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REVIEW ARTICLE

Point-contact spectroscopy of heavy-fermion systems

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Abstract. A review of heavy-fermion compounds investigated in their normal and superconducting states by point-contact spectroscopy is presented. The Joule heating appears to have a considerable influence on the conductivity in metallic point contacts in both the normal and superconducting states. Nonetheless, some spectroscopic information can be obtained from the point-contact measurements, and Josephson effects are observed for certain heavy-fermion superconductors. We review in detail very recent development in this field, including experiments using mechanically controllable break junctions covering the regime from metallic conductivity to truly vacuum tunnelling.

1. Introduction

The nature of the normal as well as the superconducting (SC) ground state for heavy-fermion (HF) compounds became subject to intensive investigation soon after the synthesis of the first HF compounds based on Ce, and interest intensified further after the discovery of the early HF superconductors; see, e.g., the review by Grewe and Steglich (1991) and references therein. A wide variety of ground states were established for HF systems, namely the non-magnetic Fermi-liquid ($CeCu_6$, $CeAl_3$), antiferromagnetic ($CeAl_2$, U_2Zn_{17}), and superconducting ($CeCu_2Si_2$, UBe_{13}) ones, as well as both types of ordering (UPt_3, URu_2Si_2, UPd_2Al_3, UNi_2Al_3). A crucial point in the understanding of the HF properties is the formation at low temperatures of a strongly correlated band of electrons that arises from the hybridization between localized f electrons and conducting electrons. This results in a huge magnification of the density of states (DOS) of the charge carriers near the Fermi level, and the appearance of quasiparticles with masses of several hundred times the free-electron mass. This point is confirmed by a very large linear coefficient of the electronic specific heat, enhanced Pauli-like susceptibility, clear-cut T^2-dependence of the specific resistivity at low temperatures, etc.

The normal-state properties of HF compounds are believed to play an important role in the formation of the SC state. Certainly, the fact that superconductivity exists in a material with such a high concentration of Ce or U local magnetic moments is very surprising. Moreover, the size of the specific heat jump at T_c, which is of comparable magnitude to the specific heat in the normal state, as well as the extremely steep slope of the upper critical field at T_c, indicate that Cooper pairs are formed by the strongly hybridized, itinerant band of heavy-mass quasiparticles. As might be expected in this case, the SC properties

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of HF compounds differ substantially from those of the usual BCS-type superconductors. Here, an exponential temperature dependence of the thermodynamic and transport properties connected with an isotropic gap in the DOS of the quasiparticles is replaced by a power-law one, indicating at least a strongly anisotropic gap, and perhaps even the presence of points or lines of zeros in the energy gap for some crystallographic directions. Hence, non-s-wave symmetries of the SC wave function, as well as non-electron–phonon mediating pairing mechanisms are required by the experimental observations. Nowadays, after nearly two decades of research, the main questions as regards the symmetry of the SC order parameter and, especially, the pairing mechanisms in HF superconductors still remain to be resolved, analogously to the same problem for high-T_c superconductors.

Point-contact (PC) spectroscopy studies non-linearities of the current–voltage (I–V) characteristics of metallic constrictions with characteristic size d smaller than the inelastic electron mean free path l_{in}. In this case a voltage eV applied to a PC defines the energy scale for scattering processes. PCs enable spectroscopy of all kinds of quasiparticle excitation in conductors to be carried out. Along with the traditional subject of investigation, namely electron–phonon interaction (Yanson 1983, Khotkevich and Yanson 1995), other types of interaction—electron–magnon, Kondo-scattering, crystal-field levels of the rare-earth ions, etc—can be studied by PC spectroscopy (Jansen *et al* 1980, Yanson and Shklyarevskii 1986, Duif *et al* 1989).

