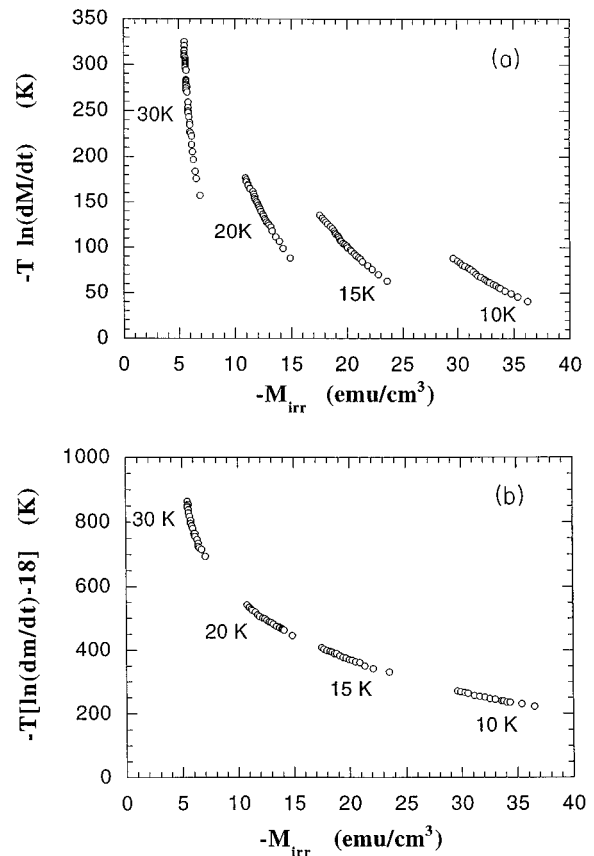


FIG. 17. Demonstration of the Maley technique for determining the dependence of the flux-creep activation energy upon current density. (a) Plots of the $T \ln|dM/dt|$ vs the irreversible magnetization for a polycrystalline YBCO at the indicated temperatures. A magnetic field of 10 kOe parallel to the c axis was used. (b) The same data replotted after adding $18T$ to each data set (T is temperature). This procedure causes all the data to fall on the same smooth curve, which represents $U(J)$. (After Maley *et al.*, 1990.)

Maley *et al.* concluded that U depends logarithmically on J , in agreement with the logarithmic dependence obtained by Zeldov *et al.* (1989, 1990) from transport measurements. Other approaches in reconstructing the dependence of U on J from magnetic-relaxation data have been discussed by Sengupta *et al.* (1993a, 1993b) and by Hu (1993).

A particularly interesting method for determining $U(J)$ was recently developed by Abulafia *et al.* (1995). This method utilizes a novel Hall-probe array technique (Zeldov, Majer, *et al.*, 1994) to measure the local induction B simultaneously at different locations as a function of time, thus enabling direct analysis of the flux creep on the basis of the flux diffusion equation [Eq. (4.2)]. Figure 18 displays typical data for the magnetic induction as a function of time at different locations in a YBCO crystal measured at 50 K in the remanent state. The inset describes the position of the various Hall probes relative to the sample. Evidently, the relaxation rate $dB/d \ln t$ is maximum near the center and decreases toward the edge. Probe 3, located near the point where $B \approx 0$ shows

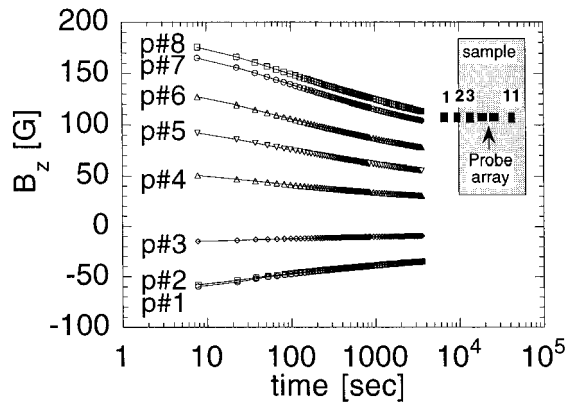


FIG. 18. Local induction vs time measured in and out of a YBCO crystal at $T=50$ K by an array of miniature Hall probes, numbered 1 through 8. Inset: Configuration of the Hall-probe array relative to the crystal. (After Abulafia *et al.*, 1995.)

approximately zero relaxation rate. This point corresponds to a contour on which vortices of different sign annihilate.¹ Due to the demagnetization effect, the induction at the sample edge increases with time, exhibiting a positive relaxation rate. These results emphasize the nonuniformity of the local relaxation (Gurevich and Brandt, 1994), thus questioning the meaning of global relaxation measurements in which negative and positive contributions are combined. In addition, it is important to realize that in global measurements one actually measures the activation energy at the surface of the sample, while the current J is averaged over the sample volume (van der Beek, Nieuwenhuys, *et al.*, 1992).

Analyzing the local data on the basis of the flux diffusion equation yields the local activation energy U as a function of time, as shown in Fig. 19. As expected, U increases with time, reflecting the decrease in J . The linear dependence of U on the logarithm of time, in the long-time limit, confirms the theoretical prediction given in Eq. (4.7). By extrapolating the linear portion backward to $U=0$, the logarithmic time scale t_0 was found to be in the range 1–10 microseconds. From the measured spatial distribution of B , Abulafia *et al.* determined the local J and constructed isotherms for the $U(J)$ dependence. Their results agree with the logarithmic model, Eq. (4.20).

Further discussion of different $U(J)$ dependences is given in Sec. IV.C.

C. Suppression of magnetic relaxation

Controlling and reducing magnetic relaxation is the goal of numerous efforts motivated by the need to understand the basic physics behind the “giant” relaxation, as well as by the need to suppress the creep in order to

¹This “neutral line” where $\partial B_z/\partial t$ changes sign was also observed in magneto-optic measurements (Forkl *et al.*, 1993; Schuster, Kuhn, and Brandt, 1995; Schuster *et al.*, 1995).

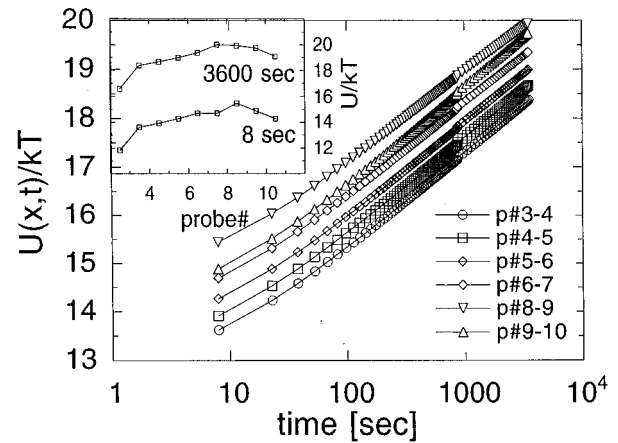


FIG. 19. Local activation energy U/kT vs time measured in a YBCO crystal at 40 K by an array of 11 Hall probes. Equation (4.7) holds for all the probes. U is almost uniform throughout the sample, as shown in the inset. (After Abulafia *et al.*, 1995.)

avoid its destructive effect on potential device applications. In the following we summarize the experimental approaches to achieving this goal.

