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Magnetic relaxation, flux pinning and critical currents in superconductors

Karl S. Lichtenberger

Iowa State University

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Lichtenberger, Karl Sanford, Ph.D.

Iowa State University, 1991
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Magnetic relaxation, flux pinning and critical currents in superconductors

by

Karl S. Lichtenberger

A Dissertation Submitted to the Graduate Faculty in Partial Fulfillment of the Requirements for the Degree of DOCTOR OF PHILOSOPHY

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CHAPTER 1. INTRODUCTION

Overview

The discovery of the high temperature superconductors (HTSC's) [1] [2], with critical temperatures ($T_c$) as high as 125 K [3], gives wide $H$-$T$ planes that include the possibility of superconducting applications at liquid nitrogen temperature ($T = 77$ K). However, experiments have shown that the $H$-$T$ plane of the HTSC's have large reversible regions [4] [5] [6] [7], where there is no pinning of flux and the critical current density $J_c$ is zero, and there exists an even larger region where $J_c < 10^4$ A/cm$^2$, so that only a small region of the $H$-$T$ plane is useful for high current applications. Even in this region, there often exists a large field dependence on $J_c$ [8] [9] which severely limits the usefulness of the material in high-field applications. It is necessary to understand the reasons for the apparent weak flux pinning in these materials before they can be employed in practical applications.

Review of Superconductivity

The magnetic phase diagram of a superconductor has evolved over time, as the level of understanding and available materials has increased over the years. In Fig. 1.1, three different magnetic $H$-$T$ phase diagrams are shown, each corresponding to the state-of-the-art of different time periods. The first diagram is for superconductors
such as Pb or Sn, where the superconducting phase is separated from the normal state by the critical field $H_C$, and represents the understanding of superconductivity around 1950. The second diagram is for superconducting alloys such as Nb$_3$Sn or Nb-Ti. Here, there are three phase boundaries; the lower critical field $H_{c1}$, the thermodynamic critical field $H_c$, and the upper critical field $H_{c2}$. This represents the level of understanding of superconductivity around 1960. The third diagram is for a HTSC, such as YBa$_2$Cu$_3$O$_7-\delta$ or Tl$_2$Ba$_2$Ca$_2$Cu$_3$O$_{10}$, with four regions; fluctuations near $T_c$, thermodynamic reversibility, thermally-activated flux motion, and rigid flux lattice. A fourth field curve, the irreversibility transition, is added to the three found in the earlier figure. This is the present level of understanding of superconductivity. It is easily seen that the increase in the level of understanding has resulted in much more complicated $H$-$T$ diagrams. Before the $H$-$T$ diagrams can be understood, it is necessary to discuss some of the history and background of
Superconductivity was discovered in 1911 by H. Kamerlingh Onnes [10] [11].
when the resistivity of metals such as Hg and Pb was found to abruptly fall below
detectable levels (< 10^{-6} \Omega\cdot cm) at a certain temperature, $T_c$, called the transition (or
critical) temperature. The transition is very sharp, in contrast to HTSC's, in which
the transition is typically more gradual [12]. Superconductivity, however, can not be
explained by a perfectly conducting state. Meissner and Ochsenfeld [13] found that
superconductors exhibit perfect diamagnetism, where the susceptibility was $-1/(4\pi)$.
regardless of magnetic field and temperature history. Superconductivity was pre­
served in magnetic fields as long as the field did not exceed a critical field $H_c$, at
which the sample becomes normal [14]. This is illustrated in the first $H-T$ diagram
in Fig. 1.1. It became necessary to find theories that explained superconductivity, as
classical behavior could not explain superconductivity.

The classical behavior of superconductors was studied by Gorter and Casimir
[15], who attempted to use thermodynamics to explain superconductivity. One result
was the two fluid model (a normal electron and a superconducting electron fluid).
which is still often used today (for example, $\lambda(T)$ is calculated with this model in Ref.
[16]). A classical electrodynamic theory of superconductivity was given by London
and London [17], who showed that currents exist in superconductors in the absence
of an electric field. The penetration depth $\lambda$ was introduced to describe the decay
of magnetic fields and currents inside the superconductor. Non-local electrodynamic
effects were introduced by Pippard [18] who introduced the coherence length $\xi$ as a
means to explain the long range interaction of superconducting electrons. These are
classical theories, however, and are not satisfactory for understanding superconduc-
tivity. The addition of quantum mechanical effects was needed to better understand superconductivity.

The existence of a superconducting energy gap was shown in tunneling \[19\] and ultrasonic attenuation experiments \[20\], and the electron-phonon interaction was shown in isotope-effect experiments \[21\] \[22\] \[23\], and discussed theoretically by Bardeen \[24\] and Fröhlich \[25\]. The electron-phonon interaction was also seen in the electronic specific heat, in which the temperature dependence was found to be exponential in \(1/T\) \[26\].

The tunneling, ultrasonic attenuation, and isotope effect experiments were explained in detail by the theory of Bardeen, Cooper, and Schrieffer (BCS) \[27\]. The BCS theory is a very general theory in which the normal state electronic wave functions are mixed coherently at one point in space to form a *Cooper pair* wave packet having the dimension \(\xi\). This theory accounts for the electron-phonon interaction in which a phonon binds the Cooper pair. Three BCS parameters give all of the electrodynamic properties of the superconductor. The superconducting wave function is still constant in space, however.

To account for variations in the superconducting wave function, Ginzburg and Landau (GL) \[28\] derived a phenomenological theory which applies well only near \(T_c\). Here, \(\lambda\) is found to be related to the electromagnetic energy of the superconducting state, and \(\xi\) is due to quantum mechanical condensation into the superconducting state. Abrikosov \[29\] showed from the GL theory that if \(\kappa \equiv \lambda/\xi > \sqrt{2}\), then the superconductor can exist at fields greater than the thermodynamic critical field, and that in this case, the field is not totally excluded from the bulk of the sample but exists in quantized units of flux called vortices. The value of \(\kappa\) defines the *type* of the
superconductor. type I for $\kappa < \sqrt{2}$ and type II for $\kappa > \sqrt{2}$. The upper and lower critical fields ($H_{c1}$ and $H_{c2}$) were introduced here, and these are shown in the second $H-T$ diagram in Fig. 1.1. Gor'kov [30] later showed that the BCS and GL theories were equivalent. This model of the superconductor is often referred to as the GLAG model.

The $H-T$ phase diagram for an ordinary type II superconductor can be derived from GLAG theory. The lower critical field $H_{c1}$ separates the Meissner state ($B = 0$) from the mixed state, where flux quanta $\Phi_0 = \hbar c/(2e)$ penetrates in sample in vortices. The upper critical field $H_{c2}$ is the field at which vortices are first nucleated (on cooling from above $T_c$) or at which the superconductivity is quenched. Above $H_{c2}$, the sample is in the normal state.

Because the resistivity of a superconductor is zero (for practical purposes), there was (and is) much interest in finding applications for these materials. Elemental superconductors are poor candidates for high-current, high-field applications because of their low critical fields, so modest current densities would drive the sample into the normal state. It was necessary to find superconductors that can carry large current densities at high magnetic fields for large-scale applications.

Matthias et al. [31] discovered that Nb$_3$Sn became superconducting at $T_c = 18$ K, which at that time was the highest yet observed. Later, Kunzler et al. [32] found that Nb$_3$Sn could carry high current densities of up to $J_c = 10^5$ A/cm$^2$ at 8.8 T. This was the first report of a high $J_c$, high $H_{c2}$ material, raising hope that superconductors could be used for high field magnets. Other high field superconductors include several A$_3$B alloys (A = transition superconducting metal, B = non-transition metal or semiconductor (i.e., V$_3$Ga) and Nb-Ti. These materials are called hard su-
perconductors because of their high critical currents [33]. Superconducting high-field solenoids using either Nb₃Sn or Nb-Ti wire are now fairly common.

**Summary of Experimental Studies on High $T_C$ Superconductors**

Type I superconductors are not very useful for practical applications, and the low temperatures of type II superconductors require the use of liquid helium, restricting superconducting applications to very low temperatures. This all changed with the discovery of the high $T_C$ superconductor (HITSC) (La$_{1-x}$Ba$_x$)$_2$CuO$_4$-δ [1], making a whole new class of superconducting materials available. This material is a doped (with Ba) semiconductor, and the crystal structure was found [34] to be a perovskite structure $K_2NiF_4$ with oxygen defects. Substitution of Sr for Ba in this material raised $T_C$ up to 40 K [35]. Wu et al. [2] showed that by replacing the La with Y, $T_C$ increased to 90 K, with a different yet similar crystal structure [36]. The later discoveries of the Bi-Sr-Ca-Cu-O [37], Nd-Ce-Cu-O, [38], and Tl-Ca-Ba-Cu-O compounds [3] showed that all of these superconductors have layered structures, with either orthorhombic (YBa$_2$Cu$_3$O$_{7-δ}$) or tetragonal (Tl-Ca-Ba-Cu-O, Bi-Sr-Ca-Cu-O, and (Nd$_{1-x}$Ce$_x$)$_2$CuO$_4$) unit cells. The only HITSC that is not layered is Ba-K-Bi-O [39], which has a cubic structure. Also, this material is the only HITSC that has no Cu. Since these materials (excepting Ba-K-Bi-O) are layered, anisotropy in the superconducting properties was expected and observed [8] [16] [40]. Single crystals or samples with good preferred orientation [41] were needed to properly understand these properties.

Many of the results of experimental studies on HITSC's were strikingly different from corresponding results on conventional superconductors. These included the
upper critical field, resistive transitions, the existence of the irreversibility transition, and superconducting fluctuations near $T_c$. A standard technique for determining the upper critical field $H_{c2}$ is to measure the resistance vs. temperature in different magnetic fields [42], and to define the transition temperature $T_{c2} = T_c(H = H_{c2})$ as either the temperature at which $R = 0$ or the temperature where the resistivity is some percentage (10, 50, or 90%) of the normal state. When this was done on HTSC's (see Refs. [12] [43], [44] and [45]), two things were immediately apparent. First, $H_{c2}$ was anisotropic, being larger for the field parallel to the basal planes than for the field parallel to the c-axis. Defining $H_{c2}$ was difficult, as the transition spread out in high fields. Moodera et al. [41] clearly showed this by plotting $H_{c2}$ using three different $H_{c2}$ criteria (10, 50, and 90% of the extrapolated normal state resistivity), and the three curves did not coincide. Palstra et al. [46] [17] analyzed the data below the first deviation from normal state behavior using a thermally activated resistance model [48], in which $p = \rho_0 \exp(-U_0/kT)$. Semi-log plots of the resistance vs. $1/T$ were very linear, showing that this model is quite adequate to explain the resistance. This also showed that the sample is superconducting in this region, so that defining $H_{c2}$ by a criterion or as $R = 0$ is probably incorrect.

Similar results for $H_{c2}$ were found by Worthington et al. [49] using AC susceptibility measurements. $T_{c2}$ was defined by extrapolating the data from below $T_c$ to a point on the temperature axis where the susceptibility was zero. A linear fit with temperature is expected from Gor'kov's theory. However, in examining their data, it can be seen that there exists considerable rounding near $T_c$, so that the linear extrapolation results in a $T_{c2}$ that is too low. More recently, this same method has been used by Welp et al. [50] in analyzing DC magnetization data. However, Hao
et al. [51] has shown that the Gor'kov result should only be used very near $T_c$, and that the data in the region where the $M \propto T$ should be analyzed using GL theory for HTSC's at low reduced fields $B/B_{c2}$.

Salamon et al. [52] found a specific heat jump at $T_c$ qualitatively different classical superconductors. The onset of the the jump is not suppressed at the same $dH_{c2}/dT$ rate as found in the resistive measurements above. As the field increased, the only noticeable effect was that the size of the specific heat jump was diminished. Junod [53] also showed this effect. Athreya et al. [54], Golub and Finnemore [55] and Sanders et al. [56] have all shown, by using the reversible magnetization near $T_c$ to derive the thermodynamic variables needed, that large values of $dH_{c2}/dT$ occur, in accordance with the specific heat data.

The question remains of what is happening very near $T_c$ in the HTSC's which causes these peculiar effects. The rounding of the resistive transitions was studied by Freitas et al. [57] using the thermal fluctuation model of Aslamazov and Larkin [58]. This model shows that Cooper pairs can exist above $T_c$, giving rise to either an excess conductivity or diamagnetism, and was used to claim that the superconductivity in these materials (in this case, YBa$_2$Cu$_3$O$_{7-\delta}$) is 3D in nature rather than 2D or 1D [59]. Lee et al. [60] showed by using an YBa$_2$Cu$_3$O$_{7-\delta}$ sample that was extremely pure (it had no Curie-Weiss paramagnetic signal), these fluctuations could be observed at temperatures as high as 200 K.

For polycrystalline or sintered samples, the measurement of $J_c$ is complicated by the inhomogeneous behavior of the material. Dimos et al. [61] studied the transport properties of a bi-crystal of YBa$_2$Cu$_3$O$_{7-\delta}$, and found that $J_c$ was much larger in the two single crystalline domains of the sample than across the internal grain
boundary. Therefore, the measured transport $J_c$ of a sintered sample with many grain boundaries will probably be much lower than the actual current density that can exist within each grain. $J_c$ can be calculated from magnetization data using the Bean model [62] [33], although this requires making assumptions (such as a constant $J_c$ throughout the sample) that may or may not be valid for the experimental conditions.

Exponential dependence on both field and temperature has commonly been observed in magnetic studies [8] [16] [40], as well as anisotropy, with $J_c \parallel \gg J_c \perp$, where the noted orientation is the orientation of the field with respect to the $c$-axis. The exponential field dependence was observed by Finnemore et al. [63] in Nb$_3$Sn wire, but was also seen by Hsiang and Finnemore [64] in SNS Josephson junctions, opening the possibility of Josephson coupling between grains [65] [9]. Another transport result is field hysteresis in $J_c$ as observed by Watanabe et al. [66].

The existence of Cooper pairs in HTSC’s was proven by Gammel et al. [67], by imaging the flux line lattice. They were able to show that the quantum of flux in HTSC’s is $\Phi_0$, instead of $\Phi = h/e$, which would be expected for single particle superconductivity. The imaging was successful at 1.2 K, but failed at 77 K (no lattice was observed), opening the possibility of flux-line-lattice melting.

Müller et al. [4] observed an irreversibility transition (the irreversibility line) at which the zero-field-cooled and field-cooled magnetization signals coincide. This transition was found to go as $(1 - t) \propto H^{2/3}$, which led them to conclude that this was a superconducting glass transition, as this is the same dependence for spin glasses. Yeshurun and Malozemoff [5] also observed this, but analyzed the data using a thermally activated critical current model, and showed that $J_c = 0$ along this same curve (the sample is reversible when $J_c = 0$). The addition of this irreversibility line.
as well as the large $dH_{c2}/dT$, shows that the magnetic phase diagram for the IIITSC’s appears to be quite different than for conventional superconductors. A typical $H-T$ diagram [68] for a IIITSC is shown in the third $H-T$ diagram in Fig. 1.1, including the irreversibility line, the three critical fields, and each region (fluctuation, reversible, flux creep, and rigid lattice) of the $H-T$ plane is labeled.

There have been many studies of flux creep on HTSC’s, so only a few will be highlighted here. One clue that flux creep effects are important was reported by Malozemoff et al. [69], where the value of “$H_{c2}$” obtained from AC susceptibility experiments was frequency-dependent. The usual methods of creep analysis is to use the method of Beasley et al. [70] and/or Campbell and Evetts [71]. Xu et al. [72] found that using this model resulted in an activation energy $U_0$ that increases with temperature and field, a common result for $YBa_2Cu_3O_7-\delta$ (see, for example, Ref. [73]). Welch [74] explains this by arguing that the increase in $U_0$ with temperature is due to the non-linear dependence on the current density $J$. Xu et al. did most of the measurements in high fields, whereas much of the literature concentrates on low fields (i.e., Ref. [5] used $H = 1$ kOe). Hagen and Griessen [75] claim that these single-barrier models are inadequate to explain flux creep data, and they give a method by which the zero-temperature distribution of activation (or pinning) energies can be extracted from experimental data.

The goals of the work reported here are to understand the differences in the $H-T$ diagrams for IIITSC’s, to find out which models apply best and where in the $H-T$ plane to apply them.
Survey of General Superconductivity Theories

London Theory

In order to describe the electrodynamic properties of a superconductor, London and London used \((j_s)\) is the superconducting current density)

\[
\nabla \times (\Lambda j_s) = -\frac{1}{c} h .
\]

(1.1)

Using Maxwell's equations [76] and vector identities [77], the Londons derived the following equation for the magnetic field within a superconductor:

\[
\nabla^2 h = \frac{4\pi}{c^2 \Lambda} h .
\]

(1.2)

Here, all quantities are written in lower case to signify that they are local quantities.

The London penetration depth is related to \(\Lambda\), and is given by

\[
\lambda_0 = (c^2 \Lambda/4\pi)^{1/2} = (mc^2/4\pi\epsilon^2 n_s)^{1/2} .
\]

(1.3)

This describes the penetration of weak magnetic fields, such as in the case of a semi-infinite superconducting slab, where \(H = H_0 e^{-x/\lambda_0}\) for an applied field \(H_0\).

Equation 1.2 also describes the supercurrent \(j_s\) (just replace \(h\) with \(j_s\)), so this also means that \(j_s\) is limited to a region of depth \(\lambda_0\) from the surface. Thus, the London penetration depth can be thought of as either a decay length for the magnetic field or the depth to which supercurrents exist.

BCS Theory

The BCS theory [27] was the first theory to propose a microscopic model for the mechanism of superconductivity. The three important parameters are the Debye
temperature $\omega_D$, the density of states $N(0)$, and a superconducting parameter such as $T_c$ or the energy gap $\Delta(0)$. Based on these three parameters, BCS worked out an entire theory for superconductivity. Since this is a very detailed and complicated theory, the results for many of the measurable quantities are listed below.

$$kT_c = 1.13\hbar\omega e^{-1/N_0 V_0} \quad T_c \text{ in terms of }$$

BCS Hamiltonian parameters

$$\frac{2\Delta(0)}{kT_c} = 3.58 \quad \text{Energy gap in terms of } T_c$$

$$\frac{c_s - c_n}{c_n} \bigg|_{T_c} = 1.43 \quad \text{Specific heat jump at } T_c$$

$$\frac{c_s}{\gamma T_c} = 1.34 \left( \frac{\Delta(0)}{T} \right)^{3/2} e^{-\Delta(0)/T} \quad \text{low temperature}$$

$$\text{electronic specific heat}$$

$$\xi_0 = \frac{\hbar v_F}{\pi \Delta(0)} \quad \text{coherence length}$$

Since $\xi \propto 1/\Delta(0) \propto 1/T_c$, the HTSC's have much shorter coherence lengths than low-$T_c$ superconductors. This is important, as many of the superconducting properties depend on $\xi$.