PC investigations can be a key to achieving a better understanding of both the low-temperature normal state and the SC ground state of HF compounds. In the latter case, they can offer a clearer view of the possibly unconventional type of superconductivity exhibited by HF compounds. PCs can be used for a direct determination of the spatial or directional dependence of the SC order parameter via the mechanism of Andreev reflection (Andreev 1964) at the S–N interface (Blonder *et al* 1982, van Son *et al* 1987) or by studying multiple Andreev reflection in S–c–S contacts (Oktavio *et al* 1983, Cuevas *et al* 1996) (here S stands for superconductor, c for constriction, and N for normal metal). The exploration of Josephson effects using S–c–S weak links can also yield additional information about the phase of the complex order parameter.

The present review concerns PC experiments on well known Ce and U HF compounds. The paper is organized as follows. Section 2 covers the regimes of the current flow in PCs between normal metals, and presents the main relations concerning PC resistance, and information which one can obtain by studying the non-linear I–V characteristics of PCs. Section 3 describes processes that influence the transport of electrons in superconducting PCs—mainly Andreev-reflection phenomena. Section 4 outlines current views on the origin and properties of the HF ground state. The experimental results are discussed in sections 5 and 6 for the normal ground state and the superconducting ground state, respectively. In section 6, the recent experiments using mechanically controllable break junctions are described. The last section summarizes both the progress and the failures of the investigations.

2. PC spectroscopy of the normal state

Depending on the relationship between the mean free path of the electrons l and the PC characteristic dimension d, various regimes of electron flow in PCs have been established. With l_{el}, l_{in} being the elastic and inelastic mean free paths, respectively, these regimes are called *ballistic* ($l_{el}, l_{in} \gg d$), *diffusive* ($l_{el} \ll d \ll \sqrt{l_{in} l_{el}}$), and *thermal* ($l_{el}, l_{in} \ll d$). There is also a *tunnel* regime, which corresponds to the small probability D of electrons passing through the constriction (Kulik 1992), which is intermediate between the regimes of metallic

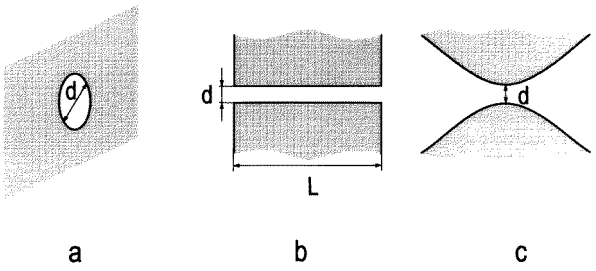

Figure 1. Models of point contacts: (a) a circular orifice with diameter d in an opaque screen; (b) a long channel with length $L \gg d$; and (c) a single-cavity hyperboloid of revolution.

constriction and pure tunnelling.

The most commonly used theoretical models present a PC as an orifice with diameter d in the non-transparent partition between two bulk metallic electrodes (figure 1(a)). Another extreme is a one-dimensional approach. That is the model of a long and narrow channel with the length $L \gg d$ (figure 1(b)). In some cases, a PC with the shape of a three-dimensional rotating hyperboloid is used (figure 1(c)). However, the shape of the PC does not influence the physics of the processes in the constriction radically, only altering some coefficients in the corresponding formulae. The current density j in the orifice with the normal parallel to the z-axis is given (see, e.g., Jansen *et al* 1980) by

$$j_z = 2e \sum_{\boldsymbol{k}} (\boldsymbol{v_k})_z f_{\boldsymbol{k}}(E).$$

Here $\boldsymbol{v_k}$ is the electron velocity, and $f_{\boldsymbol{k}}(E) = (\mathrm{e}^{(E_{\boldsymbol{k}}-E_\mathrm{F})/k_\mathrm{B}T} + 1)^{-1}$ is the Fermi–Dirac distribution function. In the ballistic and diffusive regimes, the Fermi surface splits in the PC into two parts, with the maximum energy difference defined by the bias energy eV (figure 2(a)) (Kulik *et al* 1977).