1. Modulation of flux profiles and flux annealing

One of the experimental techniques to reduce the creep rate takes advantage of the fact that the activation energy U increases as the current density J decreases (Beasley *et al.*, 1969; Feigel'man *et al.*, 1989; Maley *et al.*, 1990; Zeldov *et al.*, 1990; Malozemoff, 1991). Usually, immediately after a rapid field change, the gradient dB/dx is proportional to $J \cong J_c$. It is possible, however, to “freeze” a different flux configuration in such a way that $J \ll J_c$. This can be done, for example, by changing the field at temperature T_A and then cooling the sample to the measuring temperature $T_m < T_A$. At T_A the gradient of the flux profile is proportional to $J_c(T_A)$. This flux profile is frozen if the sample is cooled down sufficiently rapidly to T_m . Of course, $J_c(T_A) < J_c(T_m)$ and thus the slope of the flux profile obtained now at T_m is smaller than that obtained by changing the field at T_m . Using this procedure which is referred to as “flux annealing,” a dramatic decrease in relaxation rates has been achieved (Maley *et al.*, 1990; Sun *et al.*, 1990, Thompson, Sun, Malozemoff, *et al.*, 1991). The extremely slow relaxation rates found in field-cooled samples (see Fig. 11) reflect the same effect.

The flux annealing phenomenon suggests that the time evolution of J depends only on its instantaneous state and not on the path used to reach it. Thus one might expect the annealed relaxation to coincide with that observed after a long time when J decays to the same level. Indeed, this was verified by Thompson, Sun, Malozemoff, *et al.* (1991) in their study of flux-creep annealing in YBCO crystals. An example of their results is shown in Fig. 20. The solid line represents long-time relaxation when the field was applied and the temperature was maintained at $T_m=30$ K. The relaxations measured

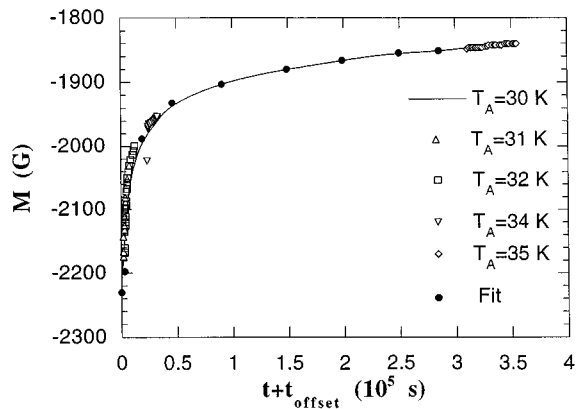


FIG. 20. “Flux-creep annealing” with subcritical J at $T_m=30$ K. The open points show relaxation measured after returning to $T_m=30$ K from $T_A=31, 32, 34,$ and 35 K. The solid line shows experimental relaxation data when a field of 1 T was applied parallel to the c axis and the temperature was maintained at $T_m=T_A=30$ K. The positions of the data points on the time axis are shifted by an empirically determined t_{offset} to match the starting magnetization, so as to align similar values of M . The solid points are fits to the interpolation formula, Eq. (4.27), with $\mu=1.9$. (After Thompson, Sun, Malozemoff, *et al.*, 1991.)

after returning T_m from $T_A=31, 32, 34,$ and 35 K, are shown by the open symbols. Their positions on the time axis are shifted by an empirically determined t_{offset} to match the starting magnetization so as to align similar values of M . The figure shows that subsequent time evolutions of M are identical, thus confirming the simple hypothesis introduced above. The time-dependent magnetization presented in this way is well described by the collective creep formula, Eq. (4.21), shown by the solid circles in the figure.

2. Effect of introducing defects

The flux-creep rate is expected to decrease when new—or denser—pinning centers are introduced. Experimentally, this is accomplished by introducing defects either chemically (doping, changing stoichiometry, etc.), mechanically, or by irradiation. The effects of these three approaches are summarized below. In general, all these techniques are capable of producing defects of widely varying sizes.

Various groups compared magnetic relaxation rates and calculated pinning potentials for materials prepared by different techniques. These experiments showed that the method of preparation did indeed have some effect on flux pinning, presumably due to different, though uncontrolled, kinds of defects. Thus, for example, plasma arc-melting and rapid-quenching methods procedures (Kamino *et al.*, 1990) are claimed to produce samples with relatively large activation energy and reduced creep rates. Similarly, a quench-and-melt growth technique produced polycrystalline YBCO without weak links (Morita *et al.*, 1990), and with reduced relaxation.

Shi *et al.* (1989), Shi, Xu, Fang, *et al.* (1990), and Murakami *et al.* (1991), among others, demonstrated that pinning can be improved by introducing nonsuperconducting precipitates into high-temperature-superconductor samples. For relatively large precipitates, pinning may be provided by either the nonsuperconducting volume or by the interface between the superconducting and nonsuperconducting materials. Other groups concentrated on relaxation rates as a function of stoichiometric changes in virtually all atoms of the unit cell. For example, studies of the relaxation in oxygen-deficient, magnetically aligned, polycrystalline $\text{YBa}_2\text{Cu}_3\text{O}_{7-\delta}$ samples (Ossandon *et al.*, 1992), showed that the normalized flux-creep rate does not change significantly with oxygen content in the range $0 \leq \delta \leq 0.2$. Its value centers about 0.02 and its temperature dependence shows the well-known “plateau” pattern (see Sec. III.B.1).

Chemical doping and substitution of the unit cell elements also affect creep rate—see for example, Paulius *et al.* (1993), Paulose *et al.* (1993), Qin *et al.* (1993), and Balanda *et al.* (1993). Chemical doping may increase the intragranular critical current density through the formation of nonsuperconducting phases that provide regions where the superconducting order parameter is suppressed and flux lines are pinned. Chemical substitutions may have a similar effect when the chemical impurity interacts locally with the superconducting holes.

A dramatic increase in flux-pinning energy after high-pressure shock-wave treatment was reported for powder compacts of high-temperature superconductors (Seaman *et al.*, 1989; Venturini *et al.* 1989; Weir *et al.*, 1990). The increased pinning energy is apparently caused by shock-induced defects.