**Ginzburg-Landau Theory**

The Ginzburg and Landau [28] (GL) theory for macroscopic superconductivity uses the free energy expansion of the complex order parameter (superconducting wave function) $\psi (|\psi|^2 = n_s)$. This theory, along with the extensions of Abrikosov [29], who showed that type II behavior occurs when $\kappa > \sqrt{2}$, gives expressions for many superconducting parameters, along with the temperature dependencies near
$T_c$. These are listed below [78]:

$$\lambda_{eff}(T) = \frac{mc^2}{4\pi e^2 n_s}$$

(1.9)

$$\xi(T) = \frac{\Phi_0}{2\sqrt{2\pi} H_c(T) \lambda_{eff}(T)}$$

(1.10)

$$\kappa = \frac{\lambda_{eff}}{\xi}$$

(1.11)

$$H_c = \frac{\Phi_0}{2\sqrt{2\pi} \lambda \xi}$$

(1.12)

$$H_{c2} = \frac{\Phi_0}{2\pi \xi^2} = \sqrt{2} \kappa H_c$$

(1.13)

$$H_{c1} = \frac{\Phi_0}{4\pi \lambda^2} (\ln \kappa + 0.5)$$

(1.11)

Here, $\Phi_0 = \hbar c/2e$ is the flux quantum. The factor of two in the denominator is due to the charge of the Cooper pair (two electrons).

The effective penetration depth $\lambda_{eff}$ and coherence length $\xi$ depends on the purity of the sample. The "clean" limit is when $\xi \gg l$, where $l$ is the mean free path of the metal (in the free electron gas approximation [79]), and the "dirty" limit is the opposite case, $l \gg \xi$. Clean superconductors are samples with high purity, as the mean free path of a metal is greatly shortened by impurities. In fact, adding enough of an impurity to a type I superconductor results in a type II superconductor, i.e., Pb-Tl [70]. The coherence length is related to the Pippard length $\xi_0$ in the clean limit by $\xi(T) = 0.71\xi_0/(1 - t)^{1/2}$, and in the dirty limit by $\xi(T) = 0.855(\xi_0 l)^{1/2}/(1 - t)^{1/2}$. The penetration depths are given by $\lambda_L(t) = \lambda_L(0)/(2(1 - t))^{1/2}$, clean limit, where $\lambda_L$ is the London penetration depth, and $\lambda_{eff}(t) = \lambda_L(t)(\xi_0/1.33l)^{1/2}$, dirty limit. The GL parameter $\kappa$ is roughly independent of temperature, and is given by $\kappa = 0.96(\lambda_L(0)/\xi_0)$, clean, and $\kappa = 0.715(\lambda_L(0)/l)$, dirty.
General concepts relating to Flux Pinning

Vortices are pinned in superconductors by grain boundaries, impurities, or other defects because of dirty limit behavior. The free energy per unit length is $\Delta F/l = H_C^2/8\pi(\pi\xi^2)$. If the superconductor is dirty with bulk mean free path $l_0$, and a defect has mean free path $l_1 < l_0$, then the difference in the free energy per unit length between the defect and the bulk is

$$\Delta F/l = \frac{H_C^2}{8\pi} \pi \Delta(\xi^2) \propto l_1 - l_0 < 0 . \quad (1.15)$$

Since the difference is negative, the vortex will remain on the defect unless forced off.

The pinning force per unit length $f_p$ is given by [48]

$$f_p = \frac{H_{c1} \Phi_0}{4\pi} \frac{\ln(d/\xi)}{\ln(\kappa)} \quad (1.16)$$

In the II-TSC's $\kappa$ is very large and $H_{c1}$ is very small, so it can be seen that $f_p$ will be smaller than for a conventional superconductor.
CHAPTER 2. THEORY-FLUX CREEP MODELS

Single-Barrier Model

Overview

The single barrier flux creep model and the critical state model are used in this work for the majority of the data analysis. The critical state model gives a method for calculating the critical current density $J_c$ from magnetization data without using a transport measurement. The flux creep model gives expressions for calculating various material-sensitive parameters, such as the the depth of the pinning well $U_0$, the activation volume $V$, and the flux hopping distance $X$, the last two parameters almost always appearing as the product $VX$. A review of these two models is needed in order to show the assumptions made and the limitations that are imposed by these assumptions.

Critical State Model

Most of the data presented here were taken in regions of the $H$-$T$ plane where magnetic flux fully penetrates the sample and the flux density is approximately uniform across the sample. In this case, the Bean [62] [33] model takes on a particularly simple form. $J_c$ can be assumed to be independent of field, because the field is nearly uniform across the sample. Using this assumption, Bean was able to calculate the
magnetization of a hard superconductor as a function of field, and compare the result to experimental data. In the following expressions, cylindrical symmetry is assumed and practical units are used (Gauss, A/cm², and cm).

As the field is increased from the initial unmagnetized state, a surface emf (voltage) is generated by Faraday’s Law [76], \( E = -d\Phi/dt \), where \( \Phi \) is the magnetic flux, and surface currents are induced. If \( H \leq H_{c1} \), the currents are screened out to a depth \( \lambda \ll R \), so the magnetization is \(-4\pi M = H\) and the internal field \( H_i = 0\).

As \( H \) increases to values greater than \( H_{c1} \), the currents will flow down to a depth deeper than \( \lambda \). This new penetration depth can be calculated using Ampere’s law

\[
\nabla \times \mathbf{H} = \frac{4\pi}{10} J_c.
\]

(2.1)

The currents are such that the field will be tangential to the surface, which allows the curl to be re-written as \( dH/dr \). Since \( J_c \) is assumed to be constant, the left side of eq. 2.1 is a constant, and can be easily integrated from \( H_{c1} \) to \( H \) to yield the field-dependent penetration depth \( \Delta \).

\[
(H - H_{c1}) = \frac{4\pi}{10} J_c \Delta
\]

(2.2)

\[
\Delta = \frac{10}{4\pi} \frac{(H - H_{c1})}{J_c}
\]

(2.3)

The physical meaning of \( \Delta \) can be seen in Fig. 2.1, which shows the magnetic field profile (more properly called the magnetic flux density profile) as a function of depth. It can be clearly seen that when \( H < H^* \), \( \Delta \) is the depth to which the currents flow, and that \( H^* \) is the field at which the flux fully penetrates to the center of the sample (\( \Delta = R \)). Bean integrated this flux profile to derive the initial magnetization of the sample.
Figure 2.1: Bean model flux profiles of a cylindrical sample with radius $R$

The flux density profiles in Fig. 2.1 are commonly referred to as flux fronts. While Fig. 2.1 seems to imply a continuous flux distribution in the sample, it is well known that flux in type II superconductors is quantized in vortices of flux $\Phi_0 = \hbar c/2\pi$. These two views can be reconciled by viewing the flux front as a gradient of vortices, so that near the surface, where the field is close to the applied field, there exists a much larger density of vortices than near the center (or at $r = \Delta(H)$, if $H < H^*$) [80]. In high fields $H \gg H^*$, the flux front is much more near to being flat, that is, $B$ is nearly uniform across the sample, so the assumption of a uniform $J_C$ across the sample is not necessarily bad.

The Bean model is much more commonly used to calculate the critical current density from magnetization data than for describing the shape of magnetization curves. This method of calculating $J_C$ from magnetization data was presented by Fietz and Webb [81].

Consider a cylindrical sample with radius $R$. In the course of a hysteresis loop, the sample will be in an external field $H_a$. $H_a \gg H_{c1}$, so $H_{c1}$ can be neglected.
Figure 2.2: Bean model flux profiles for an external field $H_e$ with both increasing and decreasing field, as well as for zero field after a field cycle. The sample is a cylinder of radius $R$. 
in the analysis. The case for which \( H = H_a \) with increasing and decreasing field histories is shown in Fig. 2.2, in which the flux profiles (that is, \( B \) vs. \( r \)) are plotted vs. radial position. The slope of each profile is related to \( J_c \) via

\[
\frac{\partial H}{\partial r} = \frac{4\pi}{10} J_c .
\] (2.4)

In high fields, \( \partial B/\partial H \sim 1 \), which results in

\[
\frac{\partial B}{\partial r} = \frac{\partial B}{\partial H} \frac{\partial H}{\partial r} \approx \frac{4\pi}{10} J_c .
\] (2.5)

Equation 2.5 gives the magnitude of the slope. When the field is increased to \( H_a \), the currents act to screen out the flux from the interior, giving the downward slope in Fig. 2.2. If the field is decreased from above \( H_a \), the direction of the current reverses because the sense of the surface emf will also reverse. This leads to flux being trapped in the interior (also referred to as the remanent flux density \( B_r \)), the amount of which is related to \( J_c \). This gives the positive slope profile in Fig. 2.2.

The flux density \( B(r) \) can then be written as

\[
B_\mp (r) = B_s \pm \frac{4\pi}{10} J_c (R - r) ,
\] (2.6)

with the subscript referring to whether the field is increasing or decreasing. At the surface, \( B_s = H_a + 4\pi M_{eq} = B_{eq}(H_a) \), because the parallel component of \( H \) at the surface must be continuous \[82\], and for \( r > R \), \( B = B_{eq}(H_a) \), where \( B_{eq}(H_a) \) is the value that \( B \) would have in the absence of flux pinning.

At the same \( H_a \), the magnetization for each profile is given by

\[
-4\pi M_+ (r) = H_a - B_+ (r)
\] (2.7)

\[
-4\pi M_- (r) = H_a - B_- (r) .
\] (2.8)
and the difference (the hysteresis) is given by \( \Delta (-4\pi M) = B_- - B_+ \). Since any measured quantity will be an average over the cross-sectional area of the sample, this quantity must also be averaged over the cross-sectional area of the sample in order to calculate \( I_c \). Since \( B_- - B_+ = 8\pi/10J_c(R - r) \) from Eq. 2.4, \( \Delta (-4\pi M) \) is given by

\[
\Delta (-4\pi M) = \frac{1}{\pi R^2} \int_{r<R} d^2 r \frac{8\pi}{10} J_c (R - r) = \frac{8\pi}{30} J_c R .
\]

which is often written without the \( 4\pi \) as

\[
\Delta M = \frac{RJ_c}{15} .
\]

Thus, the Bean critical state model can be used to calculate the critical current density from magnetic hysteresis measurements, without the need for transport current measurements.

This model assumes that \( J_c \) is constant across the sample. In most superconductors at low magnetic fields, this is not the case. To account for this several models have been proposed, including that of Kim et al. [83], Yeshurun et al. [81], and Shi et al. [85], where some field dependent form for \( J_c \) is assumed. In fact, Xu et al. [86] have derived a generalized critical state model using a two-parameter expression for \( J_c \), from which all of the previous models can be derived. Most of these models are useful mostly at low fields, whereas the data in this work were taken at high fields. Therefore, little effort was put forth to utilize these models.

The problem with using the Bean approach is that if \( J_c \) is field-dependent, the slope of the field profiles in Fig. 2.2 will not be constant. This is because the local field changes with depth. This means that Eq. 2.6 must be re-written as

\[
B_\mp (r) = B_s \pm \frac{4\pi}{10} J_c (r) (R - r) .
\]
and that the integration in Eq. 2.9 must now take in the spatial dependence on $J_c$. This is a non-trivial exercise, since in order to calculate $J_c(r)$, it is necessary to know the dependence of the local field on position, which is often not known. To get around this is, either only high field data are used, so that $B \sim H$, and $J_c$ is constant across the sample, or fine-grained samples ($R \sim \lambda$) are studied [65].

Kim - Anderson Model

The critical state model gives a method to obtain the critical current from experimental data, but does nothing on its own to account for the rather severe temperature dependence of $J_c$, even at low temperatures. Anderson [87] showed that $J_c$ is limited at finite temperature by the thermal activation of bundles of fluxons over flux pinning cites such as defects and impurities. This model predicts a logarithmic decay in $J_c$ with time, which was observed by Kim et al. [83] in measurements on Nb-Zr tubes. This means that $J_c$ will eventually decay to zero (the sample will revert to the equilibrium state), but the time required for this to happen is extremely long [11].

While it is possible that the resulting flux creep is due to the activation of single vortices, this was considered unlikely, because the field of a vortex will extend out a distance $\lambda$ into bulk of the sample, which is typically $5 \times 10^{-6}$ cm (or even longer for HTSC's [88] compared to conventional superconductors [87]). Therefore, it is unlikely that in a hard superconductor that there will not be any overlap and interaction between vortices. The flux bundles are assumed to have a size of $d \sim 10^{-6} - 10^{-5}$ cm, and the bundles can independently creep the same distance $d$.

When a flux bundle encounters a defect or impurity, there will be a difference in the free energy between the superconducting bulk material and the normal (or at
least less superconducting) defect. If the defect is normal, then the difference in free energy in a volume of \( d^3 \) is

\[
\Delta F_{\text{max}} = F_n - F_S = \frac{H_c^2}{8\pi} d^3.
\]  

(2.12)

This assumes that the entire bundle is pinned, which is unlikely. To account for this, Anderson introduced a parameter \( p \) which is the effective fraction of pinning (if the bundle is pinned over a distance \( r \), then \( p \sim r^2/d^2 \)). The actual free energy difference is then

\[
\Delta F^* = p\Delta F_{\text{max}} = p \frac{H_c^2}{8\pi} d^3.
\]  

(2.13)

In the critical state, there exists a current density \( J_c \) which is flowing throughout the bulk of the sample (here it is implicitly assumed that \( H > H^* \)). The Lorentz force (per unit volume) is \( J_c \times B \). If this is integrated over the volume of the bundle \( d^3 \) then the free energy contribution will be the Lorentz force times the hopping distance, that is,

\[
F_L = -J_c B d^1 = -J_c \Phi d^2.
\]  

(2.11)

where it is assumed that the area is \( d^2 \), so the flux \( \Phi = B d^2 = n\Phi_0 \), where \( n \) is an integer, and \( \Phi_0 = hc/2e \) is the flux quantum. Therefore, the free energy is given by

\[
\Delta F^* = p \frac{H_c^2}{8\pi} d^3 - J_c \Phi d^2.
\]  

(2.15)

This free energy difference describes the pinning potential \( U \) of the superconductor. In the case of zero driving force (no currents and/or fields), \( U = pH_c^2/(8\pi)d^3 \equiv U_0 \). If there exists a periodic array of pinning cites (i.e., \( \alpha \)-Ti precipitates in Nb-Ti), then \( U_0 \) would look something like the potential shown in Fig. 2.3, which looks like an old-fashioned washboard. If the driving force is non-zero, however, then the potential will resemble Fig. 2.1, which is the the potential in Fig. 2.3 \((U_0)\) with the term
Figure 2.3: A hypothetical non-tilted pinning potential

Figure 2.4: A tilted-washboard potential
—|F|V.X superimposed on it. The appearance of this potential gives it the name tilted washboard potential. In the presence of external fields, a vortex pinned in a well will be more susceptible to thermal activation over the edge of the well, as the tilting due to the external driving force causes the effective depth of the well to be reduced. This approximation to the activation energy is called the linear approximation, because it is assumed that $U \propto J$.

For the thermal activation of the flux bundle, the Arrhenius relation [48] is used, with the flux hopping frequency given by

$$R = R_0 e^{-\Delta F/kT},$$

where $R_0$ is an attempt frequency. Anderson then used this to show that the critical state parameter $\alpha(t) = J(t)[B(t) + B_0]$ will decay logarithmically in time.

**Beasley-Labusch-Webb Model**

The model of Beasley et al. [70] is built around the assumptions of the Bean and Kim-Anderson models, so the amount of mathematical detail here will be curtailed, with the emphasis being on the physical meaning of the formulae. This model will be referred to as the BLW model. In Table 2.1, the different notations of Beasley et al. and Anderson are shown. The notation of used in the BLW model will be used from now on.

While Anderson used $(1/c)\mathbf{J} \times \mathbf{B}$ for the Lorentz force, Friedel et al. [89] found that a more accurate expression is

$$\mathbf{F} = \frac{1}{c} \mathbf{J} \times \mathbf{B} \frac{\partial H(B)}{\partial B},$$
<table>
<thead>
<tr>
<th>Quantity</th>
<th>Anderson</th>
<th>Beasley</th>
</tr>
</thead>
<tbody>
<tr>
<td>Activation energy</td>
<td>$\Delta F^*$</td>
<td>$U$</td>
</tr>
<tr>
<td>Well depth</td>
<td>$pH_C^2/(8\pi)d^3$</td>
<td>$U_0$</td>
</tr>
<tr>
<td>Pinning length</td>
<td>$d$</td>
<td>$X$</td>
</tr>
<tr>
<td>Activation volume</td>
<td>$d^3$</td>
<td>$V$</td>
</tr>
<tr>
<td>Lorentz force</td>
<td>$</td>
<td>1/\rho</td>
</tr>
<tr>
<td>Flopping frequency</td>
<td>$R$</td>
<td>$\nu$</td>
</tr>
</tbody>
</table>

Table 2.1: Identification of the variables used by Beasley et al. and how they relate to those of Anderson

so in one dimension, $F = -\gamma/(4\pi)B\nabla B$, if the Maxwell equation $j = (c/\mu_B)\nabla \times B$ is used. In high-$\kappa$ materials at high fields, $\gamma = \partial H/\partial B \approx 1$, so the difference in these two Lorentz forces is minor.

To derive the flux creep equation, BLW starts with an expression for the flux flow density $D$, which is the amount of flux per unit length and time that crosses a line perpendicular to $B$ and $\nabla B$.

$$D = \frac{-\nabla B}{|\nabla B|} Bwv_0 e^{-U(B, \nabla B)/kT}$$

In Eq. 2.18, the first term is a unit vector perpendicular to both $B$ and $\nabla B$, and the thermal activation term explicitly contains the dependence of $U$ on fields and gradients (currents). The distance $w$ is the average distance moved in a flux jump.