The summation over $\boldsymbol{k}$-space in the above formula can be replaced by integration over the energy E and over the surface of constant energy. For a voltage-biased ballistic PC,

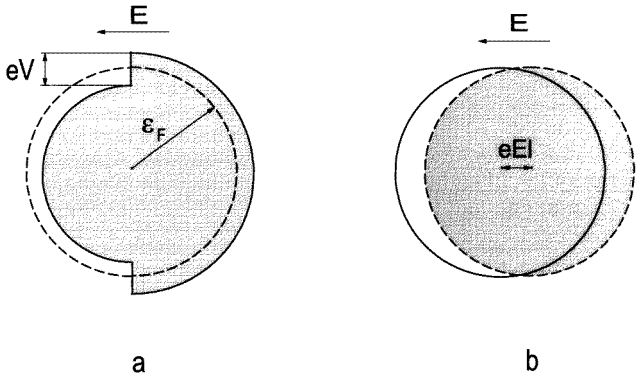

Figure 2. (a) The electron distribution function for the voltage-biased ballistic PC. (b) The same function for the thermal regime or for bulk material under the electric field $\boldsymbol{E}$ and with the mean free path l.

this gives

$$j_z = e\int_{E_F-eV/2}^{E_F+eV/2} \mathrm{d}E \int \frac{\mathrm{d}\Omega}{4\pi}\, f(E)v_z(E)N(E). \tag{1}$$

If we take into account the fact that $v(E)$ is inversely proportional to the density of electronic states $N(E)$, then we do not expect any non-linearity to be caused by the energy-dependent DOS. That is, Ohm's law is valid and the PC resistance, called the Sharvin resistance (Sharvin 1965), as follows from equation (1), is

$$R_{Sh} = \frac{16\rho l}{3\pi d^2} = \frac{16R_q}{(k_F d)^2} \tag{2}$$

where $R_q = G_q^{-1} = h/2e^2 = 12.9\ \mathrm{k\Omega}$, G_q is the conductance quantum, $\rho l = p_F/ne^2$, ρ is the specific resistivity, p_F is the Fermi momentum, n is the density of electrons, and e is the electron charge. From equation (2) one can determine the PC diameter d in the ballistic regime just by measuring its normal-state resistance at zero bias R_0. In the case of copper, $d \simeq 30/\sqrt{R_0\ (\Omega)}$ nm, which one can use for estimation of PC dimensions. It gives in general a reasonable value for most other simple metals, too.

If energy-dependent scattering processes in PCs are considered, then the absolute values of the energies of the initial and the final states are important. As a result, scattering into and out of the high-DOS region will be more probable. This can lead, in the case of an energy-dependent DOS and Fermi velocity, to structure in $\mathrm{d}V/\mathrm{d}I$ for the ballistic PC, which mimics the DOS image (Nowack and Klug 1992). Assuming this to occur, PC spectroscopy of the electronic spectrum $N(E)$ for metals is possible.

Electrons passing through the constriction gain the excess energy eV, and can scatter by creating quasiparticle excitations with this maximum energy. Electron–phonon interaction (EPI) in a ballistic PC results in backflow scattering, so some of the electrons are reflected by phonons and do not contribute to the current. This leads to a non-linear I–V characteristic, whose second derivative is expressed at $T = 0$, as shown by Kulik *et al* (1977), as (see also Kulik 1992)

$$\frac{1}{R(0)}\frac{\mathrm{d}R(V)}{\mathrm{d}V} = \frac{8ed}{3\hbar v_F}\langle K\rangle \alpha^2_{PC}F(eV) \tag{3}$$

where $\langle K\rangle$ is the geometric form factor averaged over the Fermi surface, taking into account the kinetic restriction of electron scattering processes in PCs. The K-factor depends on the scattering angle θ between the initial ($\boldsymbol{p}$) and final ($\boldsymbol{p}'$) electron momenta as $1 - \theta/\tan\theta$, and goes to infinity at $\theta = \pi$. However, this is an integrable singularity, which underlines the backscattering processes in the PC spectra.