Irradiation is a conventional technique for producing defects in a controlled way. Numerous irradiation experiments have been reported for high-temperature superconductors, using a wide range of photons (x rays, gamma rays) and particles (electron, neutrons, protons, and heavy ions) with energies ranging from eV to GeV. Effects of “internal” irradiation with fission fragments from ^{235}U were also investigated by doping YBCO with uranium and exposing the material to thermal neutrons (Fleischer *et al.*, 1989). While most experiments dealt with the effects that irradiation has on T_c and J_c , only a few focused on magnetic relaxation. In some polycrystalline materials the effect of irradiation is almost unnoticeable. For crystals, Civale *et al.* (1990) reported virtually no change in the normalized relaxation rate for YBCO after proton irradiation (see Fig. 12) but Venturini *et al.* (1990), Konczykowski, Rullier-Albenque *et al.* (1991) and Gerhauser *et al.* (1992), reported a reduction in the relaxation rate for high-temperature-superconductor crystals irradiated with heavy ions.

Heavy-ion irradiation (for example by 580 MeV Sn ions, 1–2 GeV Xe ions, 5–6 GeV Pb ions) creates special defects in the form of cylindrical amorphous tracks (“columnar defects”). These defects provide maximum possible pinning of flux lines parallel to the tracks, because they match the linear topology of vortices, and their

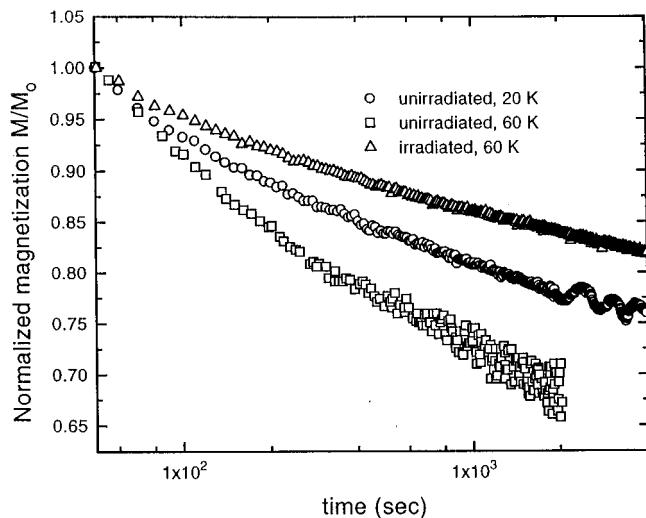


FIG. 21. Effect of heavy-ion irradiation on flux creep in YBCO crystals at 60 K. The relaxation rate, normalized to its initial value, is reduced significantly after irradiation. Also shown are relaxation data at 20 K before irradiation. In the present experiment, J at 60 K after irradiation corresponds to J at 20 K before irradiation. (After Prozorov *et al.*, 1995b.)

diameter—of order 10 nm—matches the vortex-core dimension quite well. Indeed, they induce a significant enhancement of magnetic irreversibility (Bourgault *et al.*, 1989; Hardy *et al.*, 1991; Civale *et al.*, 1992) and suppression of magnetic relaxation (Konczykowski, Rullier-Albenque, *et al.* 1991; Konczykowski *et al.*, 1993; Gerhauser *et al.*, 1992; Konczykowski, 1993).

Figure 21, from the work of Prozorov *et al.* (1995b), demonstrates the effect of heavy-ion irradiation on magnetic relaxation in YBCO crystals. To ensure the conditions of the critical state and full flux penetration, the full hysteresis loop was measured at several temperatures for both the unirradiated and the Pb-irradiated (10^{11} ions/cm²) crystals. The width ΔM of the hysteresis loop increased after irradiation, for example, by a factor of 5 at 60 K and 1 T. The figure shows that the relaxation rate, normalized to its initial value, was reduced significantly after irradiation. We note that one should be cautious in comparing these data because the relaxation process is determined by U , which is a function of J . Thus an increase in J caused by the irradiation automatically influences U and the relaxation rate. A more meaningful comparison should be made for isotherms representing similar J values. In this experiment, J at 60 K after irradiation corresponded to J at 20 K before irradiation. The relaxation data for the latter are also shown in Fig. 21 and apparently indicate that the irradiation affects U directly and not only through J .

In their columnar-defect experiments on BSCCO crystals and wires, Gerhauser *et al.* (1992) found pinning energies from relaxation that were comparable to the pinning energy calculated for a defect in a single CuO₂ double layer. This result can be explained from the suggestion of Clem (1991) that in the family of BSCCO materials, the unusually weak coupling between super-

conducting blocks across the Bi-O double layers caused the vortices to break up into stacks of “pancakes,” with individual pancakes relatively free to creep independently of the others. Another nucleation-creep model (Nelson and Vinokur, 1992) with an effective energy barrier increasing with time, i.e., with the decrease of the persisting currents, provides a more detailed description of the observed magnetic decays (Konczykowski *et al.*, 1995).

A question of interest is the comparison of the creep rate above and below the “matching field,” namely, the field at which the number of columnar defects and the number of vortices are nominally equal. Khalfin and Shapiro (1993) showed that, above the matching field, interdefect regions can also play the role of pinning centers; the potential barriers in these regions are induced by the vortices pinned in the columnar defects.

IV. THEORETICAL APPROACHES TO MAGNETIC RELAXATION IN HIGH-TEMPERATURE SUPERCONDUCTORS

Historically, the first samples of high-temperature superconductors were sintered and viewed as clusters of weakly coupled superconducting grains, forming a “glassy superconductor” as was described, for example, by Ebner and Stroud (1985) and by Morgenstern *et al.* (1987). The relaxation effects in these new superconductors were originally attributed to their glassy state, and were predicted to resemble similar phenomena observed in frustrated systems such as spin glasses (for a recent review on spin glasses, see Fischer and Hertz, 1991; for the superconducting-glass model see also Rae, 1991, and Mee *et al.*, 1991). Some authors also applied the glassy model to crystals (Deutscher and Müller, 1987; Tinkham, 1988a), though the origin of granularity in crystals is still an open question.