Using the requirement of flux conservation, and assuming that $U = U_0 - |F|V_X \ll kT$, BLW derived the following expression for the flux creep rate $R = d\Phi/d\ln t$.

$$R = \pm \frac{\pi}{3} kT \rho^3 \left( \frac{\partial U}{\partial |\nabla B|_{\rho}} \right)^{-1} (1 \pm \delta)$$

$$\delta = \rho \left| \nabla B \right|_{\rho} \frac{\partial}{\partial B} \ln \left[ \frac{1}{|\nabla B|_{\rho}} \left( \frac{\partial U}{\partial |\nabla B|} \right)_{\rho} \right]$$
Figure 2.5: Figure from Beasley et al. showing the dependence of $U$ on $J_c$. Using the linear approximation gives $U_0$ as the intercept of the tangent to the curve at the point shown.

The ± sign refers to whether the field history is increasing or decreasing. In practice, $\delta$ is eliminated by using $R_{av} = (R_+ + R_-)/2$.

To calculate $U_0$ and $VX$ from creep data, all that is needed is the dependence of $U$ on $|\nabla B|$ ($J_c$). Figure 2.5 shows a typical non-linear pinning potential function. If the linear approximation is used, then

$$\left(\frac{\partial U}{\partial |\nabla B|}\right)_{t=1, \rho} = \frac{U_0}{\frac{4\pi}{10} J_c}.$$  \hspace{1cm} (2.21)

This gives the result

$$R_{av} = \frac{\pi}{3} kT \rho^3 \frac{4\pi}{10} \frac{J_c}{U_0}.$$  \hspace{1cm} (2.22)

In the current work, the observed quantity is the average magnetization $M$. Since the flux $\Phi = BA = B(\pi \rho^2)$ and $B = H + 4\pi M$, $S \equiv dM/d\ln t = R/((1\pi)(\pi \rho^2))$, so

$$S \equiv \frac{dM}{d\ln t} = \frac{\rho J_c kT}{30 U_0}.$$  \hspace{1cm} (2.23)
This gives the material-sensitive parameter $U_0$ in terms of the experimentally measurable quantities $S$ and $J_c$.

$U_0$ is an energy, which is also a force times a distance. Because the force per unit volume is $|J_c \times B|$, the activation volume is $V$, and the pinning length is $X$. $J_c BVX$ is also an energy, so Eq. 2.2.3 can be written in terms of $VX$ as

$$S = \frac{1}{3} kT \rho \cdot \frac{1}{\gamma BVX}.$$  \hspace{1cm} (2.2.4)

This is the BLW model for flux creep, showing how the materials-sensitive parameters $U_0$ and $VX$ can be determined from the analysis of flux creep data.

**Distribution Model**

From the inhomogeneous character of the HTSC's, it was reasonable to believe that a distribution of pinning sites would be appropriate. So in 1989, Hagen and Griessen (HG) [75] worked out a method to extract this distribution from flux creep data. This model is an extension of the thermally activated $J_c$ model of Campbell and Evetts [71], where the normalized creep rate \( (1/M) dM/d \ln t = S/M = d \ln M/d \ln t \) can be written as

$$\frac{d \ln M}{d \ln t} = \left[ \frac{U(T)}{kT} - \ln \left( \frac{t_b}{\tau} \right) \right]^{-1}.$$ \hspace{1cm} (2.2.5)

where $t_b$ is the time at which creep data is taken in the critical state (i.e., $t_1$ in the BLW model), $\tau$ is a characteristic time constant for the creep process (i.e., $1/\nu_0$ in the BLW model), and $\ln t_b/\tau$ governs the initial collapse of the critical state. The magnetization that is being calculated here is not the measured signal, but is the magnetization due only to the critical current. This is an important difference, because
the equilibrium magnetization is stable in time and should not decay. For comparison to the experimental results, the equilibrium magnetization must be subtracted.

In this model, the basic assumption is that there exists a continuous distribution of independent activation energies $E$, each of which decays according to

\[ i(t, r) = iU_0(t) \left( 1 - \frac{kT}{E} \ln(1 + \frac{t}{\tau}) \right). \]  

The total magnetization is given by summing the individual magnetization contributions to obtain

\[ M(t, T) = M_0(T) \left( 1 - \frac{kT}{E} \ln(1 + \frac{t}{\tau}) \right). \]  

where $M_0$ is the value of $M$ at $T = 0$, $b(T)$ is the temperature dependence of the activation energies, $a(T)/b(T)$ is the temperature dependence of the critical current, and

\[ E_0^* = kT/b(T) \ln(1 + t/\tau) \approx kT/b(T) \ln(t_b/\tau) \]  

is the activation energy, which is assumed to be a continuous variable. Here, the energy is evaluated at $t = t_b$, and because $t_b/\tau \gg 1$, the argument of the logarithm can be simplified. Note that if a single barrier distribution $m(E^*) = \delta(E^* - t_0)$ is substituted into Eq. 2.27, Eq. 2.26 is recovered, as would be expected.

Differentiating Eq. 2.27 with respect to $\ln t$ results in

\[ \frac{dM}{d\ln t} = -M_0 \frac{kT}{a(T)} \int_{E_0^*(t_b, T)}^{\infty} m(E^*) \frac{m(E^*)}{E^*} dE^*. \]  

The distribution $m(E_0^*)$ is obtained by differentiating with respect to temperature.

\[ m(E_0^*) = \left[ \frac{d}{dT} \left( \frac{a(T)}{kTM_0} \frac{dM}{d\ln t} \right) \right] \left[ \frac{b(T)}{T} \frac{d}{dT} \left( \frac{T}{b(T)} \right) \right]^{-1}. \]
The model is not yet complete, however, as \( \ln \frac{t_b}{\tau} \) is as yet unknown. Some further manipulations of Eqs. 2.27 - 2.29 yields the following for \( \ln \frac{t_b}{\tau} \).

\[
\ln \left( \frac{t_b}{\tau} \right) = T \frac{b(T)}{a(T)} \left( \frac{d}{dT} \left[ \frac{M(T, t_b)}{b(T)} a(T) \right] \right) \left[ \frac{dM}{d\ln t} \left( 1 - \frac{d \ln b(T)}{d \ln T} \right) \right]^{-1} \tag{2.31}
\]

This term is assumed to be independent of temperature, so the proper choice of \( a(T) \) and \( b(T) \) should be made with this in mind. Equation 2.31 can be applied to data, so that \( \ln \frac{t_b}{\tau} \) can be experimentally determined.

These two equations 2.30 and 2.31 are the heart of the HG distribution model. Equation 2.31 gives the activation energy \( E_0^* \) (Eq. 2.28) and Eq. 2.30 gives the distribution function \( m(E_0^*) \), provided that the functions \( a(T) \) and \( b(T) \) are known.

The function \( b(T) \) is the temperature dependence of the activation energy as mentioned before. In the Anderson-Kim model, this was given by \( H_c^2/(8\pi) d^3 \). To generalize this to include the possibility that the coherence length is contributing to the creep process, this can be written as \( H_c^2/(8\pi) \xi^n d^{3-n} \), where \( n \) is an integer satisfying \( 0 \leq n \leq 3 \). Using the two-fluid model for \( \lambda^2 = \lambda^2(0)/(1 - \Theta^1) \) and the parabolic dependence for \( H_c = H_c(0)(1 - \Theta^2) \), where the reduced is temperature \( \Theta = T/T_c \) (\( \Theta \) is used here to prevent confusion with time \( t \)). This gives \( \xi \propto ((1 + \Theta^2)/(1 - \Theta^2))^{1/2} \), so that \( b(\Theta) \) is given as

\[
b(\Theta) = (1 - \Theta^2)^2 \left( \frac{1 + \Theta^2}{(1 - \Theta^2)} \right)^{n/2} \tag{2.32}
\]

The model of Tinkham [90] corresponds to the case of \( n = 1 \).

The other function \( a(T) \) is the temperature dependence of the Lorentz-force term \( V_X \). Using the same notation as for \( b(T) \), \( a(T) \) is expected to vary as \( \xi^m \), where
0 ≤ m ≤ 4, so $a(\Theta)$ is given by

$$a(\Theta) = \left[\frac{1 + \Theta^2}{1 - \Theta^2}\right]^\frac{m}{2}.$$ \hfill (2.33)

If $m = 0$, then the microstructure of the sample is the controlling factor the Lorentz-force term [75], but if $m = 4$, then both the activation energy and the pinning length as controlled by the coherence length. If the same volume is assumed for both $b(T)$ and $a(T)$, then the values of $n$ and $m$ are not independent of each other, that is, $m ≤ n$.

Equations 2.30 - 2.33 comprise the entire IIG distribution model. All that remains to be done is to find the temperature dependencies of $dM/d\ln t(T)$ and $M(T)$ and fit them into these equations. The exponents $n$ and $m$ are chosen in such a way as to make Eq. 2.31 independent of temperature.
CHAPTER 3. EXPERIMENTAL

SQUID Magnetometer

Magnetization measurements were performed in two commercial Quantum Design SQUID magnetometers ("QD1" and "QD2"). Both machines are computer-controlled, with full temperature and field control capability. The magnetometer QD1 is an older model of the basic QD magnetometer, with a maximum field of 2 T, while QD2 is a newer magnetometer with maximum field of 5.5 T.

One problem with these magnetometers is that the magnets are not homogeneous along the solenoid axis. If the sample exhibits hysteresis (i.e., ferromagnets and superconductors), the resulting magnetic moment will not necessarily reflect a characteristic of the sample but will reflect the response of the sample to a field gradient. Therefore, it is necessary to know the variation of the magnet along the length of the solenoid, and to limit the scanning length of the sample to a region

<table>
<thead>
<tr>
<th>Scan Length (cm)</th>
<th>Field Variation (%)</th>
</tr>
</thead>
<tbody>
<tr>
<td>1.0</td>
<td>0.001</td>
</tr>
<tr>
<td>2.0</td>
<td>0.005</td>
</tr>
<tr>
<td>3.0</td>
<td>0.048</td>
</tr>
<tr>
<td>1.0</td>
<td>0.19</td>
</tr>
<tr>
<td>6.0</td>
<td>1.1</td>
</tr>
<tr>
<td>8.0</td>
<td>5.8</td>
</tr>
</tbody>
</table>

Table 3.1: Scan lengths with the corresponding field variations for QD2
Figure 3.1: Field variation vs. position for the magnet used in QD1. The scan length is twice the position coordinate where the field has a small variation. The manual for QD2 gives the field variation for different scan lengths, shown in Table 3.1. No such information existed for QD1, however, so Steve Sanders and Jerry Ostenson of Ames Lab mapped out the field of this magnetometer by sweeping a coil connected to a Walker MF-30P integrating fluxmeter along the solenoid axis. For QD1, the operating software can calculate the magnetic moment for scan lengths no shorter than 3.5 cm, so for a 3.5 cm scan, the field variation was found to be 0.2%. Since the maximum field is only 2 T, the largest field variation is only 40 G, which is not considered to be much of a problem.

Temperature is controlled using a QD R-G bridge, heating provided by a wire-wound heater. Sample cooling is achieved by turning the heater off and drawing cold He gas into the sample chamber, which results in a large temperature undershoot. For measurements that are sensitive to temperature history, it is necessary to use a warming curve.
The magnetic field control is charged with a Kepco JQE 6-15 (M) (6 volts, 45 Amps) (QD2) or JQE 6-22 (QD1) power supply. The standard magnet charging sequence is to overshoot the set point field by approximately 20% and progressively oscillate the field down to the set point in approximately 6 cycles. For hysteretic samples, this method is unacceptable, so the no-overshoot mode can be utilized, in which the magnet is monotonically charged. When the set point is reached, a persistent current switch is shorted across the magnet terminals, and once the field is stable, the power supply is shut off. To change the field, the power supply must first be reset to the previous setting to prevent a field quench. Negative fields are obtained by use of a current reversing switch.

Samples

The primary focus of this work was to study the intragranular flux creep in a highly anisotropic superconductor Tl$_2$Ba$_2$Ca$_2$Cu$_3$O$_{10}$ and compare the results to other high $T_c$ and conventional superconductors. There was special motive in studying each sample, which are shown in Table 3.2. Tl$_2$Ba$_2$Ca$_2$Cu$_3$O$_{10}$ has the highest $T_c$ and probably has the most anisotropic structure of the HTSC’s. Nb-Ti and Nb$_3$Sn are conventional superconductors which are known to exhibit large flux pinning and critical currents. Bi$_2$Sr$_2$CaCu$_2$O$_x$ is similar to Tl$_2$Ba$_2$Ca$_2$Cu$_3$O$_{10}$ in structure, and should therefore have similar behavior. (Ba$_{1-x}$K$_x$)BiO$_3$ is the only HTSC that has a cubic structure, and (Nd$_{1-x}$Ce$_x$)$_2$CuO$_4$ is the only electron-carrier HTSC. By comparing these different superconductors, the important factors that control the flux creep can be understood.

In general, the HTSC’s are highly anisotropic (see, for example, Ref. [16]), so it
Table 3.2: Samples used in this study

<table>
<thead>
<tr>
<th>SAMPLE</th>
<th>$T_c$ (K)</th>
<th>TYPE</th>
</tr>
</thead>
<tbody>
<tr>
<td>Tl$_2$Ba$_2$Ca$_2$Cu$<em>3$O$</em>{10}$</td>
<td>120</td>
<td>grain-aligned epoxy composite, HTSC</td>
</tr>
<tr>
<td>Nb-Ti</td>
<td>9.4</td>
<td>multi-filament wire, LTSC</td>
</tr>
<tr>
<td>Nb$_3$Sn</td>
<td>17.75</td>
<td>superconducting tape, LTSC</td>
</tr>
<tr>
<td>(Ba$_{1-x}$K$_x$)BiO$_3$</td>
<td>28</td>
<td>epoxy composite (not aligned), HTSC</td>
</tr>
<tr>
<td>(Nd$_{1-x}$Ce$_x$)$_2$CuO$_4$</td>
<td>21</td>
<td>grain-aligned epoxy composite, HTSC</td>
</tr>
<tr>
<td>Bi$_2$Sr$_2$CaCu$_2$O$_x$</td>
<td>90</td>
<td>single crystal, HTSC</td>
</tr>
</tbody>
</table>

Figure 3.2: Randomly oriented grains before the application of magnetic field. The easy axis (see text) is indicated by the arrow.

Figure 3.3: Grains aligned with the east axis parallel to the field $H$. 
is necessary to obtain samples that possess a high degree of preferred orientation, and to eliminate any grain-to-grain coupling, as grain boundaries act as weak links [61]. Both of these conditions can be met by the epoxy-alignment method [41], which takes advantage of the anisotropy of the normal state magnetic susceptibility. The sample is initially a sintered pellet which is then powdered into individual single crystals, mixed with epoxy, and the conglomerate is allowed to harden in a high magnetic field, (i.e., 8 T). The mixture is usually rather dilute (< 50% superconductor) in order to prevent bad alignment caused by individual particles interfering with one another. When first mixed with the epoxy, the powder is randomly oriented (Fig. 3.2), but the field application causes a net torque on each particle, causing each particle to rotate such that the crystal axis with the largest moment (the "easy" axis, which is usually the c-axis) aligns with the field (Fig. 3.3). The sample is then kept in the field until the epoxy hardens to prevent the particles from relaxing. Once hard, the particles are fixed in place, and the magnetic properties of the sample for each direction can be determined. This method allows the experimenter to take measurements on what amounts to millions of single crystals (at the same time), without some of the inherent problems associated with single crystals, such as large demagnetizing factors [5], internal weak-links [91], and incomplete oxygenation [92].

It is necessary that each particle be a single crystal grain, because multi-grain particles will not have a well defined easy axis. If the powder is composed of such particles, it will not align well, and the measurements performed on the sample will be affected by the anisotropy. A hypothetical bi-crystalline grain is illustrated in Fig. 3.1, where the two different c-axis orientations are shown by the arrows. This grain will probably line up along direction defined by the mean of the two c-axis orientations.
resulting in no well-defined preferred orientation. Therefore, it is necessary to grind the powder carefully and screen out any large particles with an appropriately sized mesh (i.e., a 38 μm mesh should get rid of most of the multigrain particles).

The degree of alignment can be checked by performing an x-ray rocking curve on an appropriate diffraction peak. First the aligned sample is scanned normally in an x-ray diffractometer, with the x-ray beam incident along one of the defined axes, i.e., the c-axis. (This is the usual choice, so from here on, c-axis orientation will be assumed for brevity.) According to the rules of x-ray diffraction [5], only the (00l) peaks should be observed. If any other peaks are observed, the sample is poorly aligned and the alignment process must be done over. Assuming that only (00l) peaks are seen, one of these is selected (i.e., for the Tl2Ba2Ca2Cu3O10 sample studied here, the (005) peak was chosen). The sample is then re-scanned about this peak, with the sample held fixed in place. The resulting diffraction pattern looks like a resonance peak, and the quality of the alignment can be determined by the sharpness of the peak. The full-width half-maximum (FWHM) is defined as the width of the peak at one half of the peak height. A sample has good alignment if the FWHM is small. Although x-ray diffraction is most sensitive to surfaces [91], this
method gives an indication of the quality of alignment.

To examine the interior of the sample, it is necessary to cut it open and mount it in a binder for metallographic examination under a microscope. This method has the advantage of exposing the entire sample, but the distinct disadvantage of rendering the sample useless for measurement afterwards. Therefore, this should be done only after all of the desired data has been taken on the sample, at which time it is presumed that the sample would be expendable. For the case of YBa$_2$Cu$_3$O$_{7-\delta}$, optical studies of a cross-sectional cut confirm that the alignment is good throughout the sample [11].

One material that was not studied here was YBa$_2$Cu$_3$O$_{7-\delta}$. This was because YBa$_2$Cu$_3$O$_{7-\delta}$ single crystals often suffer from poor sample quality, especially in the area of incomplete oxygenation [92]. One sensitive measure of the degree of oxygenation is the “fishtail” effect which occurs in magnetization curves. Figure 3.5 (taken from Ref. [92]) shows such a curve taken at 70 K. A large maximum occurs at 4 T, which translates to a maximum in the magnetization $J_C$ at the same field. It is also shown that as the oxygen content is increased, the fishtail goes away. In the samples studied here, no fishtails were observed, indicating good oxygenation. YBa$_2$Cu$_3$O$_{7-\delta}$ samples also have twinning planes [11], and twinned samples cannot be considered as single crystalline because there is no well defined $a$ or $b$ axis direction.

