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Final "draft" 5/27/83

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CONF-830776--2

DE83 015506

SUPERCONDUCTOR STABILITY, 1983: A REVIEW*

Conf-830776--2

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INTRODUCTION

When I was invited to write this review, my first thought was to make a thorough coverage of every paper that had been written on stability in the last few years. I had not yet done this when the manuscript began to grow unmanageably large. I therefore decided to restrict myself to subjects in which I am experienced and to discuss them from my own point of view. If this will not provide an encyclopaedic review, it will, I hope, provide one that is still entertaining and informative.

Stability (or instability) is a characteristic of *steady states*. If the system (e.g., a power grid, a chemical plant, or a superconducting magnet) returns to its original steady state after a perturbation, the steady state is said to be stable. If the system does not return, the steady state is said to be unstable. Sometimes, the steady state is stable against small perturbations and unstable against large ones, we call such behavior metastable. Superconductors exhibit stable, metastable, and unstable steady states according to the conditions of operation (e.g., ambient temperature, transport current, magnetic field, degree of cooling).

Stability interests us because it is indispensable for continuity of operation. Stekly (KA65, ST65, ST66) and Laverick (LA65) were the first to build truly stable magnets, magnets that

*Research sponsored by the Office of Fusion Energy, U.S. Department of Energy, under contract W-7405-eng-26 with the Union Carbide Corporation.

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would recover the superconducting state irrespective of the size of the thermal perturbation they suffered. This they did by immersing the magnets in pools of boiling helium and reducing the overall current density by adding copper until the Joule power per unit conductor surface in the normal state was less than the minimum film boiling heat flux. This kind of stability became known as cryostability because it was achieved by overwhelming the Joule power with cooling. Cryostability is attractive because it relieves us of the burden of knowing exactly what thermal perturbations the magnet will suffer. We can rely on continuous operation of the magnet, and our reliance is not diminished by our ignorance. This idea is so appealing that three of the six magnets of the Large Coil Program are cryostable pool boilers. The drawback of cryostability is the low current density it requires (≤ 3 kA/cm² at 8 T, 5 kA/cm² at 3 T for NbTi magnets), because low current density means large size and weight and high cost.

Stekly's way of cryostabilizing penalizes the current density more than is necessary, and Maddock, James, and Norris (MA69) showed that cryostability can be preserved at somewhat higher current densities than Stekly's criterion allows. In Stekly-stable magnets, all parts of a normal zone disappear simultaneously, whereas in Maddock-stable magnets, the ends of the normal zone propagate inward and the center is the last point to disappear (the so-called "cold-end recovery"). In spite of some improvement in current density, Maddock-stable magnets suffer from low current density, too.

Other attempts have been made to increase current density. Some workers, endeavoring to maintain cryostability, have tried to improve heat transfer between conductor and helium. These efforts have taken three different roads, namely, improvement of boiling heat transfer by treatment of the conductor surface, renunciation of pool boiling in favor of internally cooling the superconductors, and use of superfluid helium as coolant. The second and third roads are discussed below. Other workers reject cryostability, considering it too conservative and expensive, and propose to build metastable magnets. Their idea is to "make it strong, make it stout, make it out of things you know about"* and so keep the thermal perturbations well below the size that causes instability. This road to the prize of smaller, lighter, and cheaper magnets often leads to gratifying success, for example as in the case of ORNL's EBT-P magnet, which operates reliably at 7.5 T with a current density of 10 kA/cm².

*To the best of my knowledge, this phrase was composed by D. Bruce Montgomery of the Massachusetts Institute of Technology.

INTERNALLY COOLED SUPERCONDUCTORS

As mentioned above, three of the six magnets of the Large Coil Program are cryostable pool boilers; the other three magnets are wound with internally cooled superconductors. Each of the latter three is different from the others (see Fig. 1). As far as stability goes, the rationales of the three coils are as different as their designs. The designs of the Euratom and Swiss coils were "guided by the intention to reduce the possible heat release as much as possible by using a rigid conductor and a monolithic winding pack." (Y082) The designers of the Westinghouse coil, on the other hand, accepted the possibility of heat release due to strand motion but sought to offset it through "a very high cooled surface [that] can be obtained by subdividing the conductor in many individual strands [and] provides the ability to rapidly remove heat input by sudden energy release and/or Joule heating from the conductor." (Y082)

The idea of increasing the cooled surface by subdividing the superconductor into thin strands is an old one. The current spate of activity on internally cooled superconductors of the Westinghouse type (also called cable-in-conduit conductors) was in part touched off by Hoenig, Iwasa, and Montgomery's having reminded us in 1975 of the advantages of subdividing (H075). But the idea had already been suggested in 1967 by Chester (CH67), who wrote, "another important parameter in the stability condition is the thermal capacity of the superconductor . . . the superconductor may be combined with another material of higher thermal capacity . . . the most effective material would clearly be liquid helium or high-pressure helium if this could be retained in close thermal contact with the superconductor, perhaps by using a hollow tubular conductor with the helium trapped inside. Clearly, excellent thermal contact is desirable between the superconductor and the thermal ballast . . . this is achieved by subdivision of the superconductor to present greater interfacial area."

Chester understood clearly that internally cooled superconductors would not be cryostable, though he did not say so explicitly, for he put his discussion of them in the section of his paper dealing with adiabatic stability rather than in the section dealing with cryostability. The reason internally cooled superconductors are not cryostable is this. Recovery from a thermal perturbation takes only tens of milliseconds, while the residence time of the helium in the coil is measured in minutes; therefore the helium inventory available for recovery is limited. This means that a large enough perturbation can quench the conductor by raising the helium temperature above the current-sharing threshold. So, strictly speaking, internally cooled superconductors are not cryostable, but rather metastable.

