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Jerry P. Gollub
Haverford College

M. R. Beasley

M. Tinkham

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Apparent Universal Behavior of Fluctuation-Induced Diamagnetism in Superconductors*

J. P. Gollub,† M. R. Beasley, and M. Tinkham
Department of Physics and Division of Engineering and Applied Physics,
Harvard University, Cambridge, Massachusetts 02138
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The dependence on temperature and magnetic field of the diamagnetism due to fluctuations above $T_c$, while showing marked deviations from the behavior expected on the basis of the Ginzburg-Landau theory, appears to follow a universal behavior if one introduces a new characteristic field $H_\phi$, which may be expressed in terms of other known material-dependent parameters.

We have previously reported\(^1\) observation of a temperature-dependent diamagnetism above the critical temperature $T_c$ due to thermal fluctuations of the superconducting order parameter $\Psi$. We have since extended these measurements to other materials (including type-II superconductors) and to much higher fields and temperatures; these results will be reported in detail elsewhere.\(^2\)

The purpose of this note is to show that these data are well described by an apparently universal function of appropriately scaled field and temperature variables. This empirically derived function deviates markedly from the behavior predicted\(^3\) on the basis of the simple Ginzburg-Landau (GL) theory, especially for $T > 2T_c$ and for fields comparable with $H_{c2}(0)$. These deviations demonstrate the expected breakdown of the GL theory in these regimes where short-wavelength fluctuations dominate, and where consequently the slow-variation approximations of the GL theory break down.\(^4\) The apparent generality of our results suggests that there may exist a reasonably simple extension of the GL theory to deal with situations in which $\Psi$ varies rapidly on the scale of $\xi(0)$, the zero-temperature GL coherence length. The attempt of Patton, Ambegaokar, and Wilkins\(^5\) (PAW) to deal with the problem by means of an ad hoc cutoff energy is shown to have qualitative, but not quantitative, success.

Prange’s calculation\(^6\) of $M'$, the magnetization due to thermal fluctuations, is exact within the framework of the GL theory. It is found that $M'$ diverges at the temperature $T_{c2}(H)$, defined by the condition $H = H_{c2}(T_{c2})$, where the energy cost of a small fluctuation toward the superconducting state vanishes. Above $T_{c2}$, $M'(T)$ is predicted to fall off roughly as $(T-T_{c2})^{-1/2}$. In the temperature range near $T_c$ where $H_{c2}(T)$ is linear in $T$, it is predicted that a field- and material-independent curve should be obtained if $M'/H^{1/2}T$ is plotted versus the dimensionless scaled temperature difference $(dH_{c2}/dT)(T-T_c)/H$. In particular, at the zero-field critical temperature $T_c$, it is predicted that

$$ M'/H^{1/2}T_c = -0.323\Phi_0^{-3/5}k_B $$

(1)

is a universal constant for all superconductors. Here $\Phi_0$ is the flux quantum $hc/2e$. While our data show some of the qualitative features of Prange’s result, they do not agree quantitatively nor do they scale so simply, except perhaps near $T_c$.

The PAW theory attempts to correct for the overestimate in GL theory of short-wavelength fluctuations by introducing an unknown energy-cutoff parameter $E$ into the fluctuation spectrum. As a result, they find that the magnetization is strongly depressed below the Prange value if $H > H^* = mcE/\hbar c$ or if $T - T_c > T^* = 4m\xi^2(0)T_cE/\hbar^2$, where $2m$ is the electronic pair mass. The expectation that the GL theory should break down when the characteristic wavelength over which the order parameter varies is smaller than $\xi(0)$ suggests that $E = \hbar^2/4m\xi^2(0)$, which corresponds to $H^* = H_{c2}(0)$ and $T^* = T_c$. Wavelengths as short as $\xi(0)$ occur for even the least energetic modes.
FIG. 1. The field dependence of \(-M'(T_c)/H^{1/2}T_c\) plotted in units of the scaling field \(H_s\) and compared with the PAW falloff from Prange’s theory.

of fluctuations when \(H > H_c(0)\) but they occur at lower fields for the more energetic fluctuations, which vary more rapidly in space. Thus in the PAW calculation the cutoff does not become fully effective in reducing \(M'\) until the energies of even the lowest modes exceed the cutoff energy.

We have measured the temperature dependence of the magnetization at fields \(H\) up to 300 Oe and temperatures up to 16 K using instrumentation based on a superconducting quantum interference magnetometer. Our samples were long cylindrical single crystals (5 \(\times\) 40 mm) of In, Pb, and PbTl alloys, grown from 99.9999% purity components. We take \(M'(H, T) = M(H, T) - M(H, T \gg T_c)\) in order to eliminate the temperature-independent magnetization of the normal state. This value of \(M'\) is quite well defined since our present measurements extend far above \(T_c\) where \(M\) is constant within experimental error over a range of several degrees Kelvin except in the highest field measurements on Pb and PbTl, where a slightly temperature-dependent moment was still observable at 16 K. These more extensive measurements almost completely eliminate the baseline uncertainty which concealed the disagreement with the Prange theory in our earlier work.