**Tl$_2$Ba$_2$Ca$_2$Cu$_3$O$_{10}$**

This class of superconductor was first discovered by Sheng and Hermann at the University of Arkansas [3], with onset temperatures as high as 120 K and zero re-
Figure 3.5: A “fishtail” magnetization curve for an YBa$_2$Cu$_3$O$_{7-\delta}$ single crystal. The predominant superconducting phases were first identified by Hazen et al. [95] as being Tl$_2$Ba$_2$Ca$_2$Cu$_3$O$_{10}$ and Tl$_2$Ca$_1$Ba$_2$Cu$_2$O$_{10}$. The highest $T_c$ phase (Tl$_2$Ba$_2$Ca$_2$Cu$_3$O$_{10}$) was found by Parkin et al. [96] at IBM Almaden to have $T_c$ as high as 125 K. The IBM group also found several other Tl-based phases [97], with $T_c$'s ranging from 65 to 125 K. These materials were found to have a layered perovskite structure, with a tetragonal rather than an orthorhombic crystal structure, so these materials do not form twin planes. Single-crystalline samples of the Tl-Ca-Ba-Cu-O system are therefore much more likely to be single-crystalline than in the case of YBa$_2$Cu$_3$O$_{7-\delta}$.

Study of the Tl-Ca-Ba-Cu-O systems is also of interest because $T_c$ is approximately 30 K higher than that of YBa$_2$Cu$_3$O$_{7-\delta}$, so the reduced temperature $t = T/T_c$ is 27% lower at 77 K ($t = 0.62$ for Tl$_2$Ba$_2$Ca$_2$Cu$_3$O$_{10}$ with $T_c = 125$ K, 0.84 for YBa$_2$Cu$_3$O$_{7-\delta}$ with $T_c = 92$ K). This would appear to make Tl$_2$Ba$_2$Ca$_2$Cu$_3$O$_{10}$
(or Bi-Sr-Ca-Cu-O, which has similar $T_c$'s) more likely to be useful in liquid nitrogen. A quick scan of the literature, however, shows that the reversible part of the $H-T$ plane is much larger for Bi-Sr-Ca-Cu-O [6] and Tl$_2$Ba$_2$Ca$_2$Cu$_3$O$_{10}$ [7] than for YBa$_2$Cu$_3$O$_{7-\delta}$ [5]. Therefore, studying the flux pinning and flux creep properties of Tl$_2$Ba$_2$Ca$_2$Cu$_3$O$_{10}$ is necessary to understand the causes and origins of such a mobile flux lattice, and to learn how to improve the flux-pinning qualities of these materials.

The Tl$_2$Ba$_2$Ca$_2$Cu$_3$O$_{10}$ sample used in this study was supplied by D. F. Farrell of Case Western Reserve University and N. P. Bansal of NASA Lewis Research Center [98]. The sample was given the name of 141-8b, and will referred to here as Tl-1. The sample was a grain-aligned epoxy composite. During the alignment process, a provision for easy sample mounting and crystal axis definition was made by the inclusion of two holes through the entire sample. The sample could be oriented with the field parallel to either the c-axis or the basal plane by mounting the sample on a quartz rod using the appropriate hole.

Before the sample was aligned in epoxy, the powder was pressed into a ceramic pellet and the resistive transition was measured, with $T_c (R=0) = 119$ K [98] (sample Tl-1 is labeled as T-9 in Tab. 2 of Ref. [98]) (this work was done before the sample was sent to Ames Lab). Fang et al. [7] measured $T_c$ magnetically at Ames Lab to be 120 K, which agrees with this work. The magnetic transition for an applied field of 100 Oe is shown in Fig. 3.6. According to Parkin et al. [96], $T_c$ for Tl$_2$Ba$_2$Ca$_2$Cu$_3$O$_{10}$ can vary from 118 to 125 K, with $T_c$ dependent on the sample preparation conditions, especially the composition of the starting material. Bansal and Farrell found that isolating the exact "2223" phase is extremely difficult and that the starting composition
Figure 3.6: Transition of Tl$_2$Ba$_2$Ca$_2$Cu$_3$O$_{10}$ sample Tl-1 in an applied field of 0.01 T with H $\parallel$ c must be off-stoichiometric.

All IITSC's have an irreversibility transition, defined as the temperature for a given field at which the difference between the ZFC and FC signals is zero. Two methods are usually employed to determine the irreversibility point, a criterion \((M_{FC} - M_{ZFC}) = 0.01M_{ZFC}\), or by curve fitting the difference of the ZFC and FC signals. The temperature at which the curve extrapolates to zero is the irreversibility temperature. The 1% criterion method seemed rather arbitrary, and was therefore not used.

It is often difficult to determine the irreversibility point accurately. In a low field of 0.01 T, the signal difference goes to zero rapidly at 110 K, so that the uncertainty in \(T_{irr}\) is only in the temperature controller, 10 mK. In higher fields, the signal
Figure 3.7: Determination of the irreversibility point for 0.5 T ($H \parallel c$)

difference goes to zero much more gradually, so it was necessary to fit a fourth-order polynomial to the signal difference and define the irreversibility point as the intersection of the fit with the temperature axis. An example of this is shown in Fig. 3.7, where the field is 0.5 T. The irreversibility point here is 79.7 ± 2 K, and the error bars in the data are reflected by the plot symbol size. The entire $H_{irr}(T)$ curve is shown in Fig. 3.8, along with $H_{irr}$ determined by a 1% criterion as a comparison. It is easily seen that the criterion method results in a lower $T_{irr}$ than the curve-fitting method for a given field.

In contrast to other reports (Refs. [4] and [5], for example), no well defined temperature dependence for $H_{irr}(T)$ was found. The low field region ($0 \leq \mu_0 H \leq 0.05 T$) can be described with a linear fit, whereas at higher field the temperature dependence is approximately cubic, i.e., $H_{irr} \propto (1 - t)^3$, where $t = T/T_c$ is the
Figure 3.8: Irreversibility transition for Tl-1 with H || c

reduced temperature. This cubic dependence has also been observed by Green et al. [6] on a sample of Pb-doped Bi$_2$Sr$_2$CaCu$_2$O$_{x-2}$ which has the same structure as Tl$_2$Ba$_2$Ca$_2$Cu$_3$O$_{10}$.

**Nb-Ti**

**General** The Nb-Ti sample used in this study (to be referred to as Nb-Ti-1) was cut from a piece of commercially available wire found in the lab. The wire was from the Magnetic Corporation of America (MCA), with the production data shown in Table 3.3. Nb is a cubic material (bcc), and because Nb-Ti is a solid-solution alloy of Nb, Nb-Ti should also be cubic. The composition of the sample is nominally 46.5 wt. % Ti and 53.5 wt. % Nb with approximately 1.5 to 2 % variation in these percentages. This sample gives a good representation of a state-of-the-art
Table 3.3: Production data for the Nb-Ti sample Nb-Ti-1 (from the MCA data sheet)

<table>
<thead>
<tr>
<th>Billet #</th>
<th>141</th>
</tr>
</thead>
<tbody>
<tr>
<td>size</td>
<td>0.030 in x 0.058 in</td>
</tr>
<tr>
<td>Cu/SC ratio</td>
<td>1.25 : 1</td>
</tr>
<tr>
<td># filaments</td>
<td>367</td>
</tr>
</tbody>
</table>

Table 3.3: Production data for the Nb-Ti sample Nb-Ti-1 (from the MCA data sheet)

multifilamentary superconducting wire.

The micrographs of Nb-Ti-1 in Fig. 3.9 show multifilamentary structure of the sample. It can be seen that the filaments are arranged in an approximately hexagonal array, and are roughly the same size. It is necessary to know the volume of the sample in order to calculate the magnetization in G (4πM) from the measured moment by the magnetometer. This can be computed by (at least) two different methods. In Fig. 3.9(b) each filament is roughly circular in cross-section with a diameter of 0.25 in. The true size is $d = 0.25 \text{ in} \times 2.54 \text{ cm/in} / 165 = 38.5 \mu\text{m}$, giving a filament radius $\rho = 19.2 \mu\text{m}$. The sample as measured consisted of three lengths of wire that had a total length of $l = 1.2268 \text{ cm}$. The superconducting volume is the number of filaments times the area of each filament times the length of the sample, that is, $V_{sc} = 367 \times \pi \rho^2 \times l = 0.0052 \text{ cm}^3$. The superconducting volume can also be calculated from the MCA data sheet. The total volume of Nb-Ti-1 is the area $(0.030 \times 0.058 \text{ in}^2)$ times the length l, which equals 0.0138 cm$^3$. The Cu/SC (copper to superconductor) ratio is a volume ratio. This means that $V_{sc} = V_{tot} / 2.25 = 0.0061 \text{ cm}^3$, which, by inverting the previous method, gives a radius $\rho = 20.8 \mu\text{m}$. These methods agree to 8.3%, with somewhat stronger confidence towards the first method, since the second requires the assumption that the data sheet is always correct. For all calculations, a radius of 20 $\mu\text{m}$ will be employed, which corresponds to a superconducting volume of
Figure 3.9: Two micrographs of sample Nb-Ti-1 with magnifications (a) 41.25x and (b) 165x
Superconducting Properties  The superconducting transition of Nb-Ti-1 is shown in Fig. 3.10. The field was 50 Oe and the data were taken after cooling to 4.5 K in zero field. The transition temperature $T_c = 9.37 \pm 0.13$ K, which is the midpoint between the last superconducting and the first normal data points. The transition is 0.61 K wide as determined from the difference in the temperatures at which the signal is 10% and 90% of the value at 4.5 K. This is a good quality sample with no second phase signals apparent. In all forthcoming discussions, $T_c = 9.4$ K.

One major difference between this material and Tl$_2$Ba$_2$Ca$_2$Cu$_3$O$_{10}$ was that here, the critical fields were small enough that they could be directly measured over
much of the $H$-$T$ plane. These are very useful in determining $\kappa$, $\lambda$, and $\xi$, as well as the $H$-$T$ phase diagram for the sample. Also, $H_c$ could in principle be determined by assuming BCS behavior, giving the free energy/unit volume $\Delta G = H_c^2/8\pi$ and $\Delta G/l$, the free energy/unit length of a vortex, but this was not done here.

The upper critical field $H_{c2}$ was determined by measuring the transition temperature in different applied fields. Because the sample is irreversible throughout the entire $H$-$T$ plane (as will be shown in the next chapter), only the field-cooled curve was measured. Quantum Design magnetometers work best for increasing-temperature measurements, so the data were taken by first cooling in field $H$ to $T = 5$ K, and then warming by at 0.25 K intervals to 10 K, after which the temperature interval was 0.5 K. The value of $T_{c2}$, defined as $T_c(H_{c2})$ was determined by first fitting the normal state data to a straight-line and subtracting the calculated normal state from
the data. The correction can be as large as 30 % in 5.5 T, and is linear with field. A polynomial fit to the superconducting data was applied, and the intersection with the fit and the temperature axis is \( T_{c2} \). An example of this is shown in Fig. 3.11.

The resulting \( H_{c2}(T) \) is shown in Fig. 3.12. The temperature dependence was found to be parabolic, with \( H_{c2} = 14.6(1 - t^2) \). This fit was found to predict the data points to an accuracy of 4.2% for fields greater than 0.7 T (the fit is much poorer close to \( T_c \) in fields of less than 0.5 T). The Gor'kov fit [99] could also be applied (this contains the parabolic dependence as well as a fourth-order polynomial in temperature), but the simpler parabolic fit was the method used. Both methods fit the data well (at \( t = 0.75 \), the differences of the parabolic fit and the Gor'kov fit are much less than 1%), but the Gor'kov fit results in \( H_{c2}(0) = 18 \) T, which seems too large for Nb-Ti. Using \( H_{c2} \), the coherence length \( \xi \) can also be calculated from
Eq. 1.14.

The lower critical field $H_{c1}(T)$ was determined by measuring the initial magnetization curves at fixed temperatures. When the applied field is less than $H_{c1}$, the magnetization is linear in $H$. The measured signal deviates from linearity at $H_{c1}$, so by fitting a straight line to the low field data, $H_{c1}$ is defined to be the field where the data first deviates by more than 1% from the fit. One thing to note was the slope $dM/dH$ at low fields was constant with temperatures, as expected for a superconductor in the Meissner state at low fields. This slope was found to be -0.00011 emu/Oe. Dividing by the volume of 0.0057 cm$^3$, and multiplying by $4\pi$, the initial slope = -0.97. This justifies ignoring any demagnetization effects, so the only correction needed for the low field data is the small, negative field trapped in the solenoid, even after demagnetization. This was done by reading the x-intercept off of an initial magnetization curve and adding it to the field at which the deviation from linearity occurs. This method of determining $H_{c1}$ is illustrated in Fig. 3.13.

Nb$_3$Sn

An attempt was made to measure flux creep in Nb$_3$Sn tape. However, the creep characteristics of the tape were quite complicated, with flux jumps superimposed on the logarithmic decay, for example. This made analysis extremely difficult, so the results presented here will be quite sketchy and preliminary.

The tape was cut from commercially available GE stock. The tape had a sandwich structure, with outer layers of Cu, and two layers of Nb$_3$Sn deposited on a non-metallic substrate. A micrograph of the sample is shown in Fig. 3.14.

In low fields, there appeared to be three superconducting transitions, at 17.76
K, 9 K, and 7 K (Fig. 3.15). The main transition at 17.76 K is $T_c$ of the Nb$_3$Sn. The 7 K transition is probably due to some Pb solder left on the tape, and the 9 K transition is some unreacted Nb. In a 1 T field, the lower $T_c$ phases disappear, and the transition is suppressed to 17.26 K, which gives a (very) rough estimate of $dH_{c2}/dT \approx 2$ T/K. No estimate of the sample volume could be made, due to the complicated low field nature of the transition.

When the observation of flux creep was attempted, flux jumps often occurred. This made the analysis extremely difficult, as the jumps were not always sudden. For this reason, the study of this sample was abandoned in favor of the Nb-Ti sample, which did not have flux jump problems. The logarithmic creep rates do appear to be nearly the same before and after the jump, as is shown in Fig. 3.16. This was not always the case, however.
Figure 3.14: Micrograph of sample Nb$_3$Sn-1, magnified 165 times. The outer layers are Cu, and the narrow dark bands near the center are the Nb$_3$Sn layers.
Figure 3.15: Superconducting transition of Nb$_3$Sn-1 in a field of 0.01 T

Figure 3.16: A flux jump during a creep measurement on Nb$_3$Sn-1. $T = 10$ K. $H = 15$ kOe
The $(\text{Nd}_{1-x}\text{Ce}_x)_2\text{CuO}_4$ sample used for comparison with $\text{Tl}_2\text{Ba}_2\text{Ca}_2\text{Cu}_3\text{O}_{10}$ was reported in more detail by Sanders et al. [5]. The transition temperature was 21 K as determined by magnetization measurements. The sample itself was a grain-aligned epoxy composite, prepared similarly to the $\text{Tl}_2\text{Ba}_2\text{Ca}_2\text{Cu}_3\text{O}_{10}$ sample. Since the presence of Nd in the sample gives rise to a large paramagnetic signal, and the normal state magnetization did not fit a simple Curie-Weiss law, the superconducting properties were deduced by subtracting the paramagnetic background. It was assumed that the paramagnetic moment of the Nd and Ce ions did not change with O$_2$ content, so that the paramagnetic component was the same in the superconducting and normal states. This was confirmed experimentally by preparing a non-superconducting sample with lower O$_2$ content and comparing the data above $T_c$, so it was also assumed to be true below $T_c$. The superconducting magnetization was then determined by subtraction of the normal state.

The irreversibility line was determined using a difference criterion of 0.1 memu, which differs from the curve fitting method for $\text{Tl}_2\text{Ba}_2\text{Ca}_2\text{Cu}_3\text{O}_{10}$. Log-log fitting revealed a power law, with the exponent being 0.28, that is, $(1 - T_{irr}/T_c) \propto H^{0.28}$. The exponent is close to 1/3, as found by Green et al. [6] for Bi-Sr-Ca-Cu-O single crystals. The upper critical field $H_{c2}$ was determined by fitting the data above $T_{irr}$ and below $T_c$ to a line and extrapolating to $M = 0$. After considerable curvature near $T_c$, $dH_{c2}/dT$ was found to be $-0.9$ T/K for $H \parallel c$ and $-1$ T/K for $H \perp c$. 

(Ba$_{1-x}$K$_x$)BiO$_3$

The cubic HTSC used for comparison with the Tl$_2$Ba$_2$Ca$_2$Cu$_3$O$_{10}$ data was a (Ba$_{1-x}$K$_x$)BiO$_3$ sample having a $T_C$ of the sample was 28 K as determined by magnetization measurements. In order to prevent grain-to-grain coupling, the sample was mixed with epoxy in the same manner as the (Nd$_{1-x}$Ce$_x$)$_2$CuO$_4$ and Tl$_2$Ba$_2$Ca$_2$Cu$_3$O$_{10}$ samples. No preferred orientation was obtained, of course, because this material is cubic.

The irreversibility line was defined by using a criterion of $|M_{ZFC} - M_{FC}| \leq 0.1$ mm. A log-log fit of $(1 - T_{irr}/T_C)$ vs. $H$ was gave a power law $(1 - T_{irr}/T_C) \propto H^{0.61}$, the exponent being quite close to the value of 2/3 found by Yeshurun and Malozemoff [5]. The upper critical field $H_{c2}$ was determined by linearly extrapolating the magnetization data below $T_C$ to $M = 0$. This resulted in a nearly linear (with a slight positive curvature) $H_{c2}$ with $dH_{c2}/dT = 0.42$ K. Fitting a different region nearer to $T_C$ resulted in $dH_{c2}/dT = 0.74$ K. More detailed analysis of this sample will be published by Sanders in his PhD. thesis. The details of the (Ba$_{1-x}$K$_x$)BiO$_3$ sample studied here have not yet been published.