The most interesting thing about a metastable system, of course, is how great a perturbation it takes to produce instability. The generally accepted standard for internally cooled superconductors is the so-called stability margin of Hoenig, Iwasa, and Montgomery (HO75), namely, the largest, sudden, uniform heat deposition from which recovery is still possible. It is usually expressed in millijoules per cubic centimeter (mJ/cm^3).

Cable-in-Conduit Conductors - Multiple Stability

The stability margins of cable-in-conduit superconductors like that in Fig. 1(a) have been studied experimentally at both MIT and ORNL. A review of this work up to the beginning of 1980, including detailed references, can be found in reference DR80. The upshot of the ORNL experiments is summarized in the three-dimensional sketch of Fig. 2. Plotted on the vertical scale is the stability margin ΔH , defined above, and plotted on the horizontal scales are the transport current I and the helium flow velocity v in the interstices. The interstitial helium flow is imposed externally, for example, by pumps. The fold in the surface is connected with occasional multivaluedness observed in the stability margin (LU80). Outside the shaded base area under the fold, the stability margin is single-valued: as we raise the energy of successive thermal perturbations we go from recovery to quench. The stability margin is the dividing line. But inside the shaded area under the fold, we experience the sequence recovery-quench-recovery-quench as we raise the energy of successive thermal perturbations. Let me emphasize again that Fig. 2 represents a summary of experimental observations (cf., e.g., Fig. 3).

According to a schematic theory that Lue, Miller and I proposed (LU80), this unusual situation is caused by the interplay of (1) an early stage of purely conductive heat transfer from the strands to the helium and (2) heating-induced flow transients that improve the convective heat transfer between the strands and the helium. The details of this theory, which applies only to the case of zero imposed flow, are too numerous for this short review, so regretfully I must refer the interested reader to the original reference.

Limiting Current - Scaling Rule

Multiple stability has a profound impact on conductor design. As you can see from the actual data of Fig. 3(a), the upper stability margin (upper recovery-quench boundary, sheet CAD in Fig. 2) can be much larger than the lower stability margin (lower recovery-quench boundary, sheet GEB in Fig. 2). *Because we do not know the perturbations to which the conductor will be exposed*, the margin of safety in the region of multiple stability must be taken as the lower stability margin. Perturbations larger than the lower stability margin *may* quench the conductor, whereas smaller

perturbations can *never* quench it. In the region of single stability to the right of line BF in Fig. 2, the margin of safety is the larger upper stability margin. So it is worthwhile to operate our conductor at currents less than that determined by the projection of line BF in Fig. 2. The current I_B at point B sets an absolute upper limit to the safe region of singly stable operation. We at ORNL call it the limiting current and have devoted a considerable effort to its practical determination.

The schematic theory of reference LU80, though too crude to allow calculation of the limiting current, does predict how it should scale with various conductor parameters (DR81). With this scaling rule and some experimental data we can estimate I_B for a wide range of conductors. The scaling rule says that

$$J_B \sim [f f_{\text{cond}} (1 - f_{\text{cond}})]^{1/2} \cdot (T_c - T_b)^{1/2} \rho^{-1/2} \cdot D^{-1} \cdot t^{-1/5} \cdot \ell^{2/5}, \quad (1)$$

where J_B is the limiting current density over the cable space, f is the fraction of copper in the metal composite, f_{cond} is the fraction of metal in the cable space, T_c is the critical temperature, T_b is the ambient helium temperature, ρ is the copper resistivity including magnetoresistance, D is the hydraulic diameter of the helium-filled interstices, t is the duration of the normalizing heat pulse, and ℓ is the length of the initial normal zone.

Several of these dependences have been checked. Lue and Miller (LU81) varied the second group, $(T_c - T_b)^{1/2} \rho^{-1/2}$, through roughly a factor of 2.5 by varying the background magnetic field and found that I_B followed the group quite closely. In another set of experiments (LU81a), they checked the dependence on t and ℓ . Here they found that $I_B \sim t^{-1/15} \ell^{2/15}$ rather than the scaling given in Eq. (1). The discrepancy in the scaling exponents, at first sight disheartening, may actually shed some light on the convective heat transfer promoted by the heating-induced flow. In the derivation of the scaling rule, it is necessary to relate the heat transfer coefficient to the velocity induced in the helium by heating. For want of a better alternative, I originally chose the Dittus-Boelter relation for this purpose: $Nu \sim (Re)^{4/5}$. Now when $Nu \sim (Re)^a$, the scaling exponents for t and ℓ are $-a/4$ and $a/2$, respectively. So when I saw the experiments of Lue and Miller, I proposed to them modifying the heat transfer correlation

to $Nu \sim (Re)^{4/15}$ to bring the theory into accord with the experiments. Finally, Lottin and Miller (LO82) checked the dependence on T_b in a pair of experiments conducted at 4.2 K and 3.6 K, respectively. The scaling rule predicts an increase of 12% in I_{lim} ; the measurements gave an increase of 9%.