Yeshurun and Malozemoff (1988) and Dew-Hughes (1988) pointed out that the giant relaxation effect observed in high-temperature superconductors may arise from thermally activated flux creep, as was first described by Anderson and Kim (1964) and later by Beasley *et al.* (1969). The basic Anderson-Kim model (see Sec II.E) was successfully applied by numerous authors to explain many of the relaxation data in high-temperature superconductors, especially in the low-temperature, low-field regimes. However, at relatively high temperatures and fields, noticeable disagreements with the basic flux-creep model have been observed. In addition, creep has been observed at ultralow temperatures where thermal activation is expected to freeze out. Various models were developed to explain the relaxation data in these problematic regimes. For example, the thermally activated flux flow model was proposed to explain the data at elevated temperatures (Kes *et al.*, 1989), and quantum creep was invoked to interpret the low-temperature data (Simanek, 1989; Blatter *et al.*, 1991; Fruchter, Hamzic, *et al.*, 1991; Fruchter, Malozemoff *et al.*, 1991; Ivlev *et al.*, 1991). A model for calculating the crossover temperature between thermally ac-

tivated creep and quantum tunneling was derived by Ma *et al.* (1993). It may be interesting to note that, in some recent works, attempts were made to describe the relaxation at ultralow temperatures as a thermally activated process—see, for example, Matsushita (1993) and Gerber and Franse (1993, 1994); see also comments on the latter by Griessen, Hoekstra, *et al.* (1994) and by Fruchter *et al.* (1994).

Interpretation of experimental data within the range in which the basic Anderson-Kim model may be applicable has led to some peculiar results. For example, a straightforward application of the Anderson-Kim model yields a barrier that increases linearly with temperature (Xu *et al.*, 1989). This has led to extensions of the basic model by introducing a distribution of activation energies (Malozemoff, Worthington, Yandroski, *et al.*, 1988; Hagen and Griessen, 1989) or a nonlinear dependence of U on J (Beasley *et al.*, 1969; Zeldov *et al.*, 1989; Maley *et al.*, 1990). A novel approach for thermally activated flux creep has been developed by Feigel'man *et al.* (1989, 1991)—see Sec. IV.C.1 below. Their “collective-creep” model assumes weak random pinning, and vortex creep is viewed as the motion of an elastic object through a random potential. Within this model, a nonlinear $U(J)$ and a barrier distribution are included.

The Anderson-Kim model and its extensions predict that the mobility of vortices freezes gradually as the temperature is lowered. A completely different approach has been proposed by Fisher (1989), who predicted a thermodynamic phase transition of the vortex system from a “vortex-liquid” state where the vortices are highly mobile into a “vortex-glass” phase where the vortices are immobile. The term “vortex-glass” relates to a situation in which the long-range order of the Abrikosov lattice is destroyed due to sample imperfections. It should be noted that total freezing of flux mobility is predicted only in the limit of infinitely small current. In this limit U is predicted to diverge, and the flux array develops a long-range glassy order. However, for finite J , creep is still expected. The static and dynamic properties of flux lines in the vortex-glass regime were extensively studied by Feigel'man *et al.* (1989), Natterman (1990), and Blatter *et al.* (1994) within a “collective-pinning” model.

Other models not included in the above categories have also been proposed. For example, the recent concept of self-organized criticality (Bak *et al.*, 1987) was applied to flux creep in high-temperature superconductors (Koziol *et al.*, 1993; Tang, 1993; see also Richardson *et al.*, 1994; Field *et al.*, 1995; Reichhardt *et al.*, 1995). These approaches are still too premature to be reviewed here. In the following we concentrate on the Anderson-Kim model and its extensions, and on the collective-creep model. Our goal here is to provide a general overview of the theory for the experimentalist. For a more detailed review, the reader is referred to Blatter *et al.* (1994).

A. The electrodynamic equation of flux creep

In this section we review the classical theory of thermally activated flux creep, following the treatment of

Blatter *et al.* (1994). We consider a simple slab geometry with flux lines or flux-line bundles aligned along the z axis and moving along the x axis (see Fig. 22). For the more complex geometry of a flat strip in a perpendicular field the reader is referred to Gurevich and Brandt (1994) and Zeldov, Clem *et al.* (1994). In the following discussion we assume that B is large compared to the magnetization, i.e., $B \approx H$. The macroscopically averaged current density J is in the y direction and is related to the vortex-density gradient, established by pinning, through the quasistatic Maxwell's equation:

$$\frac{\partial B_z}{\partial x} = -\frac{4\pi}{c} J_y, \quad (4.1)$$

in which the displacement-current term is omitted. Henceforth we drop the subscript y and z to simplify notation, but we always consider only one axis for the average flux direction. This field gradient decays with time as a result of the thermally activated motion of vortices, thus causing relaxation of the persistent current J and of the magnetization M (Anderson, 1962; Anderson and Kim, 1964). The basic equation governing the decay of the current J can be derived from the Maxwell's equation $\partial E/\partial x = -(1/c)\partial B/\partial t$ and the equation relating the electric field to the flux motion $E = (1/c)Bv$, where v is the velocity of the vortices in a direction parallel to the Lorentz force, i.e., the x direction [see Eq. (2.4) and Fig. 22]. These considerations lead to the equation of continuity for the flux-line density (Beasley *et al.*, 1969):

$$\frac{\partial B}{\partial t} = -\frac{\partial}{\partial x}(vB). \quad (4.2)$$

Using Eq. (4.1), one obtains a corresponding dynamic equation for the current density J :

$$\frac{\partial J}{\partial t} = \frac{c}{4\pi} \frac{\partial^2}{\partial x^2}(vB). \quad (4.3)$$

Assuming thermal activation over the pinning barrier $U(J)$, the velocity v in Eqs. (4.2) and (4.3) is given by

$$v = v_0 e^{-U(J)/kT}, \quad (4.4)$$

where the preexponential factor is given by $v_0 = x_0 \omega_m J/J_c$, x_0 is the hopping distance, ω_m is the microscopic attempt frequency, and the factor J/J_c is introduced to provide a gradual crossover to the viscous-flow regime in which $v \propto J$ at $kT \gg U$ (Feigel'man *et al.*, 1991; Vinokur *et al.*, 1991).

Blatter *et al.* (1994) solved Eq. (4.3) assuming complete field penetration and neglecting field dependence of the barrier $U(J)$. Their result indicates that space variation of J may be neglected throughout most of the sample region, except for a narrow region near the center of the sample where J changes sign. For a constant J , Eq. (4.1) implies that B varies linearly from the sample surface ($x=d$) to its center ($x=0$), i.e., $B = (4\pi/c)(d-x)J + H$, where H is the applied field. The geometry and dependence of various quantities on x is shown in Fig. 22.