The PC spectrum (usually measured as the second derivative of the I–V curve, i.e. $\mathrm{d}^2V/\mathrm{d}I^2(V) = R(0)\,\mathrm{d}R(V)/\mathrm{d}V$) as seen from equation (3), is proportional to the EPI spectral function $g_{PC}(eV) = \alpha^2_{PC}F(eV)$, while the voltage dependence of the differential resistance $\mathrm{d}V/\mathrm{d}I(V)$ mirrors the dependence of the scattering cross section on the energy. The function $g_{PC}(eV)$ is similar to the Eliashberg EPI function $g(\omega) = \alpha^2F(\omega)$ with $\alpha^2_{PC}(\omega)$ being the averaged EPI matrix element with kinetic restrictions imposed by the contact geometry and $F(\omega)$ the phonon density of states. $\langle K\rangle$ is equal to 1 in the ballistic regime and $\langle K\rangle \simeq l_{el}/d \ll 1$ in the diffusive regime, leading to a strong decrease of the PC spectral intensity for the diffusive regime compared to that for the ballistic one.

The case opposite to the ballistic and diffusive regimes is the thermal regime, where the PC resistance is given by the same Maxwell formula as for the diffusive regime:

$$R_M \simeq \frac{\rho}{d}. \tag{4}$$

This case arises if not only is $l_{el} \ll d$ but also $l_{in} \ll d$. Here, the excess electron energy is dissipated inside the constriction, which leads to Joule heating and an increase of the temperature T in the PC along with the voltage V. Assuming the fulfilment of the Wiedemann–Franz condition (Kittel 1971), the temperature of the PC, T_{PC}, is determined solely by the applied voltage, by virtue of the Kohlrausch equation:

$$V^2 = 8 \int_{T_{bath}}^{T_{PC}} \rho\lambda \, \mathrm{d}T.$$

Using $\rho\lambda = LT$, where L is the Lorenz number and λ is the thermal conductivity, we obtain a simple relation between the temperature at the centre of the PC, T_{PC}, and the applied voltage V:

$$T_{PC}^2 = T_{bath}^2 + \frac{V^2}{4L}. \tag{5}$$

For $T_{bath} \ll T_{PC}$ this yields a linear connection between T and V:

$$eV = (2\pi/\sqrt{3})k_B T_{PC} = 3.63 k_B T_{PC} \tag{6}$$

and so T_{PC} (K) $\simeq 3.2$ V (mV), using the standard Lorenz number $L = L_0 = 2.45 \times 10^{-8}$ V^2 K^{-2}.

In the thermal regime, the non-linearity of the I–V characteristic of a PC is caused by the temperature dependence of the resistivity $\rho(T)$. According to Kulik (1992),

$$I(V) = Vd \int_0^1 \frac{\mathrm{d}x}{\rho(T\sqrt{1-x^2}\,)|_{eV=3.63k_BT}} \tag{7}$$

which allows us to calculate an I–V characteristic using $\rho(T)$ or, conversely, to reconstruct the $\rho(T)$ dependence from a measured I–V curve of a PC.

If we suppose that the phonon part of the resistivity is small compared with the residual resistivity, then in the thermal regime the PC spectrum, i.e. the $\mathrm{d}^2V/\mathrm{d}I^2$ curve, of an Einstein oscillator with the energy $\hbar\omega_0$ looks like a smeared step with a shallow maximum at $eV = 1.09\hbar\omega_0$ (Kulik 1992). That is, in the thermal limit it is still possible to see degraded phonon features.

The thermal features in PC spectra can be quite pronounced if there is a phase transition at a particular temperature T which leads to a sharp change (i.e. a kink or hump) of the $\rho(T)$ dependence of a metal. Such a situation arises, for example, for ferromagnetic metals (Ni, Fe, Co) at the voltages corresponding to the Curie temperature T_C. In this case, a maximum in the $\mathrm{d}V/\mathrm{d}I$ curve and a N-type feature in the second derivative $\mathrm{d}^2V/\mathrm{d}I^2$ appear at about $eV_C \simeq 3.63 k_B T_C$ (Verkin *et al* 1979).