Upper Stability Margin

Using the scaling rule of Eq. (1), we can design cable-in-conduit conductors to operate at currents less than the limiting current so as to avail ourselves of the upper stability margin. How large is the upper stability margin? We can estimate it as follows. When the largest tolerable heat pulse ΔH is added to the conductor, the recovering conductor and the helium in contact with it ultimately equilibrate at the current-sharing threshold temperature T_{cs} . During recovery, some Joule heat ΔH_J is produced, so the heat absorbed by the conductor and helium is the sum of the initial heat pulse ΔH and the Joule heat ΔH_J . If this heat were absorbed *isobarically*, we would then have $\Delta H + \Delta H_J = \Delta h$, where Δh is the enthalpy difference between T_b and T_{cs} of the matter in a cubic centimeter of cable space, which is very nearly the enthalpy difference of the helium there. In a well-designed cable-in-conduit conductor, ΔH_J is comparable with ΔH , so we expect ΔH to be comparable with and smaller than Δh . Experiments (LU81, LU81a) have shown that ΔH is indeed comparable with Δh , but rather than always being smaller than Δh , sometimes it is larger. This is because the absorption of the heat by the helium is not isobaric and the work done by the expanding fluid elements can exceed the isobaric work (LU80). Nevertheless, the enthalpy difference of the helium is a useful practical estimate of the upper stability margin.

How large shall we make the stability margin? If we knew the spectrum of thermal perturbations to which the magnet would be exposed, we would make the stability margin comfortably larger than the largest perturbation. But in the absence of such knowledge, the only rational plan is to make the stability margin as large as we can, within the constraints imposed on us from outside. These constraints are usually those of magnetic field and space: we are told to make a certain magnetic field and given a certain envelope in which to build our magnet. These requirements fix the maximum field at the winding and the current density over the cable space. This leaves four principal parameters in our control: the ambient helium temperature T_b , the strand diameter, and the two composition variables f_{cond} and f .

In my own work, I assumed that T_b and the strand diameter are also given and have concentrated on finding the values of f_{cond} and f that maximize the stability margin while avoiding the region of multiple stability (TS81; see also DR77). Such stability-optimized conductors usually have low copper-to-superconductor (Cu/SC) ratios (around 2:1, rather than the very high Cu/SC ratios, 10:1 or more, proposed for the conductors of cryostable pool boilers) and helium-filled volume fractions of 30 to 40%. But the precise values of these composition variables, as well as of the stability margin itself, depend on the prescribed current density.

Recently, Lue and Miller gained some information about the perturbation spectrum in magnets wound with cable-in-conduit conductors from experiments with a small NbTi solenoid (LU82). The solenoid, capable of making 8 T of self-field, reached short sample current with no training at all at charge rates up to 0.16 T/s; no degradation due to strand motion was apparent. The measured stability margin at zero imposed flow and 7.1 T was about 45 mJ/cm³ of metal. This value (representing, incidentally, the lower stability margin), though quite low, nonetheless exceeded the perturbations in the magnet and was enough to guarantee stable operation. It may well be that in following the strategy of maximizing the stability margin, which is correct when we know nothing about the perturbation spectrum, we overdesign the conductor without being aware of it.

Lower Stability Margin

How large is the lower stability margin? According to the schematic theory of reference LU80, the lower stability margin should be comparable with the heat transferred from metal to helium during the early stage of pure heat conduction in the helium. This stage comes to an end when the fluid elements next to the conductor cross the pseudocritical line and the conductor is blanketed with low-density helium (burnout). According to calculations of mine, the time t_{bo} to burnout as a function of the heat flux q for 5-atm helium starting at 4.2 K is given by $qt_{bo}^{1/2} = 26 \text{ mW} \cdot \text{cm}^{-2} \cdot \text{s}^{1/2}$.

When we measure the lower stability margin, we add in a time t a heat density ΔH , which is transferred to the helium nearly as fast as we add it. Just as we finish adding the heat, burnout occurs. (That is why ΔH is the largest perturbation from which recovery is possible.) Since $q = \Delta H D_w / 4t_{bo}$, where D_w is the wire diameter, for 1-mm wires we expect $\Delta H t_{bo}^{-1/2} = 1040 \text{ mJ} \cdot \text{cm}^{-3} \cdot \text{s}^{-1/2}$.

Lue and Miller (LU81a) have measured the lower stability margin of a triplet of 1-mm strands soldered around a central heat wire. For pulse times of 10, 50, and 100 ms, they found $\Delta H = 60, 100, \text{ and } 140 \text{ mJ}\cdot\text{cm}^{-3}$, respectively. Because only 5/6 of the strand surface is exposed in a triplet, my estimate of $\Delta H_{\text{bo}}^{1/2}$ is 6/5 of the value given above. So we calculate $\Delta H = 125, 279, \text{ and } 395 \text{ mJ}\cdot\text{cm}^{-3}$, respectively. The experimental values do not include in ΔH the Joule heat produced during the time the central heater is on. If we assume that the theoretical values are correct and ascribe the difference to Joule heating, we find that the conductor produced 87%, 48%, and 34%, respectively, of the Joule power it would have produced if fully normal. Thus the conductor appears to have been driven only into the current-sharing range by the heating; the higher the heating rate, the higher the conductor's temperature.

Lue and Miller (LU81a) also measured the lower stability margin for a pulse time of 1 ms and found $\Delta H = 75 \text{ mJ}\cdot\text{cm}^{-3}$ compared with only $40 \text{ mJ}\cdot\text{cm}^{-3}$ expected from theory. This value is larger than that measured for a pulse time of 10 ms. Their explanation is that for high heating rates, heating-induced flow improves heat transfer and allows more heat to be transferred before burnout. However, Lottin and Miller (LO82) performed a similar experiments using a 1-ms pulse time and found a lower stability margin of only about $7 \text{ mJ}\cdot\text{cm}^{-3}$, compared with a theoretical expectation of $32 \text{ mJ}\cdot\text{cm}^{-3}$. It is not clear to me why there is this difference between two similar experiments, and my feeling about this whole matter is that while the interpretation of the lower stability margin in terms of transient conduction is quite plausible, more work needs to be done to make it absolutely sure.