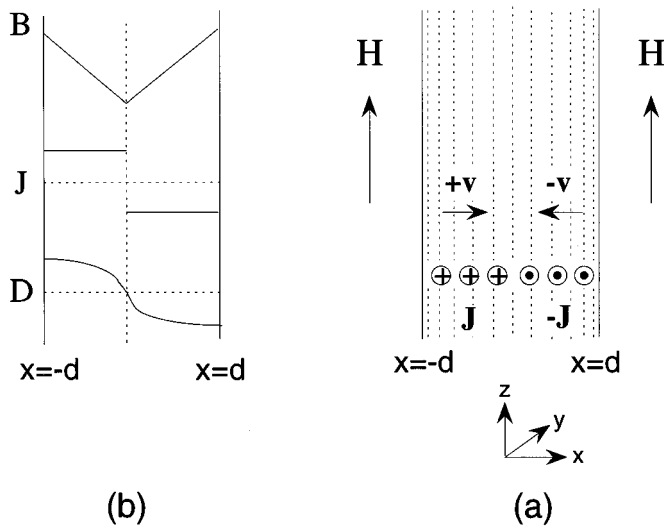


FIG. 22. Flux creep in a slab: (a) A schematic diagram showing the current density \mathbf{J} and the flux-line velocity \mathbf{v} for an infinite superconducting slab immersed in a parallel field \mathbf{H} . The induction \mathbf{B} , the current \mathbf{J} , and the velocity \mathbf{v} are parallel to the z , y , and x axes, respectively. Dotted lines represent flux lines. (b) A schematic plot of \mathbf{B} , \mathbf{J} , and the flux-line current density $D=Bv$ as functions of the distance x from the center of the slab.

Integration of Eq. (4.2) between the center and the edge of the slab yields

$$\frac{\partial J}{\partial t} \approx -\frac{cv_0 H}{2\pi d^2} e^{-U(J)/kT}, \quad (4.5)$$

where U is the activation energy at the surface of the slab. In deducing this equation we utilized two boundary conditions: (a) at the center of the slab ($x=0$) the same amounts of flux move in opposite directions, thus the net flux-current density $vB=0$; (b) at the surface of the slab, ($x=d$), B is equal to the applied field H , which is independent of time. A similar equation for the activation energy U can be obtained by substituting $dJ/dt=(dU/dt)(dU/dJ)^{-1}$ in Eq. (4.5):

$$\frac{dU}{dt} = -\frac{cv_0 H}{2\pi d^2} \frac{dU}{dJ} e^{-U(J)/kT}. \quad (4.6)$$

This equation can be solved with logarithmic accuracy (Geshkenbein and Larkin, 1989), yielding

$$U(J) = kT \ln(t/t_0), \quad (4.7)$$

where² $t_0 = 2\pi kT d^2 / (cv_0 H |dU/dJ|)$. As noted by Feigel'man *et al.* (1989), t_0 is a macroscopic quantity depending on the sample size d , and it should not be confused with the actual microscopic attempt time. Equa-

²In some publications (e.g., Schnack *et al.*, 1992; van der Beek, Kes, *et al.*, 1992; van der Beek, Nieuwenhays, *et al.*, 1992; Gurevich and Kupfer, 1993) a solution of the form $U(J) = kT \ln(t + t_{\text{initial}}/t_0)$ is adopted with $t_{\text{initial}} = t_0 \exp[U(j_{\text{initial}}/kT)]$ and j_{initial} is the current density at the beginning of the relaxation.

tion (4.7) is general and is independent of the specific functional form of $U(J)$. The time evolution of the screening current density J can be determined directly from Eq. (4.7), provided that the functional dependence of U on J is known. In the following we discuss linear and nonlinear dependence of U on J . The nonlinear behavior of $U(J)$ is of little interest for low-temperature superconductors in which the persistent current J is in the close vicinity of the critical current on any reasonable time scale. However, for high-temperature superconductors nonlinearities become important because of the giant flux creep.

B. Linear $U(J)$ —the Anderson and Kim model

Anderson (1962) and Kim and Anderson (1964) considered the simplest possible model of pinning with a linear J -dependence of the barrier energy:

$$U = U_0 - \frac{1}{c} JBx_0V, \quad (4.8)$$

where V is the volume of the “jumping” flux bundle and x_0 is the hopping distance. [We consider here only hopping in the forward direction, i.e., in the direction of the Lorentz force. The contribution of backward hopping is important for $U \leq kT$ as discussed by several authors—see, for example, Dew-Hughes (1988) and Hagen and Griessen (1990)]. The second term in the right-hand side of Eq. (4.8) is the effective reduction in the barrier due to the work done by the Lorentz force in moving the flux bundle over the distance x_0 . In this simple model V and x_0 are constants, independent of J and B , thus U depends linearly on J :

$$U = U_0 \left(1 - \frac{J}{J_{c0}} \right), \quad (4.9)$$

where $J_{c0} = cU_0/Bx_0V$ is the critical current density at which the barrier vanishes. We recall that in Sec. II.C we defined the critical current J_c as the current at which the Lorentz force is exactly balanced by the pinning force. In fact, in this critical state the barrier vanishes and thus $J_c = J_{c0}$.

As mentioned above, the linear approximation, Eq. (4.9), is a reasonable approximation near J_{c0} and is a fair description for conventional superconductors for which $U_0 \gg kT$ and therefore the persistent current is always close to J_{c0} . However, for high-temperature superconductors, because of the giant flux creep, J may be much smaller than J_{c0} and hence the nonlinear dependence of U on J may become extremely important, as will be discussed in the next section.

From Eqs. (4.7) and (4.9) one obtains the famous logarithmic time dependence of the current density:

$$J = J_{c0} \left[1 - \frac{kT}{U_0} \ln \left(\frac{t}{t_0} \right) \right]. \quad (4.10)$$

To obtain an expression that is applicable from $t=0$, it is usually accepted that one can rewrite Eq. (4.10) as

$$J = J_{c0} \left[1 - \frac{kT}{U_0} \ln \left(1 + \frac{t}{t_0} \right) \right]. \quad (4.11)$$

In the following we summarize some important conclusions from this equation. The first and most obvious one is that the term in brackets in Eq. (4.10) represents a correction to the current density; this is the “flux-creep reduction factor.” At low temperatures, where U_0 is essentially temperature independent, the flux-creep term introduces a reduction in J that is *linear* in temperature.

Equation (4.10) is valid only in the limit of $(kT/U_0) \ln(t/t_0) \ll 1$, that is, in the limit where the flux-creep reduction factor is small. One can define a crossover time t_{cr} at which $U_0 J / kT J_{c0}$ drops to unity; using Eq. (4.10), this leads to

$$t_{cr} = t_0 \exp \left(\frac{U_0}{kT} - 1 \right). \quad (4.12)$$

This is the crossover time between logarithmic and exponential relaxation (see Sec. III.B.3). Similarly, a crossover temperature T_{cr} between these two limits, for a given observation time t_{obs} is

$$\frac{kT_{cr}}{U_0(T_{cr})} = \frac{1}{1 + \ln(t_{obs}/t_0)}, \quad (4.13)$$

where we have explicitly indicated the temperature dependence of U_0 . However, in most high-temperature superconductors, U_0 is so small that T_{cr} falls in the range below $T_c/2$, where U_0 is essentially independent of temperature.