For an arbitrary relation between the mean free path and the PC diameter, a simple interpolation formula was derived by Wexler (1966) for the PC resistance:

$$R_{PC}(T) \simeq \frac{16\rho l}{3\pi d^2} + \frac{\rho(T)}{d}. \tag{8}$$

It turned out that Wexler's formula represents simply a sum of the ballistic ($l \gg d$) and diffusive ($l \ll d$) resistances. The predominance of one or another term in the PC resistance depends on the relation between d and ρ (or the mean free path l of the electrons). For high-resistivity materials and large contacts, the Maxwell contribution plays a dominant role and, conversely, the Sharvin resistance dominates for clean metals and small constrictions; see figure 3.

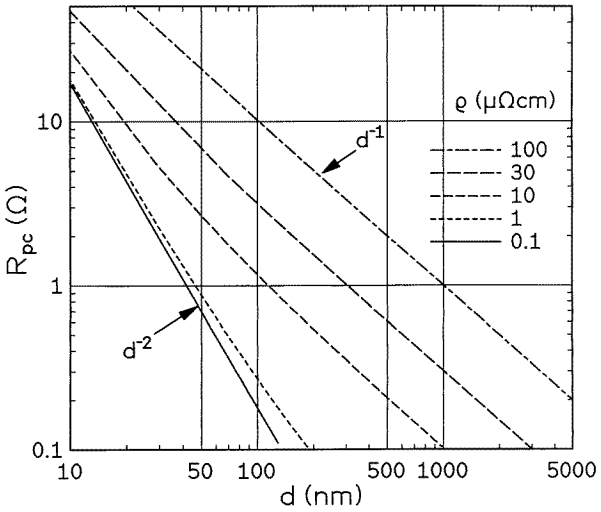

Figure 3. The PC resistance versus the diameter calculated according to equation (8), with $\rho = 0.1$, 1, 10, 30, and 100 $\mu\Omega$ cm and (standard for metals) $\rho l \simeq 10^{-11}$ Ω cm^2. It is seen that for typical PC resistances between 1 and 10 Ω, the diameter for a clean metal (the ballistic regime) is below 40 nm while those for high-resistivity materials (the diffusive regime), such as HF compounds, are between 100 and 1000 nm.

By analogy with the work of Akimenko *et al* (1982), one can determine the size (diameter) d of a PC by taking the derivative with respect to temperature of equation (8):

$$d \simeq \frac{\mathrm{d}\rho/\mathrm{d}T}{\mathrm{d}R_{\mathrm{PC}}/\mathrm{d}T}. \tag{9}$$

This formula is more acceptable than equation (8), since the residual resistance in a PC can differ from that of the bulk, ρ_0, and the decrease in transmission of the interface adds some undetermined coefficient to the Sharvin resistance. All of this limits the direct use of equation (8) for extracting d. It should be mentioned that before using relation (9) one has to be sure that the functional dependence $R_{\mathrm{PC}}(T)$ mirrors that of the bulk $\rho(T)$.

The second term of equation (8) describes diffusive electron flow in a PC, and its resistance is governed by ordinary electron transport as in the bulk. We can rewrite equation (8) as

$$R_{\mathrm{PC}}(V) \simeq R_{\mathrm{Sh}}\left(1 + \frac{3\pi d}{16 l(V)}\right) = R_{\mathrm{Sh}}\left(1 + \frac{3\pi d}{16 v_{\mathrm{F}} \tau(V)}\right). \tag{10}$$

It is seen that all of the processes resulting in the energy-dependent scattering time $\tau(eV = E) = l(E)/v_{\mathrm{F}}$ lead to an energy-dependent PC resistance. We recall that $\tau_{\mathrm{ep}}(E)$, determined by the electron–phonon interaction, is, according to Grimval (1981),

$$\tau_{\mathrm{ep}}^{-1}(E) = (2\pi/\hbar)\int_0^E \alpha_{\mathrm{tr}}^2(\omega) F(\omega)\ \mathrm{d}(\hbar\omega).$$

Then, taking the derivative of equation (10) yields

$$\frac{\mathrm{d}R_{\mathrm{PC}}}{\mathrm{d}V} = R_{\mathrm{Sh}} \frac{3\pi^2 e d}{8\hbar v_{\mathrm{F}}} \alpha_{\mathrm{tr}}^2(eV) F(eV)$$

which agrees nicely with the result from the truly microscopic calculation; see equation (3).