In Fig. 1 we present the field dependence of \(M'(H, T_c)\) for three different materials. The ordinate \(M'(T_c)/H^{1/2}T_c\) would be field independent if Prange’s theory were valid. In fact the ordinate is progressively depressed as \(H\) increases. We find that for a given \(H\), In deviates the most and Pb-5%Tl the least from Prange’s theory. However, if the field is plotted in units of a scaling field \(H_s\), which is chosen to be 2.1 Oe for In, 36 Oe for Pb, and 720 Oe for Pb-5%Tl, then the data for the three materials coincide and fall below Prange’s theory in the same fashion. With our definition of \(H_s\), the ordinate is half of Prange’s value at \(H/H_s = 1\); it is nearly zero when \(H = 20H_s\), which incidentally is about equal to \(H_c(0)\) for the pure materials. The observed fall is very slow, being roughly logarithmic with field, so that the fall from 80 to 20% of the Prange value stretches over two decades in \(H/H_s\). The PAW prediction is also shown, with \(H^*\) identified with 1.1\(H_s\) to force agreement at \(H/H_s = 1\). It is not known whether the data would coincide with Prange’s result in the limit as \(H/H_s \to 0\), since accurate data at \(T_c\) can only be obtained in fields large enough to shift \(T_c(H)\) down by at least the width of the transition due to sample inhomogeneity.

We also find that for a given value of \(H/H_s\) the temperature dependence of \(M'\) exhibits universal behavior. This is shown for two values of \(H/H_s\) in Fig. 2, where the temperature and magnetization are plotted in the scaled variables suggested by Prange’s theory. The data lie substantially below Prange’s theory but for a given \(H/H_s\) the temperature dependence is the same for all materials. We consider this universal behavior to be remarkable in view of the fact that the data
involve clean type-I superconductors with both
strong and relatively weak electron-phonon cou-
pling and also a type-II alloy superconductor.

We conclude from these comparisons that to a
good approximation, \( M'(H, T)/H^{1/2} \) depends only
on two parameters, a scaled temperature \( \Delta H_{c2}/
\) and a scaled field \( H/H_s \), where \( H_s \) is a
single parameter for each material. The PAW
theory suggests the correct sort of reduced
variables to obtain universal behavior and qualit-
atively describes the suppression of fluctuations
at high fields and high temperatures. However,
in attempting to fit \(^a\) the PAW theory to the data
we are forced to use an \( H^*(mH_s) \propto H_{c2}(0) \). With
our interpretation of the cutoff energy \( E \), this
suggests an unreasonably long characteristic
cutoff wavelength. An alternative, and perhaps
more reasonable, viewpoint is that the PAW type
of cutoff does not start to suppress the higher
energy modes soon enough. Our data on the
type-I superconductors show that for \( H > H_{c2}(0) \),
fluctuations have been almost completely sup-
pressed, whereas, as noted before, the PAW
cutoff [with \( H^*=H_{c2}(0) \)] only begins to suppress
\( M' \) strongly when \( H \sim H_{c2}(0) \).

The nature of the parameter \( H_s \) is not entirely
understood. We find that \( H_{c2}/H_{c2}(0) \) is 0.05 for
In, 0.06 for Pb, and 0.4 for Pb-5% Tl. Not only
is \( H_s \) much less than \( H_{c2}(0) \) for all three materi-
als, but the alloy differs from pure lead more
than a simple proportionality to \( H_{c2}(0) \) would
lead one to expect. Empirically, we find that to
a good approximation

\[
H_s = C(\Phi_0/2\pi\xi_0^2)(1 + \xi_0/\xi_d)^2,
\]

where \( C \approx 0.083 \). This suggests that in accounting
for the suppression of fluctuations by a magnetic

\[
\text{Table I. Superconducting material parameters for our samples.}
\]