It has been shown that, in the entire region in which the logarithmic behavior holds, the critical-state model with J replacing J_c is a good approximation (Griessen *et al.*, 1990; Feigel'man *et al.*, 1991; Schnack *et al.*, 1992; van der Beek, Nieuwenhuys, *et al.*, 1992). This leads to a magnetization M whose magnitude is proportional to J when the sample is fully penetrated by the field. For example, in a slab of thickness L , with a sufficiently large field applied parallel to the slab plane, one obtains from Eq. (2.6)

$$|M| = JL/4c \quad (4.14)$$

for a constant flux-density slope $\pm J$ (Bean, 1964; Campbell and Evetts, 1972). Therefore the magnetization relaxes with the same logarithmic flux-creep correction as the current density, given by Eq. (4.10).

From Eq. (4.14) one can derive an expression for the normalized relaxation rate

$$S = \frac{1}{M} \frac{dM}{d \ln t} = \frac{d \ln M}{d \ln t} = \frac{d \ln J}{d \ln t}. \quad (4.15)$$

This quantity has the advantage that, when using it to evaluate the barrier energy U_0 , it is not necessary to know the value of J_{c0} appearing in Eq. (4.10), which is quite uncertain in view of the very large relaxation that can occur even before the first experimental observation. A minor problem is that the relaxation rate $dM/d \ln t$ is often normalized to the initial magnetization M_i rather than by the time-dependent magnetization $M(t)$.

But since the variation in M during the experiment is small in most cases, the error introduced by this procedure is also small. The result for S from Eq. (4.10) is

$$S = \frac{-kT}{U_0 - kT \ln(t/t_0)}. \quad (4.16)$$

In the limit in which Eq. (4.10) is valid, U_0 dominates $kT \ln(t/t_0)$ in the denominator, and S is always negative.

As was pointed out by Hagen and Griessen (1990), for low-temperature superconductors in which U_0 is very large compared to $kT \ln(t/t_0)$, S reduces simply to $-kT/U_0$, which is just the coefficient in front of the \ln term in Eq. (4.10). This has made it tempting for many authors in high-temperature superconductivity to report $-kT/S$ and call it a barrier energy. Equation (4.16) shows that such an effective barrier actually represents $U_{eff} = U_0 - kT \ln(t/t_0)$. Because of this ambiguity, the experimental literature must be read very carefully to know which of these barriers, whether U_0 or U_{eff} is actually being reported.

A close relation can be established between flux creep and the I - V (or E - J) characteristics of the superconductor. Since $E \propto vB$ and the flux velocity v is proportional to $\exp(-U/kT)$, one obtains

$$E \propto B \exp \left[-\frac{U_0}{kT} \left(1 - \frac{J}{J_{c0}} \right) \right]. \quad (4.17)$$

This shows that a linear dependence of U on J , which causes a logarithmic magnetic relaxation, is associated with an exponential V - I characteristic. Finally, we note that according to Eq. (4.5), $\partial J / \partial t \propto E$. This is an expression of the conservation of energy, since JE represents power dissipation and J^2 represents a stored energy, so that $\partial J^2 / \partial t \propto JE$.

C. Nonlinear $U(J)$

Beasley *et al.* (1969) already recognized that a more realistic barrier should exhibit a nonlinear dependence upon the current density J . The importance of a nonlinear $U(J)$ has recently become apparent in relation to high-temperature superconductors (see, for example, Feigel'man *et al.*, 1989; Zeldov *et al.*, 1990). The need for a nonlinear $U(J)$ first arose when a straightforward application of the conventional formula—Eq. (4.10)—led to barriers that increase rather than decrease with temperature (Xu *et al.*, 1989).

We next discuss several barriers that have received special attention. One is the barrier first discussed by Beasley *et al.* (1969), and more recently by Griessen (1991b) and by Lairson *et al.* (1991), who found that near J_{c0}

$$U_J \propto [1 - (J/J_{c0})]^{3/2}. \quad (4.18)$$

Of more recent interest for high-temperature superconductors are the barriers proposed for $J \ll J_{c0}$: An inverse power-law barrier (see, for example, Feigel'man *et al.*, 1989),

$$U = U_0 [(J_{c0}/J)^\mu - 1], \quad (4.19)$$

and a logarithmic barrier (Zeldov *et al.*, 1989, 1990),

$$U = U_0 \ln(J_{c0}/J). \quad (4.20)$$

Both of these barriers have the peculiar property of diverging as $J \rightarrow 0$. This divergence can be understood in the context of collective pinning and will be discussed more extensively in Sec. IV.C.1.

A rather general argument can be made for the behavior described in Eq. (4.18) with a $3/2$ power of $J_{c0} - J$: for any smoothly varying potential $U(x)$, the maximum slope occurs at an inflection point and defines J_{c0} . Expanding around this point, one therefore has, to lowest order, $U(x) = J_{c0} Bx - cx^3$, where c is some constant. Differentiating this potential, along with the current-density term JBx , one immediately finds the net barrier going as $(J_{c0} - J)^{3/2}$. It should be emphasized that this dependence dominates only where J approaches J_{c0} . Since most experiments in high-temperature superconductors probe values of J far from J_{c0} , the lack of this feature is not a serious problem for Eqs. (4.19) and (4.20) except, perhaps, for low temperatures, where the relaxation is slow (Griessen, 1991b).

1. Collective-creep theory

The inverse-power-law form of Eq. (4.19) has emerged from recent theories involving collective vortex pinning. We refer here mainly to the theoretical work of Feigel'man, Geshkenbein, Larkin, and Vinokur (1989). This theory assumes weak random pinning and treats the flux-line system as an elastic medium. Contrary to the original flux-creep model, where the volume V of the thermally activated flux bundle was constant, in the collective-creep model V depends on the current density J and becomes infinitely large for $J \rightarrow 0$. Consequently, as $J \rightarrow 0$ the activation energy U diverges and the flux system is frozen. The theory of collective creep (for finite J) has been extensively reviewed by Blatter *et al.* (1994). A central result of this theory is the so-called "interpolation formula":

$$J(T, t) = \frac{J_{c0}}{[1 + (\mu k T / U_0) \ln(t/t_0)]^{1/\mu}}, \quad (4.21)$$

where t_0 is the logarithmic time scale defined in relation to Eq. (4.7). This equation is obtained by equating the activation barrier given in Eq. (4.19) and Eq. (4.7). The factor μ in the denominator is introduced in order to interpolate between the usual Anderson formula (for $J_c - J \ll J_c$ at short times) and the long-time behavior.