Considering other electron scattering processes, e.g. scattering on paramagnetic impurities (the Kondo effect), it is possible to use the Hamman (1967) expression for the temperature dependence of the resistivity to describe the voltage dependence of the PC resistance, by analogy with the work of Jansen *et al* (1980):

$$\frac{1}{R(0)}\frac{\mathrm{d}V}{\mathrm{d}I}(V) = 1 + \beta\left(1 - \frac{\ln(V/V_\mathrm{K})}{\sqrt{\ln^2(V/V_\mathrm{K}) + \pi^2 S(S+1)}}\right). \tag{11}$$

Here, the temperature is simply replaced by the voltage in the limit $eV \gg k_\mathrm{B}T$, V_K corresponds to the Kondo temperature T_K, S is the local moment of a magnetic atom, and β is the fitting constant. The Hamman formula works well for temperatures (energies) larger than the Kondo temperature (energy). For temperatures of the order of or smaller than T_K, a better approximation would be the empirical formula of Daybell (1973):

$$\frac{1}{R(0)}\frac{\mathrm{d}V}{\mathrm{d}I}(V) = 1 - \beta \ln\left[1 + (T/\theta)^2\right]_{k_\mathrm{B}T=eV} \tag{12}$$

where $\ln(\theta/T_\mathrm{K}) = -\pi[S(S+1)]^{1/2}$, S being the impurity spin, which also should be obtained empirically. This formula gives a Fermi-liquid $-T^2$-dependence for the Kondo resistance in the vicinity of the unitary limit $T \to 0$, unlike the Hamman expression which is exponential down to the lowest temperatures.

In the case of a heterocontact between two metals (1 and 2), the resistance is defined as a sum of the contributions from the two electrodes. Supposing a geometrically symmetric heterocontact and almost equal Fermi parameters of the metals, one can write

$$R_\mathrm{PC}(T) \simeq \frac{16\rho l}{3\pi d^2} + \frac{\rho_1(T) + \rho_2(T)}{2d}. \tag{13}$$

The PC spectrum in this case is a sum of the contributions of the spectra of the two metals (see equation (3)), with intensity inversely proportional to the Fermi velocity v_F. In other words, electrons moving more slowly through the PC have a higher probability of scattering in the PC by creating phonons. In the event of a large difference between the Fermi momenta $\boldsymbol{p}_\mathrm{F}$ of the electrons, the PC spectrum of the metal with the larger value of $\boldsymbol{p}_\mathrm{F}$ is transformed into narrow peaks (Shekhter and Kulik 1983), while strong elastic scattering in a PC hinders such a transformation.

For a heterocontact in the thermal regime, while the temperature of a PC increases compared to the bath temperature, the thermoelectric voltage, caused by the difference between the Seebeck coefficients of the contacted metals, results in an asymmetry of the I–V or $\mathrm{d}V/\mathrm{d}I(V)$ curves versus the bias voltage V. As was shown by Itskovich *et al* (1985),

$$\begin{aligned}\frac{1}{R(0)}\left(\frac{\mathrm{d}V}{\mathrm{d}I}\right)^{\mathrm{as}} &\equiv \frac{1}{2R(0)}\left(\frac{\mathrm{d}V}{\mathrm{d}I}(+|V|) - \frac{\mathrm{d}V}{\mathrm{d}I}(-|V|)\right) \\ &\simeq \left(\frac{\rho_1\rho_2}{L_0(\rho_1+\rho_2)^2}\right)^{1/2}(S_1(T) - S_2(T))\end{aligned} \tag{14}$$

where $S_{1,2}$ are Seebeck coefficients, and V and T are connected by the modified equation (5). Hence the voltage-dependent asymmetry of $\mathrm{d}V/\mathrm{d}I$ is defined by the difference between the thermopower of the two metals. In the case where metal 1 has high specific resistivity and thermopower, like HF materials ($\rho_1(T) \gg \rho_2(T)$, $S_1(T) \gg S_2(T)$), this gives

$$\mathrm{d}V/\mathrm{d}I(V) \propto \rho_1(T) \qquad (\mathrm{d}V/\mathrm{d}I)^{\mathrm{as}}(V) \propto S_1(T).$$

That is, the differential resistance of the PC behaves like the resistivity, and its asymmetric part like the Seebeck coefficient.