<table>
<thead>
<tr>
<th>Parameter</th>
<th>In</th>
<th>Pb</th>
<th>Pb-5% Tl</th>
<th>Unit</th>
</tr>
</thead>
<tbody>
<tr>
<td>( T_c )</td>
<td>3.409 (^a)</td>
<td>7.2 (^a)</td>
<td>7.1 (^a)</td>
<td>K</td>
</tr>
<tr>
<td>( \kappa(T_c) )</td>
<td>0.0620 (^b)</td>
<td>0.328 (^c)</td>
<td>1.19 (^a)</td>
<td></td>
</tr>
<tr>
<td>( \Delta H_{c2}(\kappa T_c)/T_c )</td>
<td>13.4 (^d)</td>
<td>111 (^c)</td>
<td>402 (^a)</td>
<td>Oe/K</td>
</tr>
<tr>
<td>( H_{c2}(0) )</td>
<td>45 (^e)</td>
<td>594 (^c)</td>
<td>1800 (^f)</td>
<td>Oe</td>
</tr>
<tr>
<td>( H_s ) (meas)</td>
<td>2.1 \pm 0.2</td>
<td>36 \pm 4</td>
<td>720 \pm 150</td>
<td>Oe</td>
</tr>
<tr>
<td>( H_s ) (from (2))</td>
<td>2.07</td>
<td>36.4</td>
<td>616</td>
<td>Oe</td>
</tr>
</tbody>
</table>

\(^a\) From this work.
\(^d\) \( \Delta H_{c2}(\kappa T_c)/T_c = \sqrt{2\kappa(T_c)\Delta H_{c2}/dT_c} \).
\(^e\) Estimated using \( \kappa(0) \approx 0.11 \) (from Ref. b).
\(^f\) Estimated from data on \( H_{c2}(0) \) vs \( \kappa \): G. Bon Mardion, B. B. Goodman, and A. Lacaze, J. Phys. Chem. Solids 26, 1143 (1965).
with GL occurs in the type-II alloy near $T_c$, this is reasonable since as one approaches $T_c$, the dominant contribution to the diamagnetism is from the very lowest modes, which are diverging there, and these lowest modes have the lowest-wavelength spatial variation. On the other hand, far from $T_c$, a great number of modes with much shorter wavelength make contributions comparable with the lowest ones, and the breakdown of the GL approximation is more severe. Because our sensitive measurements allow $M'$ to be followed out to $2T_c$ and to high fields, and because they are obtained with rather ideal bulk samples, our data should allow quite a critical test for theoretical treatments which go beyond the region near the critical point.

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†Present address: Physics Department, Haverford College, Haverford, Pa.


2J. P. Gollub, M. R. Beasley, and M. Tinkham, to be published. [For a detailed account of the results to date, see J. P. Gollub, thesis, Havard University, 1970 (unpublished).]


It is interesting to note that the PAW curve agrees quite well with the data if the horizontal scale is expanded by a factor of 2; i.e., if $(H/H_s)_{\text{t}}$ were replaced by $(H/H_s)_{\text{t}}^{1/2}$ in their results. Although this observation may well give guidance on how to improve their model, no rationale for such a change is presently available.

8For the pure materials, where $l \gg \xi_s$, $H_s$ depends only on $\xi_s$. This parameter was evaluated using the standards results $\xi(T) = 0.74\xi_s (1-T/T_c)^{3/2}$, and $H_c(\xi) = \phi_0/2\pi^2(0.74)^2\xi_s^2$, from which $\xi_s$ can be determined using the known values of $H_c(\xi)$ and $T_c$. For the alloy, $H_s$ is calculated by multiplying the value of $H_s$ computed for pure lead by the factor $(1 + \xi_s/\xi)^2$. The required ratio of $\xi_0/\xi$ can be determined from the relation $(\xi dH_c/dT)(\xi_0 H_c(\xi_0) = \chi(\xi)$, where $\chi(\xi)$ is a known function [see Ref. 5, p. 338] of $\xi_0/\xi$. These procedures yield $\xi(\ln) = 3640 \text{ Å}, \xi_0(\text{Pb}) = 870 \text{ Å}, \text{and } \xi_0/\xi = 3.12$ for the alloy, and the values of $H_s$ shown in Table I.

Hypersound Attenuation in Superconductors by Quasiparticle Creation*

M. P. Garfunkel, J. W. Lue, and G. E. Pike†

University of Pittsburgh, Pittsburgh, Pennsylvania 15213

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The attenuation of 10-GHz longitudinal sound waves in superconducting molybdenum ($T_c = 0.914$ K) and cadmium ($T_c = 0.500$ K) shows the high-frequency behavior predicted by the BCS theory. In particular, the onset of the rapid drop in attenuation with decreasing temperature that is characteristic of superconductors is shifted downward to the temperature $T^*_s$ (0.905 K in molybdenum and 0.490 K in cadmium) at which the superconducting energy gap equals the phonon energy. The analysis of the measurements indicates a large anisotropy in the energy gaps of both metals.

In the original publication of the BCS theory of superconductivity, the low-frequency limit (i.e., phonon energy $h\nu$ small compared with the energy gap $2\Delta$) for the attenuation of sound was written as

$$\alpha_s/\alpha_N = 2f(\Delta),$$

where $\alpha_s$ and $\alpha_N$ are the acoustic attenuations in the superconducting and normal states, and $f$ is the Fermi function. For higher frequencies the results include contributions to the phonon absorption not