Quench Pressure

Knowledge of the stability margin helps us to design magnets that are not likely to quench. But if they do, we must be sure that they will not be damaged. One of the hazards that internally cooled superconductors face is the very high internal pressure they may experience during a quench. The general quench problem must be solved numerically because it involves compressible flow of a near-critical fluid in the presence of strong wall friction (see section entitled "Numerical Studies"). But in the worst case of an entire flow path going normal all at once, it has been possible by approximate analytic methods to get a simple formula for the peak quench pressure p_{max} (MI80):

$$p_{\text{max}} = 0.10(Q^2 L^3 / D)^{0.36} \quad (\text{mks units!}) , \quad (2)$$

where Q is the constant heating rate of the helium ($W \cdot m^{-3}$), L is the half-length of the flow path (m), and D is the hydraulic diameter of the helium-filled volume inside the conduit (m). This relation has been tested for maximum quench pressures between 4 and 200 atm (KR81, MI80, LU82a) and found to hold well enough for design purposes.

Thermal Expulsion

Structural damage from too high an internal pressure is not the only hazard when the coil quenches: rapid expulsion of the helium from the ends of the flow paths back into the refrigeration system may also cause damage. Again, in the worst case of an entire flow path going normal all at once, it has been possible to get a simple formula for the expulsion velocity v (DR81a):

$$v = 0.95c \left(\frac{\beta Q}{\rho c_p} \right)^{2/3} \left(\frac{Dt}{fc} \right)^{1/3} \quad (3)$$

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Here c is the sonic velocity, β the volume coefficient of thermal expansion, Q the volume heating rate of the helium, ρ the density, c_p the specific heat, D the hydraulic diameter, t the elapsed time, and f the Fanning friction factor [about three times the smooth-tube value (HO76, LU79)]. This relation is valid only for early times; its limitations are explained in reference DR81a. It, too, has been tested experimentally (LU82a) and within the limits of its validity found accurate enough for design purposes.

Numerical Studies

As mentioned above, the general quench problem must be solved numerically because it involves compressible flow of a near-critical fluid in the presence of strong wall friction. Several authors have attacked this problem, making similar assumptions and finding similar results (AR78, KR79, MA79, BE80, KR80, AG81, RI81, SH81, IB82, TU82). A good review of this work can be found in reference RI81. Most of the numerical computations were undertaken to substantiate the workability of a particular magnet design and so did not result in general and widely applicable results like Eqs. (2) and (3).

All authors noted strong flows induced in the helium by heating during the quench. These flows have a marked effect on heat transfer from the metal to the helium. For want of a better alternative, the heat transfer coefficient has been related to the local flow velocity with the Dittus-Boelter-Giarratano correlation, but as we have seen in the section "Limiting Current - Scaling Rule," at least one shred of information indicates that this is not entirely correct.

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The computer programs can also be used to study stability. So far, none of them has been able to model the phenomenon of multiple stability. The reason may have to do with the way they handle transient conduction in the helium. Most programs add to the Dittus-Boelter-Giarratano heat transfer coefficient a transient term $h \sim \sqrt{k\rho c_p/t}$, where k is the thermal conductivity of helium, ρ its density, c_p its specific heat, and t the elapsed time. This means the influence of transient conduction declines slowly. On the other hand, according to the schematic theory of reference LU80, transient conduction should decrease abruptly as the fluid elements next to the conductor cross the pseudocritical line (burnout; also see the section called "Lower Stability Margin"). If the reader consults section IV of reference LU80, he will see that it is the deep dip in Fig. 19 that is responsible for multiple stability. Describing transient conduction without taking account of burnout makes the dip less deep and may well be the reason that the programs fail to predict multiple stability.

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One interesting phenomenon noted by Ries in his computations and reported in his review (RI81) is this: when a very long normal zone spanning several turns recovers, it may do so by first breaking up into several shorter zones. The breaks come at places where the magnetic field is low. At these places, the Joule power is lower owing to the lower magnetoresistance, and the cooling capacity of the helium is higher owing to the higher current-sharing threshold temperature.

In general, the computer programs are time-consuming to operate. The program YAQHEL, used at KfK/Siemens, requires 9.6 s of real time to advance 1 ms of problem time. The computations reported in reference KR79, which followed a quench in the Euratom Large Coil Test conductor for 3 s, therefore took 8 hours. Such a program is ill adapted to broad parametric studies and must be simplified and speeded up. Some of the authors referenced above have already taken steps in this direction (MA79, KR80).

SUPERFLUID HELIUM (He-II)

Superfluid helium (also called He-II) is a low-temperature liquid phase with unusual properties that make it behave differently from ordinary helium (He-I) when used as a coolant for superconducting magnets. The difference can be understood on the basis of three things: the phase diagram, the Kapitza law of interfacial heat transport, and the Gorter-Mellink law of heat transport in bulk He-II.

Phase Diagram

Figure 4 shows part of the phase diagram for He-II. The line LL' separating the two liquid phases originates at a point L on the

saturation curve called the lambda point (2.18 K, 0.05 atm). The line LL', sometimes called the lambda line, is nearly vertical (slope = -0.01 K/atm).

Superfluid helium is made from liquid He-I by lowering the pressure and removing the latent heat of the vapor, thereby cooling the remaining liquid. Sometimes the low-pressure, saturated superfluid helium is used for experimental purposes, but more often it is used as a coolant in a heat exchanger to make He-II at higher pressures, usually 1 atm. A typical working temperature is 1.8 K.