Bi$_2$Sr$_2$CaCu$_2$O$_x$

This Bi$_2$Sr$_2$CaCu$_2$O$_x$ sample used for comparison with Tl$_2$Ba$_2$Ca$_2$Cu$_3$O$_{10}$ was reported by Shi et al. [100]. The sample was a single crystal prepared by the flux growth method [101] and annealed at 800 °C for seven days in air on a Au substrate. The crystal was cooled slowly to room temperature after annealing. The dimensions of the crystal were $3 \times 0.5 \times 0.1 \times \text{mm}^3$. The transition temperature was found to be 86 K by measuring the magnetization of the sample in a S.H.E. SQUID magnetometer.
In measuring the resistivity, a current density of 50 A/cm² was used. The field could be oriented either parallel or perpendicular to the $a-b$ plane of the crystal.

In Ref. [100], the resistive transitions in magnetic fields were analyzed using a thermally activated resistivity model. The resistivity was plotted as $\ln \rho$ vs. $1/T$, which resulted in a linear fit. $U_0$ was found to obey the temperature and field dependence of

$$U_0 = A(H)(1 - T/T_c)^m.$$  \hspace{1cm} (3.1)

The parameters found for this sample with both field orientations are given in Table 3.4, and the dependence of $A(H)$ on field is shown in Fig. 3.17.
Figure 3.17: Field dependence of the parameter $A(H)$ obtained from the Arrhenius fitting
CHAPTER 4. SINGLE BARRIER RESULTS

\textbf{Tl}_2\textbf{Ba}_2\textbf{Ca}_2\textbf{Cu}_3\textbf{O}_{10}

\textbf{Overview of Results}

This section will be organized as follows. First, standard hysteresis measurements are presented. From these data, the critical current \( J_c \) is determined as a function of field and temperature using the Bean model. The flux pinning force per unit volume \( F_p = |J_c \times B| \) is studied as a function of \( H \) and \( T \). Then, general results of flux creep measurements will be presented. These will entail the general features of creep measurements, such as time dependence, magnitude of creep (i.e., 5\% change in one hour), and the logarithmic creep rate \( S = dM/d\ln t \) is also shown as a function of field and temperature. These data are then combined to calculate the effective pinning potential \( U_{\text{eff}}(H,T) \) in the range where the BLW model applies. In the analysis it is seen that \( U_{\text{eff}}/kT \) is the most relevant parameter for creep, and \( U_{\text{eff}}/kT \) is determined at many fields and temperatures. A method for plotting lines of constant \( U_{\text{eff}}/kT \) in the \( H-T \) plane is presented, which shows that there exists a band in the \( H-T \) plane where \( 2 \leq U_{\text{eff}}/kT \leq 20 \), and that \( U_{\text{eff}}/kT \leq 2 \) over a large portion of the \( H-T \) plane.
<table>
<thead>
<tr>
<th>$B$ (T)</th>
<th>$\Delta B$ (T)</th>
</tr>
</thead>
<tbody>
<tr>
<td>0</td>
<td>0.02</td>
</tr>
<tr>
<td>0.1</td>
<td>0.05</td>
</tr>
<tr>
<td>0.3</td>
<td>0.10</td>
</tr>
<tr>
<td>0.4</td>
<td>0.2</td>
</tr>
<tr>
<td>1.0</td>
<td>0.5</td>
</tr>
<tr>
<td>1.5</td>
<td>0.4</td>
</tr>
</tbody>
</table>

Table 4.1: Field cycle used in hysteresis measurements on sample T1-1. The field was stepped up with increment $\Delta B$ until the next $B$ value was reached. For the decreasing-field cycle, read up.

**Hysteresis Results**

The behavior of the critical current density $J_c$ throughout the $H$-$T$ plane characterizes the flux pinning properties of the sample. Hysteresis measurements were performed to find the field and temperature dependence of $J_c$, which in turn is used for calculating the volume flux pinning force and $U_{eff}$ from the BLW model.

The hysteresis measurements were performed at several temperatures in the SQUID magnetometer QD1. The measurement sequence normally used was to first apply a $-2$ T field to fully soak the sample with flux, reduce the field to zero, and then cycle the field from zero to 2 T back to zero. To reduce field overshoot to values less than 0.4% of the set point field, the field was increased in small increments noted in Table 4.1. The initial cycling of the field put the sample into the critical state, so only a half-loop was required.

After the field was set, five scans were recorded at each field to allow instrument to set the proper SQUID range for transients to die out. The critical current density $J_c$ was determined via use of the Bean model (Eq. 2.10). Usually, the fourth scan at each field was used, as this was often the first good measurement (as determined by...
Figure 4.1: \( J_c \) vs. \( T \) for \( \mu_0 H = 0.1, 0.25, 0.8, \) and \( 1.0 \) T of Ti-1 (\( H \parallel c \))

examining the signal printout from the computer) at each field. This corresponded to a time of approximately 10 minutes after each field was set, in which \( J_c \) has relaxed typically by 5%, consequently, any quantity that uses \( J_c \) (\( U_{eff} \) and \( F_p \), for example) will be underestimated by at least this amount. One possible method of recovering the initial value of \( J_c \) would be to extrapolate the logarithmic time dependence of these measurements back to some convenient reference time such as 1 minute, and use these values of \( M \) to calculate \( J_c \). However, because there were often only two scans that were not clipped, this was not deemed practical, and the underestimated values were used as is.

In order to show the temperature and field dependence of \( J_c \), \( J_c \) is plotted vs. \( T \) in Fig. 4.1, and vs. \( H \) for in Fig. 4.2, with \( H \parallel c \). The major features of \( J_c \) are that \( J_c \) is approximately exponential in \( T \), \( \left( J_c \sim \exp\left(-T/T_0\right) \right) \), and that \( J_c \) is approximately
Figure 4.2: $J_c$ vs. $H$ for $T = 10, 15, 20, 35, 50$, and $78$ K of Tl-1 ($II \parallel c$)

exponential in field at low temperatures ($T \leq 35$ K), but drops off faster in higher temperatures.

A maximum in $J_c$ occurs at non-zero field, which decreases with increasing temperature. Using $dH/dx = (4\pi/10)J_c$, and $H = 0$ at the center of the sample when $H = H^*$, it can be seen that when $H/(dH/dx) = R$, $H = H^*$. In Table 4.2, values of $H/(dH/dx)$ at the field of maximum $J_c$ are shown, and since the grain

<table>
<thead>
<tr>
<th>$T$ (K)</th>
<th>$J_{c,\text{max}}$ (A/cm$^2$)</th>
<th>$\mu_0 H^*$ (T)</th>
<th>$H/(dH/dx)$ ($\mu$m)</th>
</tr>
</thead>
<tbody>
<tr>
<td>10</td>
<td>$2.000 \times 10^6$</td>
<td>0.225</td>
<td>9.0</td>
</tr>
<tr>
<td>20</td>
<td>$8.991 \times 10^5$</td>
<td>0.15</td>
<td>13.4</td>
</tr>
<tr>
<td>35</td>
<td>$3.537 \times 10^5$</td>
<td>0.06</td>
<td>13.5</td>
</tr>
<tr>
<td>50</td>
<td>$1.882 \times 10^5$</td>
<td>0.03</td>
<td>12.7</td>
</tr>
</tbody>
</table>

Table 4.2: Data showing that the field of maximum $J_c$ is the full penetration field $H^*$. The sample grain size is 10-15 $\mu$m.
radius is 10-15 \( \mu m \), it can be concluded that \( J_c \) is maximizes at \( H = H^* \).

Two reasons for the maximum in \( J_c \) are that either in the low field region \( H < H^* \), Bean's model does not apply, or that there is a pinning mechanism that increases with field such as the "fish-tail" effect [92]. The first reason is probably more likely, because Beasley et al. [70] observed the same effect in Pb-Tl alloys, where the pinning mechanisms are impurities and defects, and not oxygen content.

These measurements were performed for two field orientations, \( H \parallel c \) and \( H \perp c \), in order to see the effect of anisotropy on \( J_c \). The emphasis in this work, however, is for \( H \parallel c \) for two reasons. First, if \( H \parallel c \), then only currents in the Cu-O sheets are induced, so that the measured \( J_c \) is \( J_c^{aa,c} \), where \( J_c^{aa,c} \) refers to the critical current density in the basal planes when the field is in the \( c \) direction. When \( H \perp c \), the induced currents will include a contribution \( J_c^{aa,aa} \) from the Cu-O sheets and \( J_c^{c,aa} \) along the \( c \)-axis, making the data harder to interpret. The other reason is that a slight misalignment in the grains causes \( B \) to align with the \( c \) axis. This is because these grains are usually very thin along the \( c \) axis, which causes them to be much more sensitive to misalignment, and the anisotropy of the effective mass tensor [102]. Figure 4.3 shows the critical current density as a function of angle, clearly showing the effect of misalignment [103]. This effect has been observed in other works as well [104]. For these reasons, most of the work was done for the case of \( H \parallel c \).

To calculate \( J_c^{c,aa} \) from the \( H \perp c \) data, the approach of Sauerzopf et al. [105] gives \( J_c^{c,aa} \) by

\[
J_c^{c,aa} = \frac{3cJ_c^{aa,aa}}{2a} \left[ 1 - \left( 1 - \frac{16\Delta M}{3cJ_c^{aa,aa}} \right)^{1/2} \right]. \tag{1.1}
\]

where \( c \) is the dimension along the \( c \)-axis and \( a \) is the dimension along the \( a \) axis (basal plane). If it is assumed that \( J_c^{aa,aa} = J_c^{aa,c} \), then the \( H \parallel c \) data can be used
Figure 4.3: Angular dependence of $J_c$ in a 7 T field for 1.2, 10, and 60 K (from Roas et al. Phys. Rev. Lett. 64, 479 (1990))
in Eq. 4.1. While this assumption is probably not valid [106], the only other method is to use several samples of different sizes or to decrease the sample size, which for Tl-1 can’t be done. Therefore, it will be assumed that $J_c^{aa,aa} = J_c^{aa,cc}$.

$J_c^{c,aa}$ and $J_c^{aa,cc}$ are shown as functions of temperature and field in Fig. 4.4. It is easily seen that $J_c^{aa,cc}$ is always greater than $J_c^{c,aa}$ for the field ($0.1 \leq H \leq 1.0$ T) and temperature ($10 < T < -50$ K) ranges studied. The current density anisotropy does decrease with increasing field. Further study of $H \perp c$ data was not pursued, however, because of the assumptions involved in extracting $J_c^{c,aa}$ from the data.

The field dependence of $J_c$ can be seen to be at least approximately exponential in field, $J_c \sim \exp(-H/H_0)$ for $T \leq 35$ K. This is a common feature often seen in HTSC’s, both in ceramic samples [107], where it is attributed to the granularity of the samples, and in aligned samples [40], where it is found to be an intrinsic property of the sample. Because this sample is composed of single crystals, this means two things. Either the individual particles are not single grains, or this exponential behavior is intrinsic to HTSC’s. Single crystals also display exponential field dependence [108], so it would appear that it is intrinsic, and that there exists within HTSC’s some internal weak-link mechanism. The decrease of $J_c$ with field at higher temperatures $T > 50$ K is due to the highly reversible nature of the sample in this region of the $H-T$ plane.

At low temperatures ($10$ K), $J_c^{aa,cc} \sim 10^6$ A/cm² for $B < 1$ T. Anisotropy remains as a concern for practical applications of Tl$\text{Ba}_2\text{Ca}_2\text{Cu}_3\text{O}_{10}$, however, because it would require some sort of preferred orientation in the sample, or the useful critical current would be limited by that of the lower orientation. The temperature dependence of $J_c$ shows that this material would only be practical for high-current
Figure 4.4: $J_c$ of Tl-1 vs. (a) $T$ for $\mu_0 H = 0.1$ and 1.0 T; (b) $H$ for $T = 15$ and 50 K with both field orientations, with $J_c^{c,a}$ extracted from the data as described in the text.
applications (such as superconducting magnets) at low temperatures, and the field dependence, even at low temperatures, makes this sample a poor candidate to replace conventional superconductors, such as Nb-Ti. Also, the anisotropy cross-over indicates that the flux pinning in this sample is very complicated.

In order to examine the flux pinning characteristics of this sample, it is useful plot the volume pinning force \( F_p = |J_c \times B| \) vs. \( B \), as shown in Fig. 4.5. For conventional superconductors, the resulting curve shows a peak in \( F_p \) at either \( B_{c2}/5 \) (Nb$_3$Sn) or \( B_{c2}/2 \) (Nb-Ti) when in fields near \( B_{c2} \). In Fig. 4.5 peaks in \( F_p \) are observed, however, these peaks occur at fields that are far below \( B_{c2} \). Also, the shape of these curves do not look anything like the expected field dependencies \( b^{1/2}(1-b)^2 \) or \( b(1-b) \), where \( b \) is the reduced field \( B/B_{c2} \). This result is not surprising, because \( B_{c2} \) is extremely large in this sample, and the reversible region is very wide, so it is impossible to have flux pinning near \( B_{c2} \). The irreversible transition is the cut-off point of \( F_p \) for this sample [109], and the field at which \( F_p \) peaks decreases as the temperature is increased, indicating that the pinning in these materials is being modulated by \( H_{irr}(T) \). Although the field dependence of \( F_p \) could not be determined, a plot of the peak pinning force vs. the field at which the peak occurs reveals a power law scaling, \( F_{p,max} \propto B_{max}^2 \) (Fig. 4.6), as expected in conventional superconductors [110].

**Flux Creep Results**

**General** To study the time dependence of the magnetization, flux creep was observed in both ZFC and hysteresis measurements over a wide range of temperatures and fields. The following discussion will center on the results of the ZFC measurements, because these data were recorded over a period of 1 hour rather than
Figure 4.5: Volume pinning force $F_p$ of Ti-1 vs. $H$ ($H \parallel c$) for (a) 10, 15, and 20 K; (b) 35 and 50 K; (c) 78 K. Note the change of scale as the temperature increases.
Figure 4.6: Scaling of $F_{p,max}$ vs. $B_{max}$. The straight line is a log-log fit, with slope $n = 2$.
Table 4.3: The region of the $H-T$ plane in which flux creep is logarithmic in time throughout the entire measurement.

<table>
<thead>
<tr>
<th>$T$ (K)</th>
<th>$\mu_0 H^*$ (T)</th>
<th>$\mu_0 H_{max}$ (T)</th>
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<tr>
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<td>0.25</td>
<td>&gt; 2</td>
</tr>
<tr>
<td>20</td>
<td>0.15</td>
<td>&gt; 2</td>
</tr>
<tr>
<td>35</td>
<td>0.08</td>
<td>0.8</td>
</tr>
<tr>
<td>50</td>
<td>0.06</td>
<td>0.4</td>
</tr>
<tr>
<td>78</td>
<td>0.02</td>
<td>0.08</td>
</tr>
</tbody>
</table>

Although creep was observed throughout the $H-T$ plane, the range in which the creep is linear in $\ln t$ was somewhat limited. In low temperatures $T \leq 20$ K and all fields $H^* \leq H \leq 2T$, the time dependence of the creep was logarithmic, but at $T \geq 35$ K, there existed a field above which the creep becomes non-logarithmic at longer times. It is useful to discuss each of these effects separately. The logarithmic flux creep "envelope" is shown in Table 4.3, where the field $H_{max}$ is the largest field at that temperature where the creep is still logarithmic.

An example of logarithmic flux creep is shown in Fig. 4.7, which is a ZFC creep
measurement at $T = 20$ K and in an external field of $\mu_0 H = 0.15$ T. In this measurement, the first two scans had autoranging errors, giving an initial "dead" time of 5.3 minutes. By extrapolating back to at $t = 1$ min. (this is the time to which all measurements were referenced) it was found that the initial magnetization was $-286.72$ emu ($-0.0282$ G). If extrapolated out to a total time of one hour, the magnetization has decayed to 86% of the initial value ($-216.25$ emu). The creep rate $S$ was $9.883$ emu ($20.78$ G), and the change in magnetization in one decade of time, $\frac{dM}{d\ln t} \times \ln 10 = 2.303 \times S$, was $17.8$ G per decade. Now, $(1/2)\Delta M = 299.6$ G at this temperature and field, so it would take approximately 6.3 decades of time ($10^6.3$ minutes, which is approximately 3.5 years) for the current to decay away and the sample to reach the reversible equilibrium state. This is a long time, so even though the creep rate is quite strong, the metastable state is in fact quite stable.

Outside of the logarithmic region, the creep took on a different character. Figure 4.8 shows three different measurements at $T = 78$ K, the fields being 0.04 T, 0.08 T, and 0.2 T. The 0.04 T data is logarithmic in time throughout the entire measurement after the first three readings, but the higher field data are logarithmic in time only for about 20 minutes, after which the data bend away, decaying at a much slower rate which appears to be exponential at long times. Exponential behavior is shown in Fig. 4.9, where the field was 0.1 T. For 78 K, creep as seen in Fig. 4.8(b) was observed for all fields greater than 0.06 T.

For high temperature (78 K) measurements, the amount of change in the magnetization is much less than in Fig. 4.7. The 0.04 T data decayed 6.3% in one hour (slightly less than half that of the low temperature measurement), while the 0.08 T decayed 2.9% and the 0.2 T data decayed 1.1% in the same time. $S$ was much
Figure 4.7: Low temperature flux creep measurement at 20 K and 0.15 T, with $H \parallel c$
smaller at this temperature, being 0.5870 memu for 0.04 T, 0.217 memu for 0.08 T, and 0.0803 memu for 0.2 T. In the cases of non-logarithmic creep, $S$ is calculated from the initial, logarithmic data.

The flux creep measurements typically showed three kinds of time dependence, which are shown in Fig. 4.8. The data in the region 0 is usually observed at low fields $0 \leq H \leq H^*$, and is due to either the sample not being fully penetrated with flux (although in Fig. 4.8(a), this is not the case) or ambiguity in defining the time origin (which is uncertain by approximately 15 seconds), which affects the first few measurements but not the later ones. Data in region 1 are described well by the BLW model, and comprise the majority of the creep observed in this study. At high temperatures and fields, data was observed in region 2. Here, the BLW model does not describe the time dependence, and at long times the time dependence is closer to $e^{-t}$ than $\ln t$, as shown in Fig. 4.9. This exponential dependence is due to the effective pinning barrier being so low that vortices are able to diffuse backwards as well as forwards [112].