It should be noted that Malozemoff and Fisher (1990) derived the same formula on the basis of the vortex-glass model (Fisher, 1989). In this model, the flux system undergoes a thermodynamic phase transition from a "vortex-liquid" state, where the vortices are highly mobile, into a "vortex-glass" state in which vortices are localized in a metastable state created by interactions with the pinning centers and the other vortices. In this phase, vortex motion is possible only in the presence of a current. This can be described in terms of a diverging creep barrier at low temperature and $J=0$.

In its simplest version (Fisher, 1989), the vortex-glass model predicts μ to be a universal exponent less than one, while Feigel'man *et al.* (1989) predicted a complicated dependence of μ on field and temperature. For example, in three dimensions $\mu=1/7$ in the low-field, low-temperature region where the creep is dominated by the motion of individual flux lines; at higher temperatures and fields $\mu=3/2$ due to collective creep of small bundles; finally at still higher fields and temperatures, where the bundle size is much larger than the London penetration length, $\mu=7/9$. For two-dimensional creep $\mu=9/8$. These results are valid for hopping distances much shorter than the Abrikosov lattice constant. Nattermann (1990) considered the opposite case and found $\mu=1/2$.

Since $M \propto J$, Eq. (4.21) immediately leads to a normalized relaxation rate

$$S = kT/[U_0 + \mu kT \ln(t/t_0)], \quad (4.22)$$

which is obviously different from the Anderson-Kim prediction, Eq. (4.16). Equation (4.22) predicts that the normalized relaxation rate decreases with time. Another interesting prediction of this equation is that, with increasing temperature, the second term in the denominator dominates U_0 , and S approaches the limit (Malozemoff and Fisher, 1990; Natterman, 1990)

$$S = 1/\mu \ln(t/t_0). \quad (4.23)$$

Unless μ is dependent on T , or t_0 is strongly dependent on T , this formula predicts that S will have a plateau, much as is observed in region III of Fig. 12. Data indicating the plateau in a variety of YBCO samples, both films and crystals, are summarized by Malozemoff and Fisher (1990). All these data show values for S that fall in the range of several percent (0.022–0.026 in the case of Fig. 12). A value of S in this range can be obtained from Eq. (4.23), assuming μ of order one, t of order 1000 s as in most flux-creep measurements, and t_0 of order 10^{-10} s. A more reasonable value of t_0 of order 10^{-6} s (Blatter *et al.*, 1994) gives a value for S of about 0.05.

We consider the ability to predict a plateau in $S(T)$ to be a major success of the collective-pinning theories. However, many problems still remain. For example, as already mentioned in Sec. III.B.1, recent data show a significant field dependence of the plateau; the plateau value for S in YBCO is shifted from 0.05 down to 0.015 with field increased from 0.2 T to 1.7 T. In the context of Eq. (4.23) such a low value of S requires $\mu > 2$; otherwise the ratio t/t_0 becomes implausibly large. A value of $\mu > 2$ is beyond the present predictions of either the vortex-glass theory or that of Feigel'man *et al.* (1989) (see also the discussion in Sec. III.B.1).

A more direct support for the collective-creep model emerges from long-time relaxation measurements by Thompson, Sun, and Holtzberg (1991) on a YBCO crystal. These authors were able to fit the deviation from logarithmic behavior with the collective-pinning formula, Eq. (4.21), as shown in Fig. 2. Of particular interest in this work is the first experimental effort to deduce the temperature dependence of μ . The inset to Fig. 2

shows that the μ values extracted from the fit to Eq. (4.21) show a pronounced peak around 30 K. Qualitatively this peak mimics the trends predicted by the collective-pinning theory (Feigel'man *et al.*, 1989), though the μ values at the peak are somewhat larger than predicted. Also, the temperature dependence of μ obtained in this work implies a strong temperature variation in S . This is, of course, in conflict with the experimental plateau.

Another point related to the prediction of Eq. (4.21) is the long attempt-time t_0 extracted from the data (Svedlindh *et al.*, 1991; Brawner *et al.*, 1993a). This was found to be "macroscopic," in the range of milliseconds to microseconds, much larger than the inverse of the actual attempt frequency expected to be in the nanosecond to picosecond range. In a recent work, Feigel'man, Geshkenbein, and Vinokur (1991) explained the macroscopic nature of the attempt times in the following way: while the magnetization is proportional to the sample volume, its rate of change is proportional to the surface area (any change in the total trapped flux involves vortices entering or leaving the sample through its surface). Consequently, t_0 depends on the sample size d as discussed in Sec. IV.A above [see also Blatter *et al.* (1994)].

Finally, an interesting consequence of the plateau in S should be noted. As was pointed out in Sec. III.B.6, magnetic relaxation is closely related to energy dissipation in the material and therefore manifests itself in I - V (current-voltage) or J - E (current density/electric field) measurements. The current-voltage characteristic is empirically well approximated by a power law: $E = \alpha J^n$, with a large exponent ($n \sim 10$ – 60) (Gilchrist, 1990). Sun *et al.* (1991) proposed a model in which the exponent n is directly related to the quantity U_0/kT measured in relaxation experiments. Assuming that the sample is in a critical state, they derived the relation $n = 1 + 1/S$. Thus a temperature-independent S implies a constant n in $E \sim J^n$. Sun *et al.* compared n values obtained from magnetic relaxation and from transport studies of YBCO thin films, and demonstrated a consistent trend in their temperature dependence. In particular, at the intermediate temperature range, both data sets show only weak temperature dependence.

2. Logarithmic barrier

The logarithmic J -dependent barrier of Eq. (4.20) was first proposed by Zeldov *et al.* (1990) in order to explain their magnetoresistance data. More recently, this prediction was supported by the extensive flux-creep studies of Maley *et al.* (1990), McHenry *et al.* (1991), and Ren and de Groot (1992) in aligned grains of YBCO and in LaSrCuO and YBCO crystals. In analyzing their data they started with the standard Arrhenius relation for thermally activated hopping and its effect on the time decay of current density. Based on Eq. (4.5), one has

$$dJ/dt \propto \exp(-U/kT). \quad (4.24)$$

From this equation they recognized that

$$U = -kT[\ln(dM/dt) + c], \quad (4.25)$$

with c a temperature- and field-dependent constant. Thus a plot of this quantity versus M (with $M \propto J$) should directly exhibit the J dependence of U . They found a dependence that supports the logarithmic dependence of Eq. (4.20)—see Fig. 17.

The procedure outlined by Maley *et al.* (1990) involves a determination of $c(T, H)$ by way of trial and error. The best values for c are considered to be those that lead to a "smooth" functional form of $U(J)$. In the original procedure described by Maley *et al.*, c was assumed to be constant at low temperatures. This was justified by the fact that c is a logarithmic function of slowly varying parameters of the diffusion equation. However, it should be noted that, in general, the field- and temperature-dependence of c is quite complicated; moreover, c is sample dependent (van der Beek, Kes *et al.*, 1992; van der Beek, Nieuwenhays, *et al.*, 1992). A scaling relationship for these dependences has been recently suggested by McHenry *et al.* (1991).