We are concerned here with heat transfer from a solid surface to superfluid helium. If we are using saturated, low-pressure He-II as coolant, then a vapor film may exist between it and the heated surface. If we are using atmospheric He-II, the film may consist of vapor and liquid He-I. If the temperature of the heated surface is low enough, liquid He-II may contact it directly. These conclusions can be read right off the phase diagram of Fig. 4.

It may also happen that if the temperature of the heated surface is to the right of the lambda line LL', but low enough, liquid He-II is in direct contact with the surface. One would not expect this from the phase diagram, and the reason is to be found in the Kapitza interfacial resistance.

Kapitza Resistance

Kapitza observed that there is a temperature discontinuity at a liquid helium-solid interface related to the heat flux q being transported:

$$q = \alpha(T_s^4 - T_h^4) . \quad (4)$$

Here T_s is the solid temperature, T_h is the liquid temperature at the interface, and α is a constant characteristic of the surface. Schmidt (SC81) has measured α for a copper surface and found it to be $0.043 \text{ W} \cdot \text{cm}^{-2} \cdot \text{K}^{-4}$. Formula (4) shows, for example, that $T_s = 3.0 \text{ K}$, $T_h = 1.8 \text{ K}$, and $q = 3.0 \text{ W} \cdot \text{cm}^{-2}$ are mutually consistent, so in this case no film of He-I would form even though the solid temperature $T_s > T_\lambda = 2.18 \text{ K}$.

The experimental data on Kapitza heat conduction are scattered and confirm Eq. (4) only roughly. For the uses to which it is put in this paper, namely, rough estimation, Eq. (4) is accurate enough. Detailed discussions of Kapitza conduction can be found in references SC81a and SE81.

As in all solid-liquid heat transfer problems, conveyance of the heat from heated surface to sink involves not only interfacial heat transport but also heat transport through the bulk of the liquid. The interesting thing about He-II and "the main feature which makes He-II attractive as a coolant is its unique thermal property which allows heat to be carried over long distances within a very small temperature differences and *without convection or fluid displacement*" (LO79) [*italics mine*]. This rapid conduction is described by the Gorter-Mellink law.

Gorter-Mellink Law

The law of heat transport in superfluid helium is not Fourier's linear law, but rather a nonlinear law proposed by Gorter and Mellink:

$$q = -K(dT/dx)^{1/3}, \quad (5)$$

where q is the heat flux, dT/dx the temperature gradient, and K a thermal conductance parameter. Except very close to the λ temperature, $K \sim 10 \text{ W}\cdot\text{cm}^{-5/3}\cdot\text{K}^{-1/3}$ as long as $T > 1.8 \text{ K}$. Thus a temperature gradient of only 10^{-3} K/cm will support a heat flux of $1 \text{ W}\cdot\text{cm}^{-2}$ in He-II, compared with $2 \times 10^{-7} \text{ W}\cdot\text{cm}^{-2}$ in 1-atm He-I! But, as Lottin and Gerard (LO79) point out, and as we might expect, "as the heat flow increases, there is a sudden increase in thermal resistance due to the formation of a thin film of poorly conducting fluid at the solid surface."

We can estimate when this will happen, and it is worthwhile to do so. We consider a channel of length ℓ and assume that the sudden increase in thermal resistance has its onset when the temperature of the (1-atm) He-II at the interface reaches T_λ . Then a steady-state heat balance requires

$$q = \alpha(T_s^4 - T_\lambda^4) = K\left(\frac{T_\lambda - T_b}{\ell}\right)^{1/3}, \quad (6)$$

where T_b is the sink temperature at the end of the channel (heat-exchange surface or open bath of He-II). With this equation we obtain the results shown in Table 1 for the maximum heat flux q_m the channel can support without the appearance of a film of He-I and the corresponding temperature of the solid T_s .

Table 1. The Maximum Heat Flux q_m a Channel of Length l Can Support Without Appearance of a He-I Film at the Hot End. T_b is the Ambient He-II Temperature and T_s is the Temperature at the Solid Surface at the Hot End.

l (cm)	T_b			
	1.8 K		2.0 K	
	T_s (K)	q_m (W/cm ²)	T_s (K)	q_m (W/cm ²)
10	3.17	3.35	3.02	2.57
100	2.77	1.56	2.66	1.19
1000	2.5	0.72	2.44	0.55

Note that q_m depends on the length of the channel, a feature quite different from heat transfer in He-I. Note also that the solid surface temperatures are quite low, possibly less than the current-sharing threshold temperature. *Thus the usual situation following a normalizing thermal perturbation will be the appearance of a He-I film on the conductor surface.* Once the He-I film has appeared, the interfacial heat transfer coefficient is that of He-I, namely, about $0.1\text{--}0.3 \text{ W}\cdot\text{cm}^{-2}\cdot\text{K}^{-1}$.