**Creep Rates** The attention will now turn to the creep rate as functions of temperature and field. The existence of three different regimes of flux creep behavior in ZFC measurements is established: onset of creep at low fields, a sharp peak in the creep rate $S$ at intermediate fields, followed by saturation and/or gentle drop-off at higher fields. These three regimes are clearly seen in the data for $T = 20$ K. The creep measurements are shown in Fig. 4.10, and $S$ vs. $T$ is shown in Fig. 4.11.

For very low fields ($\mu_0 H < 0.06$ T), little if any creep was observed, which implies that few if any vortices are in the sample, that is, $H \leq H_{c1}$. It is very tempting to
Figure 4.8: High temperature flux creep measurement at 78 K with (a) $\mu_0 H = 0.01$ T, (b) 0.08 T, and (c) 0.2 T, with $H \parallel c$. 
Figure 4.9: Creep measurement at 78 K and 0.4 T showing the initial logarithmic creep followed by exponential decay at long times.
Figure 4.10: Flux creep measurements on Tl-1 for $0.02 \leq \mu_0 H \leq 1.9$ T at $T = 20$ K
identify the field for which flux creep first occurs as $H_{c1}$. But as $H_{c1}$ is extremely difficult to determine, this field will be referred to as $H_{fe}$, that is, the field for which flux first enters the sample ("fe" $ightarrow$ flux entry).

Above $H_{fe}$, $S$ increases rapidly to a maximum near 0.25 T, followed by monotonic, nonlinear decrease with increasing field. This peak is probably associated with the arrival of the Bean flux front at the center of the sample, however, it does not occur at $H^*$, which at 20 K was 0.15 T. At 0.25 T, the value of $H/(dH/dx)$ is 27.7 $\mu$m, which is nearly three times the radius, so the sample is fully penetrated with flux before this peak. Also, derivations of the field dependence of $S$ by expanding the Bean model to account for a field-dependent $J_c$ [84] [85] usually result in a power law dependence of $S$ on $H$, such as $S \propto H^2$ [85], or $S \propto H^3 - H_{c1}^3$ [84]. The data diverge from the fit when $H \geq H^*$, and in Ref. [84], this occurs at fields below the
Figure 4.12: Low field creep rate at 20 K; the curves are attempts to fit to the models as discussed in the text.

Attempts to determine the low-field dependence for the 20 K data are shown in Fig. 4.12. Here, the applied fields have been corrected for demagnetization effects by assuming a spherical sample and using the formula $H = H_0 - N4\pi M$, where the demagnetization factor $N$ for a sphere is 1/3. The cubic dependence of Ref. [84] appears to describe the data best. This model also gives $H_{c1}$ as the field for which creep first occurs. Defining $H_{c1}$ in this way, at 20 K, $\mu_0 H_{c1} = 58.7$ mT. (The cubic field dependence was also observed at 10 K, giving $\mu_0 H_{c1} = 71.3$ mT.) The divergence of the data from the cubic fit occurs at approximately 0.15 T, which is very close to the hysteresis peak field. The log-log fit in Fig. 4.12 was done to find the “best” power law. However, this does not describe the very low field behavior, as
the log-log fit goes to zero at $H = 0$, and not at $H > 0$ as observed. The quadratic
fit of Ref. [85] is clearly seen not to fit the data. The conclusion of the matter here
is that $S$ peaks at $H > H^*$, making the determination of $H^*$ from this sort of plot a
nontrivial problem, and that the model of Yeshurun et al. [84] is the most useful for
describing the low-field behavior of this sample.

At high fields, $S$ either saturates or falls much more slowly than near the peak.
The creep rate will now reflect the dependence of $J_c$ and $U_{eff}$ on field, and will
not be affected by incomplete flux penetration. Figure 4.11 would indicate that for
$\mu_0 H > 0.8T$, $J_c(H)/U_{eff}(H) = 30S(H)/(rkT)$ is very weakly dependent on field,
or equivalently, $J_c$ and $U_{eff}$ have roughly the same field dependence.

Maxima occur in plots of both $S$ vs. $H$ and $S$ vs. $T$, which are shown in Fig. 4.13
and 4.14. An $H-T$ diagram can be constructed from the resulting field-temperature
Figure 4.14: Flux creep rate vs. $T$ for 0.1 T, 0.25 T, and 1.0 T points, and is shown in Fig. 4.15, along with $H^*$ as determined from the maxima in $J_C$. The error bars in field for the $J_C$ and $S$ vs. $H$ methods are typically less than 250 Oe, while for the $S$ vs. $T$ method, the error bars in temperature are less than 2.5 K. While the different curves are not within error bars of each other, they do not differ greatly, so taking $H^*$ from the peak in $J_C$, $S$ vs. $H$, or $S$ vs. $T$ is certainly reasonable. From now on, $H^*$ will be defined as the peak field in $J_C$ vs. $H$, because $H/(dH/dx)$ at the peak is close to the grain radius.

**Single Barrier Potential** To determine $U_{eff}$, the creep rates $S$ were combined with the hysteresis data in Eq. 2.23. It is more convenient to use Eq. 2.10 to replace $rJ_C$ with $\Delta M$, which gives the result

$$S = \frac{dM}{d\ln t} = \frac{\Delta M}{2} \frac{kT}{U_{eff}}.$$  \hspace{1cm} (4.1)
Figure 4.15: Full penetration field $H^*$ determined by hysteresis measurements along with the peak fields of the $S$ vs. $T$ and $S$ vs. $H$ plots.
Here, the symbol $U_{eff}$ is introduced to contrast the potential calculated from Eq. 4.2 with the potential shown in Fig. 2.3. $U_{eff}$ is a parameter obtained from creep data which gives the slope $dU(x)/dx$ that gives the measured creep rate $S$, as shown in Fig. 2.5. $U_{eff}$ would be the same as $U_0$ only if the true driving potential is linear in $J$. Eq. 4.2 could also be used for a single sample or a sample composed of grains of all the same size, and was used exclusively to calculate $U_{eff}$ (and $U_{eff}/kT$) from the creep data.

Inserting measured values of $\Delta M$ and $S$ into Eq. 4.2 often gives unreasonable values of $U_{eff}/kT$ in regions where the BLW model should not apply. For example, measurements for $H < H^*$, especially at low temperatures (10 and 20 K), usually gave $U_{eff}$ of approximately a few eV, which were 2 orders of magnitude larger than the results of measurements for $H > H^*$. In order to determine whether this was due to the measurement method, the creep in the hysteresis measurements were also analyzed. During the hysteresis measurements, only four or five measurements were taken at each field during, making it difficult to determine if the creep was actually logarithmic in time. The creep rates agreed well between the hysteresis and ZFC measurements, however, so this was not deemed to be a problem. In general, $S$ did tend to be larger for the hysteresis case, resulting in smaller $U_{eff}$ by 10 to 20%. Another method employed was to first apply a $-2$ T field, rather than starting from zero, followed by the field at which creep was to be observed. The comparison of the measurements are shown in Fig. 4.10, and it can be easily seen that the results agree within 20% (or less) as long as $H \geq H^*$. The differences here are in $S$, as the same $\Delta M$ is used for each calculation.

When $H < H^*$, the results were seen to be dependent on the method, which
Figure 4.16: $U_{eff}$ vs. field, $T = 20$ K, for the three measurement methods

means that this simple creep model can not be used to analyze data for $H < H^*$. This is a reasonable conclusion, because Eq. 4.2 was derived by assuming that the sample is fully in the critical state. This is not the case when $H < H^*$ during a ZFC measurement, so these results (in fact, most low field ZFC results) should be disregarded. The fact that $U_{eff}$ calculated from the hysteresis and initial $-2$ T application results do not agree indicates that it is extremely difficult to apply this model to any low field $H < H^*$ data, regardless of field history. Therefore, the only data from which meaningful results can be extracted is data at high fields $H > H^*$.

Given the approximate nature of $U_{eff}$ as calculated from Eq. 4.2, the 10 to 20 % variation is not significant. Therefore, the ZFC data is used in all discussions to follow. The ZFC data are considered to be most reliable because more data were
Figure 4.17: $U_{eff}$ vs. $B$ at 50 K and 78 K, with the fits to an inverse field law shown.

At low temperatures (especially 10 and 20 K), $U_{eff}$ varies slowly with field, and the dependence is qualitatively linear (Fig. 4.16, but at higher temperatures, the field dependence is much stronger. In Fig. 4.17, $U_{eff}$ is plotted vs. $H$ at $T = 50$ and 78 K. The curves are fits to an inverse field behavior, that is, $U_{eff} = A/B$. Here, the fields are corrected for demagnetization effects, using a spherical demagnetization factor ($N = 1/3$) (because the grains are all of different sizes and the exact dimensions are not...
known, a more accurate determination of \( N \) was not possible). Also, the exact value of \( H \) was seen to be of lesser importance, because \( U_{\text{eff}} \) can also be fit to \( 1/H_a \), which is probably due to the low value of the magnetization (< 50\% screening), even in low fields. Tinkham [90] has shown that if thermal activation occurs within a volume of \( (\Phi_0/B)\xi \), then \( U_{\text{eff}} \propto B^{-1} \), as seen here. This effect appears only at higher temperatures (50 and 78 K), where \( J_c < 10^5 \text{ A/cm}^2 \), so this field dependence would appear to hold only if the pinning is weak \( (U_{\text{eff}}/kT \leq 10) \). In fact, the data shown in Fig. 4.17 extend out into the region where the creep becomes non-logarithmic, where \( U_{\text{eff}}/kT < 2 \), and the creep model may not apply well here, as will now be discussed.

The low-field limit of applicability of Eq. 4.2 is \( H^* \). There also exists a high-field, high-temperature limit, governed by the onset of thermally activated flux flow (TAFF), where the rate of back diffusion of vortices is comparable to the rate of forward diffusion. By applying Eq. 4.2 to the data in Fig. 4.8, \( U_{\text{eff}}/kT = 4.2 \) for 0.04 T, \( U_{\text{eff}}/kT = 2.35 \) for 0.08 T, and \( U_{\text{eff}}/kT = 0.95 \) for 0.2 T. A very general result for this sample is that the sharp deviation from logarithmic behavior is seen to occur when \( U_{\text{eff}}/kT \approx 2 \), so that the data in Table 4.3 mark out the region of the \( H-T \) plane where \( U_{\text{eff}}/kT = 2 \). This transition is not a sharp one (the creep is not exactly linear when \( U_{\text{eff}}/kT = 4 \), for example), but this very sharp deviation always occurs when \( U_{\text{eff}}/kT \approx 2 \) and not before. The creep model predicts logarithmic behavior, so it stands to reason that it should not be applied to data where the creep is non-logarithmic.

This discussion has established that there exists in the \( H-T \) plane a limited region of applicability for the creep model. This region is bounded at low temperatures and
Figure 4.18: The region of the $H$-$T$ plane in which the creep model applies low fields by the full penetration field $H^*(T)$ and at high temperatures and high fields by a line of constant $U_{\epsilon \text{ff}}/kT = 2$. This region is shown in Fig. 4.18, and it can be seen that much of the $H$-$T$ plane lies outside of this region. The creep model can therefore only be applied to a relatively narrow part of the $H$-$T$ plane.

**Phase Diagram** One of the goals of this work was to establish the behavior of $U_{\epsilon \text{ff}}/kT$ in the $H$-$T$ plane. In Fig. 4.19, $U_{\epsilon \text{ff}}$ and $U_{\epsilon \text{ff}}/kT$ are both plotted vs. field for $T = 10$ and 20 K, and while $U_{\epsilon \text{ff}}$ is nearly equal at both temperatures, $U_{\epsilon \text{ff}}/kT$ is much larger at 10 K. This could explain why $J_c$ (Fig. 4.1) and the flux pinning force $F_p$ (Fig. 4.5) are greater at 10 K than at 20 K. Therefore, $U_{\epsilon \text{ff}}/kT$ is more meaningful as a parameter than $U_{\epsilon \text{ff}}$, and the behavior of $U_{\epsilon \text{ff}}/kT$ in the $H$-$T$ plane can be seen from the following $H$-$T$ phase diagram method. This is an
Figure 1.19: (a) $U_{eff}$ vs. $H$ (b) $U_{eff}/kT$ vs. $H$ at $T = 10$ K and 20 K. Note that while $U_{eff}$ is nearly the same for each temperature, $U_{eff}/kT$ is much larger at 10 K.
important distinction, because much attention is paid to the actual value of $U_{\text{eff}}$ in the literature (for example, Ref. [73]).

In order to study the behavior of $U_{\text{eff}}/kT$ in the $H$-$T$ plane, it is desired to find curves for which $U_{\text{eff}}/kT$ is constant. It has already been shown that the upper bound for applying the creep model is a line of constant $U_{\text{eff}}/kT$, in this case a line where $U_{\text{eff}}/kT = 2$. To do this, $U_{\text{eff}}/kT$ was plotted vs. temperature (Fig. 4.20), and by (visually) fitting a curve to the data in Fig. 4.20 and finding the intersection points of these curves with certain values of $U_{\text{eff}}/kT$ (i.e., $U_{\text{eff}}/kT = 10$), a line of constant $U_{\text{eff}}/kT$ can be found. These curves can then be plotted in the $H$-$T$ plane, as seen in Fig. 4.21.

These lines of constant $U_{\text{eff}}/kT$ characterize the flux creep in the $H$-$T$ plane. Several observations can be made. First, the curvature (qualitatively, the shape) of
Figure 4.21: $H-T$ phase diagram, showing the irreversibility transition, lines of constant $U_{\text{eff}}/kT$ ($5 \leq U_{\text{eff}}/kT \leq 20$), and the boundary curves $H^*(T)$ and the non-logarithmic transition ($U_{\text{eff}}/kT = 2$).
the lines is the same as that of the irreversibility transition. Because it has already been seen that $F_p$ is cut off at $H_{irr}(T)$, it is reasonable to expect that a quantity that is related to flux pinning should follow the same trend as $H_{irr}(T)$. Also, the behavior of $J_c$ as a function of field and temperature is reflected in this diagram. If the temperature is fixed at a low temperature (i.e., 10 K) and the field is increased, $U_{eff}/kT$ does not change much. As shown in Fig. 4.1, $J_c$ drops a factor of 5.8 at this temperature between $H_m$ and 1.9 T. If the field is held fixed at some field (i.e., 1.0 T) and the temperature is increased, $U_{eff}/kT$ falls rapidly from 20 $(T=10 K)$ to 2 $(T=35 K)$: in Fig. 4.1, $J_c$ falls about two orders of magnitude in this temperature interval.

The phase diagram approach shows that the creep model gives very meaningful results if applied within the proper region of the $H-T$ plane, even though this is the simplest model for creep. This method also gives an interpretation of the irreversibility line, which by virtue of the shape of the line of constant $U_{eff}/kT$, would also appear to a line of constant $U_{eff}/kT$. Here, it would be expected that $U_{eff}/kT \approx 1$, for in this case, the thermal energy would be equal to the well depth, which would effectively de-pin the vortices and allow the sample to become thermodynamically reversible, meaning that the irreversibility line is a flux-line lattice melting or depinning transition.

**Technical Difficulties** One technical difficulty in observing flux creep was discovered early on in this study. The scanning length used in early measurements was 9.0 cm, which was based on the default number of steps (128) and length of each step (0.07 cm) as programmed by Quantum Design. However, it was found that over
Figure 4.22: Comparison of flux creep measurements at 78 K and 0.4 T with 3.5 cm and 9.0 cm scan lengths, with H \perp c
this length, the magnetic field varies by 8.6% from the endpoints of the scan, whereas by limiting the scan length to 3.5 cm (the minimum length that the software used by QD1 can reliably calculate the magnetization), the field varies only by 0.2%. This difference had quite a striking effect on the behavior of flux creep, as can be seen in Fig. 4.22, in which two different creep measurements, both at 78 K and 0.2 T, are compared (here, \( H \perp c \)). The 9 cm scan data are extremely linear in \( \ln t \), whereas the 3.5 cm data are initially logarithmic, followed by exponential decay, as in Fig. 4.8. For a 9 cm scan, the sample is initially located in a field that is 8.6% lower than at the center. The magnetization of the sample will therefore be larger (if \( H > H^* \)) for the 9 cm scan than the 3.5 cm scan, as shown in Fig. 4.22. Logarithmic “creep” is observed all the way up to 2.0 T, even though the sample becomes reversible around 0.8 T. It appears that this “creep” is due only to the field gradient, and not anything uniquely related to the sample. This is a general result seen in all measurements for \( T = 78 \) K and \( \mu_0 H > 0.2 \) T, with \( H \perp c \).

To explain this effect, consider that during each scan, the sample is swept through the field gradient four times. Subjecting a sample that exhibits any hysteresis to such treatment will have the effect of cycling the field through a minor hysteresis loop between the central field and the initial field, which is 8.6% below the central field. Such a cycle is shown in Fig. 4.23(a). For the first scan of a measurement sequence, the cycle is: \( a \rightarrow b \rightarrow c \rightarrow d \) for the upstroke, and \( d \rightarrow a \rightarrow b \rightarrow c \rightarrow d \) for the downstroke. This leaves the sample in the state at \( d \), that is, on the reverse-field side of the hysteresis curve after each scan. For all subsequent scans, the cycle followed is the same as that of the downstroke.

Each scan took approximately 57 seconds to complete. The upstroke was stepped
Figure 4.23: Effects of a field gradient (a) Cycle followed for a sample scanned in a field gradient; (b) Time dependence of the field seen by the sample; (c) Time dependence of the magnetization corresponding to the field above

Note: no data are taken on the return stroke.
Table 4.1: Magnetization differences assuming the Meissner slope between hysteresis curves for 3.5 cm and 9.0 cm scanning lengths

<table>
<thead>
<tr>
<th>$\mu_0 H$ (T)</th>
<th>$\Delta(\mu_0 M)$ (T)</th>
</tr>
</thead>
<tbody>
<tr>
<td>3.5 cm</td>
<td>0.001 0.043</td>
</tr>
<tr>
<td>0.5</td>
<td>0.002 0.086</td>
</tr>
<tr>
<td>1.0</td>
<td>0.003 0.129</td>
</tr>
<tr>
<td>1.5</td>
<td>0.004 0.172</td>
</tr>
<tr>
<td>2.0</td>
<td></td>
</tr>
</tbody>
</table>

with a 0.035 cm interval, with 20 readings of the SQUID voltage at each step, while the downstroke was a continuous motion, so the time for the upstroke was 52 seconds and the downstroke 5 seconds. This results in a time-dependent field experienced by the sample. This is shown in Fig. 4.23(b), and the resulting time dependence of the magnetization is shown in Fig. 4.23(c).