Further evidence, though indirect, for the logarithmic barrier comes from the widely reported quasiexponential dependence of the measured critical current density on temperature (Sennoussi *et al.*, 1988). As pointed out by McHenry *et al.* (1991), Eqs. (4.20) and (4.7) lead to

$$J = J_c \exp(-(kT/U_0) \ln(t/t_0)), \quad (4.26)$$

which explains this exponential temperature dependence for a fixed experimental time window.

The logarithmic law, Eq. (4.20), does not explain the plateau in S ; Eq. (4.26) predicts

$$S \equiv d \ln M / d \ln t = -kT/U_0, \quad (4.27)$$

with an upward-curving temperature dependence for S .

3. Barrier distribution

Many models for flux creep have been proposed involving a distribution of barriers (Malozemoff, Worthington, Yandrofski, *et al.* 1988; Hagen and Griessen, 1989; Ferrari *et al.*, 1990; Griessen, 1990, 1991b; Gurevich 1990; Martin and Hebard, 1991; Niel and Evetts, 1991). These kinds of models also offer an explanation of the plateau in the normalized relaxation rate. Using the standard Anderson-Kim treatment, a plateau in S implies a barrier that increases in a manner proportional to T . With an appropriate distribution of barriers, one can understand the increase of the average barrier with temperature, since the small barriers are overcome at lower temperatures. In fact, these models can explain arbitrary temperature dependences of magnetic relaxation; Hagen and Griessen (1989) have developed an inversion procedure for extracting the required distribution from the given temperature dependence. (For a recent example of utilization of the inversion procedure for YBCO and BSCCO crystals see Theuss, 1993 and Hardy *et al.*, 1994.) In effect, such models have an infinite number of parameters that can be fit to the data.

Indeed, a distribution of barriers in these complex materials is plausible and, although in a sense inelegant, is hard to argue against. Malozemoff and Fisher (1990) at-

tempted to do so on the basis of the S -plateau universality, pointing out the implausibility that barrier distributions in films, crystals, melt-textured rods, etc., would all be the same. Nevertheless, subsequent work has made the universality less clear, leaving the correct interpretation largely open.

V. CONCLUDING REMARKS

The many studies of magnetic relaxation phenomena in high-temperature superconductors have revealed properties that initially appeared unconventional in the context of low-temperature superconductors. These include unusually large relaxation rates, a nonlogarithmic decay, and complex field and temperature dependence of the relaxation. Subtle details of the relaxation in high-temperature superconductors have been accurately investigated over broad temperature and field ranges, thanks to the sheer size of the effect in these materials. Many of these interesting features are now becoming apparent.

The conventional model of thermally activated flux creep explains many of the relaxation data in high-temperature superconductors. However, because of the giant flux creep, the current density J changes considerably during the relaxation process and the nonlinear dependence of the activation energy U upon J must be taken into account. A proper description of $U(J)$ is still a challenge for both theory and experiment. Among the several models that have been proposed, the collective flux-pinning model is apparently an important step in deepening our understanding of the complex phenomena of flux creep in high-temperature superconductors. It predicts some of the main features observed, such as the deviation from logarithmic relaxation and the plateau in $S(T)$ in the intermediate temperature range. All in all, we judge this model to be most promising at present in integrating the largest amount of experimental data. However, not all the data are consistent with this picture either—for example, the relaxation data at ultralow temperatures call upon a different model involving quantum tunneling of vortices. Moreover, detailed predictions of the collective flux-pinning model, such as the different exponents μ in the various fields and temperatures, are not yet confirmed experimentally. At the same time, none of the competing models are able to explain the variety of the data.

While it is not possible to derive $U(J)$ directly from the relaxation data, each of the theoretical models gives a specific relaxation behavior that can be compared with the experimental results. Such an approach for evaluating $U(J)$ is model-dependent and involves fitting several parameters. The technique proposed by Maley *et al.* (1990) bypasses this difficulty by analyzing global magnetic-relaxation data on the basis of an integrated form of the flux-diffusion equation. However, because of the global nature of this technique, one actually measures the activation energy at the surface of the sample, while the current density J is averaged over the sample volume. The method recently proposed by Abulafia

et al. (1995 and 1996) allows direct, model-independent determination of the local U and J in the bulk by local magnetic measurements. The novel Hall-probe-array technique utilized in their experiment allows measurements of the time evolution of the field profile in the sample, thus enabling direct analysis of the flux creep on the basis of the flux-diffusion equation. This method is a great improvement compared to the indirect method available up to now to obtain the important characteristics of flux creep in high-temperature superconductors. It has the potential of giving a definite test of many as yet speculative models on collective pinning, collective creep, surface barriers, “fishtail” effects, and so on.

The remarkable development in high-temperature superconductors, both experimentally and theoretically, has prompted renewed activity in low-temperature superconductors, with researchers reexamining relaxation phenomena and other irreversible properties in light of the newly developed models (Yeshurun, Malozemoff, Wolfus, *et al.*, 1989; Mota *et al.*, 1990; Suenaga *et al.*, 1991; Tomy *et al.*, 1991; Schmidt *et al.*, 1993).

As we pointed out in Sec. III.B.6, there is an intimate relationship between magnetic relaxation and transport properties. Thus a more complete assessment of the various models should take into account transport data, which are not discussed in this review. We just note here that transport data reveal features that are consistent with a thermodynamic phase transition of the flux structure (Koch *et al.*, 1989; Charalambous *et al.*, 1992; Safar *et al.* 1992), though other interpretations, based on thermal activation, are still to be considered (Brandt *et al.*, 1989).

Finally, we note that many phenomena described in this review have implications for potential applications of high-temperature superconductors. In general, magnetic relaxation limits the stability of the persistent current density and affects the sharpness of the I - V curves. A question of a great practical importance is how to control the magnetic relaxation, in particular how to optimize the defect pinning. Perhaps the most promising approach has been that of irradiation using high-energy, heavy ions, which produce well-defined columnar defects. In addition to the purely empirical approach of introducing different pinning centers and measuring the relaxation, an important effort has been directed to understanding the nature of the pinning and, to this end, deducing information from the magnetic-relaxation data itself. But before this effort can be successful, there must exist an adequate phenomenological and microscopic theory by which to interpret the data and deduce useful information about the defects. This is where the field has foundered, because there has been no agreement, even about the phenomenological framework, not to mention a microscopic theory. At the same time, the profusion of novel theoretical concepts has added a richness and excitement to the field unequaled in the earlier development of low-temperature superconductivity. It is hoped that the present status report can provide a sound basis for further progress in this field.

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