Normal Zone in a Long Channel

Let us suppose now that the channel is partly filled with a superconductor in the center of which a normal zone of length L has developed. The total Joule power being produced in the normal zone is $\rho I^2 L / A_{\text{Cu}}$, where ρ is the resistivity of copper (including magnetoresistivity), I is the transport current, and A_{Cu} is the cross-sectional area of the copper. Can this heat be conducted away by the superfluid helium in the channel? The channel can conduct away at most $q_m A_{\text{He}}$ watts. Thus,

$$\rho J^2 L < f_{\text{He}} f_{\text{Cu}} q_m ; \quad (7)$$

$J = I/A$ is the overall current density (A is the area of the channel), and $f_{\text{He}} = A_{\text{He}}/A$ and $f_{\text{Cu}} = A_{\text{Cu}}/A$ are the respective volume fractions of helium and copper. The right-hand side is on the order of a few watts per square centimeter, as we have already seen. At 8 T,

$\rho \sim 5.5 \times 10^{-10} \Omega \cdot \text{m}$ (RRR = 100), and typically for a cryostable fusion magnet $J \sim 5000 \text{ A/cm}^2$. Thus $\rho J^2 = 1.4 \text{ W} \cdot \text{cm}^{-3}$, and L must then be on the order of a few centimeters. This is certainly not the case in large fusion magnets, where a turn may be several meters long. So in general we may expect that during recovery a superfluid-filled channel will not be able to carry away the Joule power as quickly as it is produced. If the channel is to clear away the heat by conduction after recovery, the helium temperature must remain below T_λ . Thus the enthalpy available for recovery is that of the superfluid helium between T_b and T_λ , about $300 \text{ mJ/cm}^3_{\text{He}}$ when $T_b = 1.8 \text{ K}$.

Transient Heat Transfer

The entire enthalpy difference just mentioned above is not available for recovery because the helium next to the heated surface reaches T_λ before the rest of the helium does (burnout). Claudet and Seyfert (CL82) have calculated the ratio of the heat transferred up to burnout ΔE to the available enthalpy ΔE_0 using the nonlinear diffusion equation that arises from the Gorter-Mellink relation (5). The boundary condition they chose was constant heat flux through the heated surface. Their results, together with their measurements, are shown in Fig. 5.

The good agreement of Claudet and Seyfert's calculations with their measurements strongly supports the Gorter-Mellink relation as a practical basis for describing heat transport in He-II in the range of heat fluxes of interest in building magnets ($\sim 1 \text{ W/cm}^2$). Van Sciver (SC79), interested in the details of this heat transport process, measured at various times t the temperature profiles (ΔT vs z) in a long channel arising from a constant heat flux through the heated surface. Using the Gorter-Mellink law, I found (DR82) a similarity solution in terms of the variables $\Delta T/\sqrt{t}$ and $z/\sqrt{t}$, use of which collapses all of van Sciver's profiles to a single curve. The calculated curve was in excellent agreement with van Sciver's experimental points, lending further support to use of the Gorter-Mellink law.

The constant-flux problem is one of a classical trilogy of one-dimensional diffusion problems, the other two being the constant-temperature problem and the pulsed-source problem. Lottin and van Sciver (LO82a) studied the latter problem by measuring the temperature profiles from a plane, pulsed source in a long channel. This problem, too, has a similarity solution, which can, moreover, be expressed in simple, algebraic terms (DR81b). The similarity variables $\Delta T \cdot t^{3/2}$ and $z/t^{3/2}$ reduce all of Lottin and van Sciver's

profiles to a single curve with which the calculated curve agrees extremely well (DR83). The similarity solution for this problem can be used to study the late stages of the disappearance of a heated zone in a long channel.

The last of the three problems, the constant-temperature problem in a long channel, also has a similarity solution that can be expressed in simple, algebraic terms (DR81b). Seyfert et al. (SE82) have argued that this problem describes heat transfer from a solid surface to He-II after burnout. In their words, "At the onset of burnout, formation of the thermal barrier starts. The He-II near the heated surface experiences a phase transition. A He-II-He-I interface appears *which has its temperature locked at T_λ* We assumed that this barrier had a negligible thickness and that it only affected heat transport in He-II by the condition of a constant temperature, i.e., $T = T_\lambda$, at the hot end of the channels in our test section."

[italics mine] Seyfert et al. drive their sample beyond burnout with a heat pulse E_1 and then see whether or not it cools below T_λ again while producing a steady (post-heating) flux q . According to the similarity solution (which represents the limit of large $qL^{1/3}$) the largest tolerable heat pulse ΔE_1 should scale inversely as q^3 . This expectation is borne out by the data of Seyfert et al. Furthermore, the value of $\Delta E_1 \cdot q^3$ calculated from the similarity solution is in fair agreement with the experimental value (DR83). This last problem provides the most realistic model for discussion of the stability of magnets cooled with He-II.

Recovery in a Long Channel

The recovery of a normal zone in a long channel will leave a bubble of warm helium behind. If the central temperature of the bubble does not exceed T_λ to start with, the disappearance of the bubble can be calculated using the nonlinear Gorter-Mellink diffusion equation. The late stages can be described by the similarity solution to the pulsed-source problem mentioned above, according to which the central temperature rise $T - T_b$ reaches the value $\Delta T \ll T_\lambda - T_b$ in a time t , given by

$$t = 0.21 \frac{\rho c_p}{K} [(T_\lambda - T_b)^2 L^2 / \Delta T]^{2/3}, \quad (8)$$

where ρ is the helium density, c_p is the specific heat of the helium, and L is the initial length of a bubble heated to T_λ (DR83). When $T_b = 1.8$ K, $\rho c_p = 0.41$ J·cm⁻³·K⁻¹ and $K = 10.4$ W·cm^{-1/3}, so for a

1-meter bubble, ΔT reaches $0.1(T_\lambda - T_b)$ in about 10 s. If the central temperature T_1 of the bubble exceeds T_λ , the situation is different, and the bubble disappears by inward propagation of its ends at a velocity

$$v = \frac{q_m}{\rho c_p (T_1 - T_\lambda)} \quad , \quad (9)$$

where q_m is the maximum heat flux in the channel and ρc_p is the volumetric heat capacity of He-I (DR83). If a bubble with $T_1 - T_\lambda = 2$ K is created in the center of a 100-m-long channel, as might occur in a He-II-filled cable-in-conduit conductor, then $q_m \sim 0.33 \text{ W}\cdot\text{cm}^{-2}$, $\rho c_p \sim 0.35 \text{ J}\cdot\text{cm}^{-3}\cdot\text{K}^{-1}$, and $v \sim 0.47 \text{ cm}\cdot\text{s}^{-1}$. So a 1-m-long bubble takes about 200 s to disappear, and a 10-m-long bubble takes more than half an hour.