To quantify the effect shown in Fig. 4.23, assume that while traversing from the lower to upper branch of the hysteresis curve, $|dM/dH| = 1/(4\pi)$, so $4\pi\Delta M = \Delta H = gH$, where $g$ is the field variation. Table 4.1 shows these numbers for 0.5, 1, 1.5, and 2.0 T, and it is very apparent that the effect is much, much stronger for the 9 cm scan. Even if $dM/dH$ is only a fraction of $1/(4\pi)$, the 9 cm scan would still subject the sample to a large variation in $M$ during each scan.

It can be easily seen that any appreciable field gradient will cause problems, especially in samples that have field-dependent hysteresis curves. Also, samples that exhibit large hysteresis will be affected more adversely than those with little hysteresis, meaning that measurements on strong-pinning superconductors are much more likely to suffer from these effects. As a result, all measurements were confined to scan lengths in which field gradients were minimized whenever measuring a sample that exhibits hysteresis (meaning any superconductor of interest here). For QD1, 3.5
cm has become the standard scan length, while for QD2, any length 3 cm or less is acceptable.

A second difficulty was with the temperature controller for QD1, which often made lower temperature measurements difficult, because it was not uncommon for the temperature to become unstable for temperatures less than 10 K. This is due to the percentage criterion for temperature stability. This means that if the temperature drifts by 10 mK at 10 K ($\Delta T/T = 0.001$), the temperature is considered unstable, while the same variation at 20 K is considered stable. For this reason, none of the data in this study were taken in this region. Also, there were troubles with the cooling mechanism, which resulted in large temperature overshoots.

**Nb-Ti**

**Overview**

As with the Tl$_2$Ba$_2$Ca$_2$Cu$_3$O$_{10}$ results, the Nb-Ti results will be presented in the same order: hysteresis, creep measurements, and analysis using the single barrier model. However, only one measurement method is used here, that being creep observed during hysteresis measurements, so there will not be any discussion as to the importance of field history. The sample studied was Nb-Ti-1, as described earlier. Because this sample is a classical superconductor, there will be more emphasis placed on comparing these results to those of Beasley et al. [70] on Pb-Tl alloys.

**Hysteresis Results**

To establish the field and temperature dependence of $J_c$ for Nb-Ti-1, hysteresis measurements were performed using QD2 at 4.5, 5, 6, 7, 7.5, 8, and 8.5 K, with
field intervals of 0.5 T. As in the Tl$_2$Ba$_2$Ca$_2$Cu$_2$O$_{10}$ measurements, the sample was first soaked with flux by applying the field cycle $0 \rightarrow -5.5T \rightarrow 0$, followed by the hysteresis loop. At each field, either 10 or 20 scans were recorded, which gave enough data to reliably observe creep. The scan length was limited to 2 cm, limiting the field variation to 0.005 %, and the moment was calculated from the data using the iterative regression algorithm. This scan length was less than that for Tl$_2$Ba$_2$Ca$_2$Cu$_3$O$_{10}$ because a 4 cm length was found to adversely affect the measurements, which will be discussed in detail at the end of this section. The sample was carefully centered to prevent spurious values of the magnetization resulting from poor centering and drift in the SQUID detector. Low fields $H < H^*$ were avoided, as time constraints and the questionable value of low field data made them impractical.

To calculate $J_c$ from the hysteresis data, Eq. 2.10 was used as before, using
Figure 4.25: $J_c$ vs. $H$ of Nb-Ti-1 at 1.5, 5, 6.7, 7.5, and 8 K

the filament radius of 20 μ. Because the filaments should be decoupled from each other in high fields, the total current in the sample is the sum of the currents in each filament. The magnetization curves are shown in Fig. 4.21, and $J_c(H)$ are shown in Fig. 4.25. The current densities at low temperature are well within the requirements for high-current applications, being $5.6 \times 10^5$ A/cm$^2$ at 4.5 K and 1.0 T, and $J_c$ falls only by a factor of three at 5 T. As a comparison, Finnemore et al. [63] found that at 4.2 K and 1.0 T for a Nb$_3$Sn sample specially prepared for high current capability, $J_c \approx 10^6$ A/cm$^2$, with similar field dependence. The reduced temperature here is $t = 4.5/9.1 = 0.18$. For Tl$_2$Ba$_2$Ca$_2$Cu$_3$O$_{10}$, the closest temperature that corresponds to this is would be 50 K, where $J_c$ falls from $1.88 \times 10^5$ A/cm$^2$ at 0.04 T to $1.18 \times 10^2$ A/cm$^2$ at 1.9 T, which is a drop of three orders of magnitude. This sample (Nb-Ti-1) qualifies as a typical high pinning superconductor.
as would be expected for a commercially available wire.

The field dependence of $J_c$ was found to obey $(1 - b)^2$, where $b = B/B_{c2}$ is the reduced field, and was approximately the same at each temperature. At low temperature (4.5 K and 5 K), this field dependence broke down at high fields ($B \sim 4$ T), but at higher temperatures all of the data for $H > H^*$ are well described by this fit. The low-field deviations are due to using the Bean model, $J_c = 15\Delta M/r$, when $H < H^*$, which is probably not valid (the maximum in $J_c$ at non-zero field is not a "fishtail" effect due to incomplete oxygenation). In Fig. 4.26, $J_{c0}^{1/2}$ is plotted vs. reduced field $b = B/B_{c2}$, showing linear behavior. $B_{c2}(T)$ was calculated from $1.6(1 - t^2)$, as found earlier.

The volume flux pinning force $F_p = |J_c \times B|$ was determined for each temperature. The upper critical field could be reliably determined, so it was possible to determine the field dependence of $F_p$. In Fig. 4.27, $F_p$ is plotted vs. field, showing that as the temperature is increased, both the maximum pinning force decreases and the peak in $F_p$ moves to lower field. By scaling the data to their maximum values ($F_p/F_{p,\text{max}}$) and plotting this vs. the reduced field $b = B/B_{c2}$, as in Fig. 4.28, it can be seen that the data follow the same field dependence, peaking at the same reduced field $b = 1/3$ for each temperature. The field dependence for these data was found to be closest to $b(1 - b)^2$, as shown for the 6 K data in Fig. 4.29. This is a reasonable result, given the definition of $F_p = J_c B$ and that $J_c \propto (1 - b)^2$.

The usual field dependence expected for Nb-Ti is $b(1 - b)$ [113] rather than $b(1 - b^2)$ as observed here. The usual field dependence for Nb$_3$Sn is $F_p \propto b^{1/2}(1 - b)^2$ [110] [114]. However, Fietz and Webb [81] found a scaling law in Nb-Ti alloys in which the peak pinning force $F_{p,\text{max}}$ is found to occur at the same reduced field regardless
Figure 4.26: Field dependence of $J_c$ shown by plotting $J_c^{1/2}$ vs. $b$ for (a) 4.5, 5 and 6 K and (b) 7, 7.5 and 8 K.
Figure 4.27: Volume flux pinning force $F_p$ vs. field for $4.5 \text{ K} \leq T \leq 8.5 \text{ K}$

Figure 4.28: Normalized volume flux pinning force $F_p/F_{p,\text{max}}$ vs. reduced field for $4.5 \text{ K} \leq T \leq 8.5 \text{ K}$, showing a common field dependence
Figure 4.29: Flux pinning force at 6 K. with the field dependence $b(1 - b)^2$ shown of temperature, as seen in Fig. 4.28. The scaling law is given as

$$F_p = [B_{c2}(T)]^n f(b).$$

where $f(b)$ is a function only of the reduced field $b = B/B_{c2}$ and $n$ ranges from 2 to 3. According to Kramer [110], $f(b)$ is determined by the sample microstructure resulting from the processing conditions, so that samples of the same material can obey the same scaling law yet have different field dependencies. For Nb-Ti, a very common value for the exponent in Eq. 4.3 is $n = 2.5$ [81]. To determine $n$, it is necessary to plot $F_{p,max}$ vs. $B_{c2}(T)$ on a log-log scale, which is shown in Fig. 4.30. The maximum measured values of $F_p$ were used at each temperature, and $B_{c2}(T)$ is calculated from $B_{c2} = 14.6(1 - t^2)$. The slope of the resulting straight line is $2.51 \pm 0.01$, which is very much in line with the findings of Fietz and Webb. Therefore, although the field
dependence of $F_p$ is not the same as other Nb-Ti samples, the scaling of $F_p$ with $B_{c2}$ is. This scaling behavior is also similar to the low temperature ($T \leq 35$ K) power law observed in Tl$_2$Ba$_2$Ca$_2$Cu$_3$O$_{10}$. although the exponent was $n = 2$ and the field was scaled to something other than $B_{c2} \,(T)$. (Fig. 4.3).

An irreversibility transition was observed in Nb-Ti-1, by extrapolating to $J_c = 0$ using the $J_{c}^{1/2}$ fitting. In Fig. 4.26, the fits do not extrapolate quite to $b = 1$, so the irreversibility transition was defined by multiplying the reduced field at which $J_c = 0$ by $B_{c2}$ at each temperature as determined by the parabolic fit to $B_{c2}$. A plot of $B_{irr}$ determined this way is shown in Fig. 4.31. The transition is quite narrow, where a typical value for $\Delta T \equiv T(B_{c2}) - T(B_{irr}) \sim 0.2$ K. This "transition" may not even be a true irreversibility transition, but may be due to using the parabolic fit.
100

Figure 4.31: Irreversibility transition of Nb-Ti-1, along with $B_{c2}(T)$

for $B_{c2}$ (which varies from the data by approximately 4%). It can be seen, however,
that any reversible region of the $H$-$T$ plane is very small. Irreversibility transitions
0.5 K to 1 K wide in Nb-Ti and Nb$_3$Sn have been observed by Suenaga et al. [113],
so while the reversible regions in conventional superconductors are not as wide as in
HTSC's, they do exist.

Flux Creep Results

Creep Measurements and Rates Creep measurements were taken during
the hysteresis measurements, with either 10 or 20 measurements were taken at each
field. ZFC measurements were not performed due to time constraints on the usage of
QD2 and because it has already been shown that as long as $H > H^*$, the results of the
creep measurements are independent of field history. Also, this was the measurement
Because Nb-Ti is a conventional superconductor, it was expected that the creep effect would be much smaller than in Tl₂Ba₂Ca₂Cu₃O₁₀, resulting in smaller changes in $M$, making the determination of $S$ somewhat less certain. In Fig. 4.32, creep measurements at 4.5 K and 4.0 T are shown for both increasing and decreasing field histories. These data have more scatter than values observed in Tl₂Ba₂Ca₂Cu₃O₁₀, but the logarithmic trend can still be clearly seen. Also found was that the creep rates are very small. The change in the magnetization per decade of time is about 6.3 G per decade for the increasing field measurement, and 3.9 G per decade for the decreasing field measurement. Here, $(1/2)\Delta M = 204$ G, so (taking the average of these two rates) 50 decades of time ($10^{50}$ minutes) would be required for the current to decay away and the sample to reach the equilibrium reversible state. This is much longer than in the case of Tl₂Ba₂Ca₂Cu₃O₁₀ at 0.15 T and 20 K (in fact, this is $10^{32}$ ages of the universe, assuming the universe is 20 billion years old), showing how much weaker the creep effect is in Nb-Ti.

The creep rates do not appear to have the same characteristics when plotted vs. field as in the case of Tl₂Ba₂Ca₂Cu₃O₁₀, as can be seen by comparing Fig. 4.11 to Fig. 4.33. Because the data were taken during hysteresis measurements, no incomplete penetration effects were observed, as in Tl₂Ba₂Ca₂Cu₃O₁₀. Also, the behavior of $S$ vs. $H$ changed with increasing temperature. At low temperature (4.5 K), the creep rate starts at minimum near $H^*$ (which is approximately 1.25 T), followed by a slow rise with field up to 5 T, as shown in Fig. 4.33(a). As the temperature increases, the creep rate tends to become relatively constant with field (Fig. 4.33(b)), and as $T \rightarrow T_c$, the creep rate is largest at zero field, and falls off to
Figure 4.32: Creep measurement on Nb-Ti-1 at 4.5 K and 4.0 T with measurements for both increasing and decreasing field histories shown.
Figure 1.33: \( S \) vs. \( H \) for Nb-Ti-1 at \( T = \) (a) 1.5 K, (b) 6.0 K, (c) 8.5 K
zero as $H \to H_{c2}$. No peak in $S$ near $H_{c2}$ is observed, in contrast to the observations of Beasley et al. This is probably due to the fact that $\kappa$ is much larger for Nb-Ti ($\sim 40$, which is more comparable to Tl$_2$Ba$_2$Ca$_2$Cu$_3$O$_{10}$ than to Pb-Tl) than the Pb-Tl alloys studied by Beasley et al.

**Single Barrier Model Analysis** In analyzing the data with Eq. 4.2, the average creep rate $\langle |S_+| + |S_-| \rangle / 2$ is used, as in Ref. [70]. This could be done for this sample, as opposed to Tl$_2$Ba$_2$Ca$_2$Cu$_3$O$_{10}$, where no reliable decreasing-field creep data were available. The field interval of 0.5 T was too coarse for an accurate determination of $H^*$, so $H^*$ could only be estimated from the peaks in the magnetization curves. Values of $H^*$ ranged from around 1.5 T at 1.5 K to less than 0.5 T at 8.5 K. The calculation of $U_{eff}$ was limited to fields above $H^*$.

The pinning barrier $U_{eff}$ is plotted as a function of field and temperature in Fig. 4.34. The values of $U_{eff}$ are not extremely large, ranging anywhere from 0 (at $H_{c2}$) to 90 meV. These values are of the same order as found in Tl$_2$Ba$_2$Ca$_2$Cu$_3$O$_{10}$, but because the data for Nb-Ti were taken at much lower temperatures than for Tl$_2$Ba$_2$Ca$_2$Cu$_3$O$_{10}$, $U_{eff}/kT$ is larger by a factor of 5 for Nb-Ti. Beasley et al. found $U_{eff}$ to be of order 1 eV for Pb-Tl alloys, which have a relatively low $\kappa$. If it is assumed that $U_{eff} \propto \xi$ [90], and for Nb-Ti, $\xi = 50 \AA$, and for Pb-Tl alloys, $\xi \sim 500 \AA$, then

$$\frac{U_{eff}(\text{Nb-Ti})}{U_{eff}(\text{Pb-Tl})} \approx \frac{\xi(\text{Nb-Ti})}{\xi(\text{Pb-Tl})} = 0.1 \ ,$$

and the factor of 10 difference in $U_{eff}$ between the two materials is very reasonable. While $U_{eff}$ is not especially large, $U_{eff}/kT$ is more than large enough to guarantee that $U_{eff}/kT \gg 1$, i.e., at 5 K and 5 T, $U_{eff} = 21$ meV, so $U_{eff}/kT = 18.7 \gg 1$. 

Figure 4.34: $U_{eff}$ vs. (a) field and (b) temperature for Nb-Ti-1
Figure 4.35: Creep model parameter $V_X$ as a function of reduced field for the temperatures shown.
Figure 4.36: Pinning length $X$ vs. $H$ for the temperatures shown

From Eq. 2.24, the parameter $VX$ can be calculated from the creep rates. This parameter is the product of the activation volume $V$ and the pinning length $X$. The $VX$ data are shown in Fig. 4.35. These values of $VX$ are a factor of from $10^{-5}$ to $10^{-3}$ smaller than found by Beasley et al., which again is due to the shorter coherence length of Nb-Ti (a typical factor would be $\xi(Nb-Ti)/\xi(Pb-Tl)^3 = 10^{-3}$). $VX$ also decreases with decreasing temperature, as found by Beasley et al.

Within Tinkham’s model, $V = \frac{a_0^2 \xi}{\Phi_0} = (\Phi_0/B)^{1/2}$, so $X$ may be estimated from these data. Although the data are noisy, it is apparent from Fig. 4.31 that $U_{c,eff}$ is a non-linear, decreasing function of $B$, so Tinkham’s model will be assumed to hold true. Dividing $VX$ in Fig. 4.35 by $V$ as defined above gives $X(B,T)$, which is shown in Fig. 4.36. In calculating the activation volume, $\xi(T)$ was calculated using the parabolic fit for $H_{c2} = 11.6(1 - t^2)$ (T) and the GL expression for $H_{c2}$. 
Inspection of Fig. 4.36 shows that \( X \) increases both as a function of temperature and field. The data for fields very near to \( H_{c2} \) are not shown, as in these cases, \( X \) became very large (several \( \mu m \)). In fact, these data may be somewhat questionable, as near \( H_{c2} \), \( X \) becomes comparable to the diameter of the filaments (40 \( \mu m \)).

The pinning length \( X \) is also the geometric width of the pinning barrier. \( X \) increases with both field and temperature, meaning that the potential barrier is stretched outwards with increasing field and/or temperature. For a fixed temperature, this means that as the field is increased, which increases the degree of tilting of the potential in Fig. 2.4, the individual pinning wells are being moved farther apart (the same is true for a fixed field and increasing temperature). This can explain the decrease in \( J_c \) with field and temperature, because \( F_p = \Delta U/\Delta x \) decreases with temperature. \( \Delta U \) decreases because \( H_c \) decreases \( (\Delta U \propto H_c^2) \), and \( \Delta x \) increases because the characteristic lengths \( \xi \) and \( \lambda \) diverge at \( T_c \).

**H-T Phase Diagram** The \( H-T \) phase diagram was determined in the same manner as before for \( \text{Tl}_2\text{Ba}_2\text{Ca}_2\text{Cu}_3\text{O}_{10} \), except that in this case, \( U_{eff}/kT \) was plotted vs. field for different temperatures. This was done solely because it was more convenient, and is shown in Fig. 4.37. Here, only the data for which \( U_{eff}/kT < 100 \) are shown, as only in this case did the data follow a regular trend (there was much scatter in the data for \( U_{eff}/kT > 100 \)). At 4.5 K, this condition is only met for 5 T, so the 4.5 K data were not used for this analysis. The best curve was visually fit to the data, and the intersection points of these curves at constant values of \( U_{eff}/kT \) gave the lines of constant \( U_{eff}/kT \), which are shown in Fig. 4.38.

The phase diagram is very different than that of \( \text{Tl}_2\text{Ba}_2\text{Ca}_2\text{Cu}_3\text{O}_{10} \). The lines
Figure 4.37: $U_{eff}/kT$ vs. $B$ for different temperatures (shown)

of constant $U_{eff}/kT$ lie very close to the $H_{c2}$ line, and the relative change in the pinning strength is much faster. Also, the values of $U_{eff}/kT$ is much higher than for Tl$_2$Ba$_2$Ca$_2$Cu$_3$O$_{10}$, even at low temperature where the reduced temperature for Tl$_2$Ba$_2$Ca$_2$Cu$_3$O$_{10}$ is a factor of 13 smaller than for Nb-Ti. Again, the phase diagram reflects the field and temperature dependence of $J_c$, because a measurable $J_c$ is observed throughout the $H$-$T$ plane all the way out to $H_{c2}$. Also, the lines have the same curvature and follow the same trend as the $H_{c2}$ line. Since the $H_{c2}$ line is the “irreversibility” line for this sample, it appears that in general the lines of constant $U_{eff}/kT$ for any sample will follow the irreversibility line, whether this be an actual thermodynamic reversible transition or the upper critical field.

All of the above data discussed above were taken with 2 cm scans. When a 4 cm scan length was used, the field variation, while small, had a large effect. The
Figure 4.38: \( H-T \) phase diagram showing \( B_{c2}(T) \) and the lines of constant \( U_{eff}/kT \) for \( 10 \leq U_{eff}/kT \leq 70 \)
magnetization was smaller than the 2 cm measurements on the decreasing-field part of the hysteresis curve, while for the increasing-field part, the data overlapped. This resulted in $\Delta M$ being reduced by anywhere from 2% at 2 T to 20% at 5 T, and the lines of constant $U_{\text{eff}}/kT$ shifted to lower temperature (for example, the temperature at which $U_{\text{eff}}/kT = 30$ with $B = 4$ T is depressed from 6.7 K for 2 cm to 4.7 K for 4 cm). The magnetization curves for 5 K are shown in Fig. 4.39(a), and the $H$-$T$ diagram for 4 cm scan is in Fig. 4.39(b).

**Comparisons**

The phase diagram approach was used to compare the different superconductors. The $H$-$T$ diagrams are plotted with the reduced temperature, because $T_c$ varies anywhere from 0.4 to 120 K. As a baseline for comparison, particular note of the position and curvature of the lines of $U_{\text{eff}}/kT = 10$ will be made.

**Tl$_2$Ba$_2$Ca$_2$Cu$_3$O$_{10}$ - Nb-Ti**

Although the effective pinning energies for Tl$_2$Ba$_2$Ca$_2$Cu$_3$O$_{10}$ are similar to Nb-Ti, the $H$-$T$ diagrams for both samples are very different, as shown in Fig. 4.10. It is quite apparent that the shorter coherence length of Tl$_2$Ba$_2$Ca$_2$Cu$_3$O$_{10}$, combined with the anisotropy in $J_c$ and structure have had extremely detrimental effects on the flux pinning.

For Tl$_2$Ba$_2$Ca$_2$Cu$_3$O$_{10}$, the lines of constant $U_{\text{eff}}/kT$ follow $H_{irr}(T)$ and have positive curvature, but for Nb-Ti, the lines have negative curvature and follow $H_{c2}(T)$. The $U_{\text{eff}}/kT = 10$ line for Tl$_2$Ba$_2$Ca$_2$Cu$_3$O$_{10}$ crosses into a region of high field (i.e., 1 T) at a reduced temperature of $\sim 0.15$, which corresponds to
Figure 4.39: (a) Comparison of hysteresis curves at 5 K for 2 cm and 4 cm scans; (b) $H-T$ diagram for 4 cm scan
Figure 4.10: $H-T$ diagram for Nb-Ti and Tl$_2$Ba$_2$Ca$_2$Cu$_3$O$_{10}$ plotted on a reduced temperature scale

$T = 18$ K, while for Nb-Ti, the lowest reduced temperature that this line reaches is $t = 0.85$ ($T = 8$ K, $\mu_0 H = 4$ T). If Tl$_2$Ba$_2$Ca$_2$Cu$_3$O$_{10}$ is to be used for practical applications, some method must be implemented to give a line of $U_{c f f}/kT = 100$ at higher reduced temperature than the lines in this study occur.

Tl$_2$Ba$_2$Ca$_2$Cu$_3$O$_{10}$ - (Nd$_{1-x}$Ce$_x$)$_2$CuO$_4$

Flux pinning and critical current studies in (Nd$_{1-x}$Ce$_x$)$_2$CuO$_4$ provide another interesting contrast to Tl$_2$Ba$_2$Ca$_2$Cu$_3$O$_{10}$. While both samples were quite reversible, they differed greatly in $T_c$ (120 K vs. 21 K) and carrier type (Tl$_2$Ba$_2$Ca$_2$Cu$_3$O$_{10}$ is a $p$-type superconductor while (Nd$_{1-x}$Ce$_x$)$_2$CuO$_4$ has both $n$ and $p$-type carriers), so it is interesting to study how the $H-T$ diagrams compare.
Figure 4.41: $H$-$T$ diagram for $\text{Tl}_2\text{Ba}_2\text{Ca}_2\text{Cu}_3\text{O}_{10}$ and $(\text{Nd}_{1-x}\text{Ce}_x)\text{Cu}_4$ plotted on a reduced temperature scale.
Figure 4.11 shows the $H$-$T$ diagrams for both samples, with $H \parallel c$. Even though $T_c$ differs by a factor of 6 between these two samples, the diagrams look surprisingly similar on a reduced temperature scale, the lines possessing positive curvature and nearly coinciding. This is a rather surprising result, as the coherence length would be expected to be a factor of 6 longer for $(\text{Nd}_{1-x}\text{Ce}_x)_2\text{CuO}_4$ and $U \propto \xi$. Both of these superconductors, however, are anisotropic, which would indicate that $U_{\text{eff}}/kT$ is affected more strongly by anisotropy than by $\xi$. The $U_{\text{eff}}/kT = 10$ lines for both samples nearly coincide.

The anisotropy in the lines of constant $U_{\text{eff}}/kT$ for $(\text{Nd}_{1-x}\text{Ce}_x)_2\text{CuO}_4$ is shown in Fig. 4.42. While the curvature of the lines are the same, it is easily seen that $U_{\text{eff}}/kT$ is consistently larger for $H \perp c$ than for $H \parallel c$. In particular, the
Figure 1.13: $H$-$T$ diagram for $(\text{Nd}_{1-x}\text{Ce}_x)_2\text{CuO}_4$ and $(\text{Ba}_{1-x}\text{K}_x)\text{BiO}_3$ plotted on a reduced temperature scale.

$U_{\text{eff}}/kT = 10$ line for $H \perp c$ practically falls right on top of the $U_{\text{eff}}/kT = 2$ line for $H \parallel c$.

$(\text{Nd}_{1-x}\text{Ce}_x)_2\text{CuO}_4 - (\text{Ba}_{1-x}\text{K}_x)\text{BiO}_3$

To further test the effects of anisotropy, the comparison is made between two samples of similar $T_c$ (and $\xi$), but different crystal structures, $(\text{Ba}_{1-x}\text{K}_x)\text{BiO}_3$ being cubic and $(\text{Nd}_{1-x}\text{Ce}_x)_2\text{CuO}_4$ being tetragonal. The $H$-$T$ diagram is found in Fig. 4.13.

The effect of anisotropy is very apparent in this figure, the reversible region of the $H$-$T$ plane being much smaller for the cubic $(\text{Ba}_{1-x}\text{K}_x)\text{BiO}_3$ than the tetragonal $(\text{Nd}_{1-x}\text{Ce}_x)_2\text{CuO}_4$. The lines for $(\text{Ba}_{1-x}\text{K}_x)\text{BiO}_3$, while having similar curvature
to those of $(\text{Nd}_{1-x}\text{Ce}_x)_2\text{CuO}_4$, lie well above those of $(\text{Nd}_{1-x}\text{Ce}_x)_2\text{CuO}_4$. There
is no $U_{\text{eff}}/kT = 10$ line shown on this figure, as $U_{\text{eff}}/kT$ was so large for this
sample the $U_{\text{eff}}/kT = 10$ was extremely difficult to estimate. The $U_{\text{eff}}/kT = 20$
line is much closer to that of the Nb-Ti sample than any other of the materials studied
here. This material may be a used in the future as a magnet wire material, if the
granularity and other metallurgical problems can be solved.

\textbf{Tl}_2\text{Ba}_2\text{Ca}_2\text{Cu}_3\text{O}_{10} - \text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_x

The final comparison to be made is between two materials with similar crystal
structures and relatively large $T_c$'s. Both \text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_x and \text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_x
have tetragonal structures. Although the values of $U_{\text{eff}}/kT$ for \text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_x
were derived from a different experiment, it can be seen in Fig. 4.14 that these num­
bers are comparable. The constant $U_{\text{eff}}/kT$ lines have the same positive curvature,
but the \text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_x lines are consistently above the \text{Tl}_2\text{Ba}_2\text{Ca}_2\text{Cu}_3\text{O}_{10} lines.
This may be due to the lower $T_c$ of \text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_x.

\textbf{Summation}

To conclude this section, an $H$-$T$ diagram with lines of $U_{\text{eff}}/kT = 10$ is
shown in Fig. 4.45. The cubic $(\text{Ba}_{1-x}\text{K}_x)\text{BiO}_3$ line lies below the Nb-Ti line, fol­
lowed by the \text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_x line near the center, and the \text{Tl}_2\text{Ba}_2\text{Ca}_2\text{Cu}_3\text{O}_{10}
and $(\text{Nd}_{1-x}\text{Ce}_x)_2\text{CuO}_4$ lines near the bottom of the diagram. Both Nb-Ti and
$(\text{Ba}_{1-x}\text{K}_x)\text{BiO}_3$ are cubic, lacking any sort of layered structure, while the layered
superconductors have Cu-O planes, two for \text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_x and three for each of
\text{Tl}_2\text{Ba}_2\text{Ca}_2\text{Cu}_3\text{O}_{10}$ and $(\text{Nd}_{1-x}\text{Ce}_x)_2\text{CuO}_4$. Therefore, it is apparent that struc-
Figure 4.44: $H$-$T$ diagram for Tl$_2$Ba$_2$Ca$_2$Cu$_3$O$_{10}$ and Bi$_2$Sr$_2$CaCu$_2$O$_{x}$ plotted on a reduced temperature scale.
Figure 4.45: $U_{\text{eff}}/kT = 10$ lines for all samples discussed here. For the anisotropic superconductors, only the $H \parallel c$ lines are shown.
tural anisotropy is the most important factor in characterizing the flux creep in these samples, with $T_c$ (or $\xi$) having lesser importance. (Sample microstructure could also be an important factor in the flux creep characteristics, however.) If the HTSC's are to be used for practical applications, the $\beta_{eff}/kT = 10$ for these materials must be shifted to higher reduced temperature.
CHAPTER 5. DISTRIBUTION MODEL RESULTS

The Hagen-Griessen (HIG) [75] model will be used here to extract the distribution of activation energies $m(E_0^*)$ from the creep data for Tl$_2$Ba$_2$Ca$_2$Cu$_3$O$_{10}$, with $H \parallel c$. The distributions will be calculated for three fields: 0.1 T, 0.25 T, and 1.0 T, to observe any field dependence in $m(E_0^*)$. As in the case of the single barrier analysis, the data used for this analysis were taken at $t_b = 10$ minutes after stabilization of the field and temperature to ensure the proper SQUID range and to allow any transients to die out.

Curve Fitting

In order to apply the HIG model, it is necessary to obtain the temperature derivatives of $M(T)$ and $S = dM/d\ln t(T)$. To do this, creep data were recorded every 5 K from 10 K to near $T_{trr}(H)$ for each field on QD1. These data were taken from ZFC measurements, to be consistent with Ref. [75]. Curves were then fit to the data using polynomial least-squares fitting. The number of terms in the polynomial was chosen by finding the largest number that did not result in oscillations between the data. After the curves were fit, the actual values of the data were ignored. The derivatives were easily calculated once the functional dependence was known, and the curves as fit to the data are shown in Fig. 5.1. It was necessary to fit the $S$ vs. $T$ data in...
Figure 5.1: (a) $M(T)$ and (b) $S(T)$ with the fitted curves shown
two parts, but a single fit was possible for the magnetization. The magnetization fitted here was $(1/2)\Delta M$, to eliminate the possibility of including any normal state or reversible superconducting signal in the analysis.

**Determination of $\ln(t_b/\tau)$**

Equation 2.31 was used to calculate $\ln(t_b/\tau)$, using Eqs. 2.32 and 2.33 for the functions $a(T)$ and $b(T)$. The choice of the exponents $n$ and $m$ was made by using both the value of $\ln(t_b/\tau)$ (HIG found typical values of $\sim 30$) and the degree to which $\ln(t_b/\tau)$ is independent of temperature. Unfortunately, while values of $\sim 30$ are easily obtained, $\ln(t_b/\tau)$ is rarely independent of temperature, as is shown in Fig. 5.2. The “best” choices of $n$ and $m$ were $n = 1$ and $m = 1$. Again, $n = 1$ corresponds to the model given by Tinkham [90], and since this model was shown to approximate the data, this reinforces this choice of $n$.

The physical interpretation of $n = 1$ is that the activation volume $V$ is proportional to $\xi$. Because $m = 1$, and because $a(T)$ describes the temperature dependence of $V\xi$, and it is assumed that the volume $V$ that appears in $a(T)$ is the same volume as in $b(T)$, the pinning length $X$ is independent of $\xi$. meaning that the sample microstructure determines the size and structure of the pinning potentials.

To calculate the activation energy $E^{\ast}_{0}$ (Eq. 2.28), the average value of $\ln(t_b/\tau)$ was used. This varied by about 15% from field to field, so the average value of $\ln(t_b/\tau)$ for each field were averaged again. This gave $\ln(t_b/\tau) = 30.07$, which corresponds to $\tau = 5.2 \times 10^{-11}$ sec.
Figure 5.2: $\ln(t\gamma/\tau)$ vs. $T$ for (a) fixed $n = 1$ and (b) fixed $m = 1$ for $\mu_0 H = 0.1$ T. Similar behavior was observed for the other fields.
Distributions

The distribution function $m(E_0^*)$ was calculated using Eq. 2.30, using $E_0^*$ from Eq. 2.28 and $\ln(t_b/\tau) = 30.07$, and are shown in Fig. 5.3. The distributions were calculated down to zero energy (corresponding to zero temperature) by extrapolating the fits for $S(T)$ and $M(T)$ to $T = 0$. Since these fits are questionable (at best) at low temperature, the very low energy results should be ignored.

Although the distributions look qualitatively like those in Ref. [75], there are some important differences to note. While Ref. [75] reported zero occupation at low energy, here it can be seen that for 0.1 T and 0.25 T, there appears to be a negative occupation of states at low energies, which is clearly an unphysical result. The scale of the 0.1 T distribution is lower than that of 0.25 and 1.0 T, but is similar to the 0.1 T distributions in Ref. [75]. All three distributions have a peak at some non-zero energy which decreases with increasing field.

Because there can not be a negative occupation of states ($m(E_0^*)$ is either zero or positive), some assumptions of the model must disagree with the data. The peaks in the distributions occur at energies corresponding to the temperature at which $S(T)$ peaks, so that at temperatures below this peak, the field is probably less than $H^*$ and the sample is not fully penetrated with flux. This model assumes that the sample is fully penetrated, so data below the $S(T)$ peak should not be used in this calculation. This means that there is no information about the low energy occupation of states to be gained from this analysis. A more reasonable expectation is that the $m(E_0^*)$ is constant at low energy [111] up to the peak energy, and falls off with increasing energy, as is indicated in Fig. 5.3 by the dashed lines.

One interesting fact to note is that it is possible to force $m(E_0^*)$ to remain
Figure 5.3: HHG distribution functions \( m(E^*_0) \) vs. \( E^*_0 \) for (a) \( \mu_0 H = 0.1 \) T, (b) \( \mu_0 H = 0.25 \) T, and (c) \( \mu_0 H = 1.0 \) T.
Figure 5.1: IIG distribution function with the data altered as mentioned in the text positive by adding a constant fudge factor to $S$ (for example, replacing $S$ with $S + 11$ will cause the 0.1 T distribution to remain positive). Some manipulation of the IIG equations shows that with $m = 0$, if

$$
\frac{S}{T} \geq \frac{dS}{dT}.
$$

then $m(E^*_0) \geq 0$, as shown in Fig. 5.4 for 0.1 T. A log-normal distribution (Ref. [75]) could then be fit to these “data”.

While this model is incomplete for low energies, there still is useful information to be obtained from these distributions. The peak activation energy is depressed in higher fields, indicating that $E^*_0$ is a decreasing function of field. Also, $m(E^*_0)$ as shown is for $T = 0$. At finite temperature $T$, states with energies corresponding to temperatures less than $T$ will be unavailable. If $E^*_0(T)$ is above the peak in $m(E^*_0)$,
then the majority of pinning energy states will not be available, and the ability of the sample to pin flux will be greatly reduced. This is reflected in the considerable temperature dependence of $J_C$, and the low values of $J_C$ that occur at temperatures above the $S(T)$ peak.

To conclude this discussion, the IIG distribution model can be applied to creep data to extract the distribution function $m(E^*_0)$. However, since the assumption of a constant $\ln(t_b/\tau)$ is not observed, and the model yields unphysical results for low temperature-low field data, the utility of this model may be questioned. One fact that is apparent is that neither the BLW single barrier nor the IIG distribution model is adequate for the analysis of data for which the sample is not fully penetrated with flux. Further theoretical modeling is necessary before this can be rectified. Since the BLW model yielded satisfactory results with fewer assumptions necessary, this model is preferred over the IIG model.
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[80] C. P. Bean. Talk given at the 1990 March Meeting of the APS, Anaheim CA.