This performance is actually not very different from that of cable-in-conduit conductors filled with supercritical He-I, where the residence time of the helium in the conduit may be several tens of minutes. Furthermore, both coolants offer good stability (stability margin $\sim$ hundreds of millijoules per cubic centimeter). Supercritical helium must be pumped, whereas He-II need not be. But as noted earlier, pumping is not necessary for stability, recovery being promoted by the strong flow induced in the supercritical helium by the perturbation itself. So only a little pumping is necessary to clear out the occasional heat produced during recovery, as well as the steady nuclear and ac loss heat produced in tokamak magnets. Since the dissipation varies roughly as the cube of the fluid velocity, the pumping power consumed may actually be very small. He-II will have higher refrigeration costs, probably by a factor of 3, so there is a trade-off that must be worked out carefully in choosing a coolant.

METASTABLE MAGNETS

As we have seen earlier, internally cooled superconductors are metastable because the local inventory of helium available for recovery is limited. Another class of metastable superconductors includes those that are pool-cooled (so the helium inventory available for recovery is unlimited) but operated at current densities beyond that prescribed by the Maddock criterion. If a long normal zone is created in such conductors, they quench, but if a short enough normal zone is created, they recover. So the natural question about such conductors is how great is the largest *local* heat pulse from which they can still recover. The energy of this pulse is called the quench energy.

Martinelli and Wipf (MA72) and Wilson and Iwasa (WI78) both noted the existence of a localized normal zone that is steady but unstable, which they called the minimum propagating zone (MPZ). (By steady we mean that if unperturbed, the zone will endure unchanged indefinitely. By unstable we mean that if perturbed however slightly, the zone will change drastically, never returning to its initial form. A stick balanced on its lower end is a good, homely example of a steady, unstable state.) Wilson and Iwasa, who wrote in 1978, "advance[d], but [were] unable to prove, the hypothesis that the energy required to create the MPZ is the minimum energy to quench the superconductor." In spite of the passage of five years, this hypothesis, so tantalizingly simple to state, has been neither proved nor disproved. If it is so, they say, then once we know the disturbance spectrum, we should be able to design by calculation conductors that will be stable against the disturbances. Unfortunately, that time has not come yet.

In the absence of such knowledge, the only rational plan is again to make stability as large as we can within the constraints imposed on us from outside. Using the MPZ energy as the measure of stability, Elrod et al. (EL81) determined in this way the best Cu/SC ratio in certain metastable conductors. They found that small ratios, $\sim 2-3$, led to higher MPZ energies than the larger ratios, say 10 or more, that are used in cryostable pool-cooled conductors. Similar trade-off studies were done by Wipf (WI78a).

Also using the MPZ energy as the measure of stability, I determined the maximum allowable resistive fault in a metastable conductor (DR82a). I observed that in the presence of a resistive fault, the heat transport equation has either two steady localized states or none (according to the resistance of the fault). When two states exist, the one with the smaller energy of formation is stable; the one with the larger energy of formation, the MPZ, is unstable. As the fault resistance increases, the larger energy decreases and the smaller energy increases. Following Wilson and Iwasa, I proposed using the energy difference between the upper, unstable state and the lower, stable state to get an idea of how stability (against a heat pulse at the fault, of course) declines as the size of the defect increases. Beyond the fault resistance at which the two energies become equal, no steady localized states exist, so this fault resistance is the maximum allowable.

Meuris and her coworkers (ME81, SE81a) have used a similar approach to study the influence of spacers (uncooled regions) on the stability of metastable conductors. Their situation was somewhat more complex than the preceding, there being altogether four steady solutions that occur in various combinations according to the parameters of the problem. Again, the quench energy is estimated as the energy of formation of the unstable steady state. One interesting result

of Meuris's work is that, under certain circumstances, a stable localized resistive zone is possible. Such zones have occasionally been observed (CL79).

Potted Magnets

Potted magnets are another kind of metastable magnet. The situation for these magnets was summed up at the Saclay Workshop by Scott, who said, "Measurements [of the quench energy] on this type of winding have not been reconciled with theoretical results." In his paper, Scott (SC81b) attempted such a reconciliation using the MPZ idea of Martinelli and Wipf and of Wilson and Iwasa.

Treating the winding pack as an anisotropic continuum, we find by dimensional analysis that the quench energy must have the form

$$E = \frac{S k_r \sqrt{k_z} (T_c - T_b)^{5/2}}{(\lambda_w \rho_c J^2)^{3/2}} F(i), \quad (10)$$

where S is the volumetric specific heat of the winding pack; k_r and k_z are its transverse and longitudinal thermal conductivities, respectively; T_c is the critical temperature; T_b is the ambient helium temperature; λ_w is the volume fraction of superconducting composite in the winding pack; ρ_c is the normal resistivity of the composite; J is the current density in the composite; i is the ratio of the transport to the critical current; and F is a function that cannot be determined by dimensional analysis.

Wilson (WI83) has estimated the function $F(i)$ in an ingenious manner. He wrote down steady solutions of the three-dimensional heat balance equation having fictitious cold outer boundaries, calculated their energies of formation, and chose the minimum formation energy (with respect to outer radius) as the quench energy. That his formula fits Scott's data is beyond doubt; why it works so well is still moot. To free the theory of Wilson's imaginative but questionable assumptions, I calculated the quench energy from the time-dependent solutions of the heat balance equation (DR83a). I found a formula like Eq. (10) with $F(i) = 6\pi^2(1-i)i^{3/2}$, which also fits Scott's data well.

If k_r approaches zero, so does the quench energy E given by Eq. (10). Now this cannot be correct for the following reason. When $k_r = 0$, the conductor can still recover from finite energy pulses small enough that their central temperatures fall, owing to one-dimensional conduction, below the threshold temperature for

current sharing before too much Joule heat is created. Considering only the superconducting composite, we find (DR81c) by dimensional analysis that the quench energy must have the form

$$E = \frac{SA(T_c - T_b)^{3/2}\sqrt{k}}{(\rho_c J^2)^{1/2}} G(i) , \quad (11)$$

where A is the cross-sectional area of the composite, G(i) is another function of i not determined by dimensional analysis, and all other symbols have the same meaning as before, except that S and k refer only to the composite. According to reference DR83a, $G(i) = \pi(1 - i)i^{1/2}$.

The functions F(i) and G(i) should take into account the temperature dependence of both the specific heat and the thermal conductivity, as well as any spatial and temporal extension of the heat source used in the experiments. Some hints of the effects of these phenomena are given in reference DR83a. My personal feeling is that work remains to be done on this problem.

CONCLUDING REMARKS

Three main topics have been discussed in this paper, namely, internally cooled superconductors, cooling by superfluid helium, and metastable magnets. The discussion of each has centered around a dominant idea, and it is fitting to highlight these ideas by way of conclusion. With regard to internally cooled superconductors, most of what we have learned in the last few years centers on the strong motion caused by the thermal expansion of helium. How naive were our early calculations that treated the helium as though it were incompressible! Our discussion of He-II was organized around the Gorter-Mellink relation and the solutions of the nonlinear diffusion equation it gives rise to. And our discussion of metastable magnets revolved around the fruitful concept of the MPZ. These three ideas are sturdy trunks that support much of the thought about superconductor stability that has flowered in the past several years.

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FIGURE LEGENDS

1. The three internally cooled superconductors of the Large Coil Program. The Westinghouse conductor is a 486-strand cable encapsulated in a square, stainless steel jacket filled with supercritical helium. The Swiss conductor consists of a solder-filled cable compacted around a central, helium-filled tube. The Euratom conductor is a flat 23-strand cable in a stainless steel, helium-filled box; the individual strands are large (2.35 mm $\times$ 3.1 mm) and well separated (0.9 mm). Redrawn from reference Y082.
2. Schematic representation of the stability margin as a function of imposed helium flow and transport current (not to scale). The fold in the stability surface is connected with the occasional multivalued nature of the stability margin. Points under the surface correspond to recovery, points above it correspond to a quench.
3. Experimental measurements of multivalued stability margins (a) from reference LU81 and (b) from reference LU80. If we slice the stability surface of Fig. 2 through the fold with a plane parallel to the ΔH -I plane, we get a Z-shaped curve like that of (a). If we slice through the fold with a plane parallel to the ΔH -v plane, we get a smooth curve lying above a pair of intersecting segments like those of (b).
4. The phase diagram of He⁴.
5. Ratio of heat transferred ΔE to available enthalpy ΔE_0 . Here q is the heat flux and L the length of the He-II-filled channel. The measured points and the results of a numerical calculation based on the Gorter-Mellink equation (dashed curve) are taken from reference AY81. The solid curve is the limit of large $qL^{1/3}$ given in analytic form by the similarity solution of reference DR82. The dashed-dotted curve is a graphical interpolation between the long-channel limit and the point (0,1).

Fig. 1. The three internally cooled superconductors of the Large Coil Program. The Westinghouse conductor is a 486-strand cable encapsulated in a square, stainless steel jacket filled with supercritical helium. The Swiss conductor consists of a solder-filled cable compacted around a central, helium-filled tube. The Euratom conductor is a flat 23-strand cable in a stainless steel, helium-filled box; the individual strands are large (2.35 mm x 3.1 mm) and well separated (0.9 mm). (Redrawn from reference Y082)

Fig. 2. Schematic representation of the stability margin as a function of imposed helium flow and transport current (not to scale). The fold in the stability surface is connected with the occasional multivalued nature of the stability margin. Points under the surface correspond to recovery, points above it correspond to a quench.

(a)

(b)

Fig. 3. Experimental measurements of multivalued stability margins. ~~Figure (a) is taken from reference LU81, figure (b) from reference LU80.~~ Figure (a) is taken from reference LU81, figure (b) from reference LU80. If we slice the stability surface of Fig. 2 through the fold with a plane parallel to the ΔH - I plane, we get a Z-shaped curve like that of (a). If we slice through the fold with a plane parallel to the ΔH - v plane, we get a smooth curve lying above a pair of intersecting segments like those of (b).

Fig 4 The phase diagram of He^3

Fig. 5. Ratio of heat transferred ΔE to available enthalpy ΔE_0 . Here q is the heat flux and L the length of the He-II-filled channel. The measured points and the results of a numerical calculation based on the Gorter-Mellink equation (dashed curve) are taken from reference AY81. The solid curve is the ~~long-channel~~ limit, given in analytic form by the similarity solution of reference DR82. The dashed-dotted curve is a graphical interpolation between the long-channel limit and the point (0,1).