One should be cautious as regards identifying the thermal regime in a HF point contact with the true thermal equilibrium between conduction electrons, f electrons, and a lattice, since inelastic scattering occurs predominantly between conduction electrons and f electrons which are in thermal equilibrium and could have effective temperatures much higher than that of the lattice. That is, some non-linearities in the I–V characteristics might be due to the energy dependence of the 'bottle-neck' between the electron and phonon subsystems.

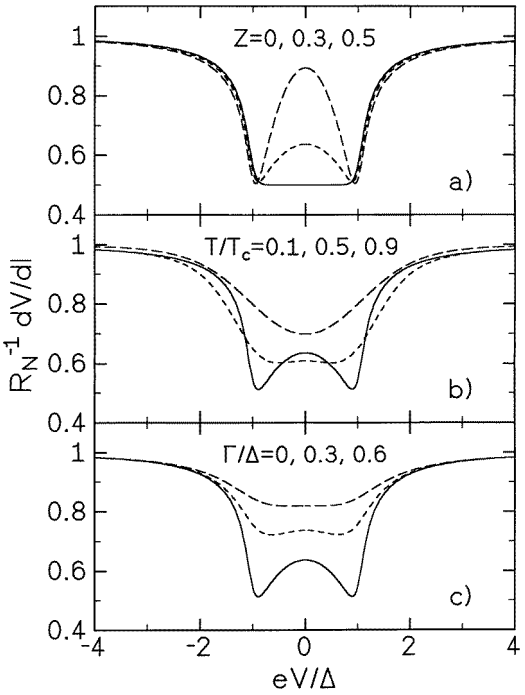

Figure 4. $\mathrm{d}V/\mathrm{d}I(V)$ curves calculated according to equations (15) and (16) (a) for different values of the barrier parameter Z and $\Gamma = 0$ at $T/T_c = 0.1$, (b) for $Z = 0.3$, $\Gamma = 0$, and different temperatures, and (c) for $Z = 0.3$ and different broadening parameters Γ at $T/T_c = 0.1$. (The fitting program used was developed by R Häussler.)

3. Point contacts with superconductors

Let us consider a S–c–N contact. Here the contribution from the Maxwell term in the superconducting electrode vanishes because $\rho_S = 0$, while the ballistic resistance can decrease to half its initial size as a result of Andreev reflection (AR) (Andreev 1964) at the N–S boundary. Blonder, Tinkham, and Klapwijk (BTK) (Blonder *et al* 1982), using a one-dimensional (figure 1(b)) scattering model for the contact with dimensions smaller than the SC coherence length ξ, have derived a simple formula for the I–V characteristic taking AR into account:

$$I(V) \sim \int_{-\infty}^{\infty} T(E)[f(E - eV) - f(E)]\,\mathrm{d}E \tag{15}$$

where

$$T(E) = \begin{cases} \dfrac{2\Delta^2}{E^2 + (\Delta^2 - E^2)(2Z^2+1)^2} & |E| < \Delta \\ \dfrac{2|E|}{|E| + \sqrt{E^2 - \Delta^2}(2Z^2+1)} & |E| > \Delta \end{cases}$$

where $f(E)$ is the Fermi distribution function, Δ is the SC order parameter, and Z is the barrier strength at the N–S interface. Equation (15) leads to a flat zero-bias minimum in $\mathrm{d}V/\mathrm{d}I(V)$ at $T \to 0$ and $Z = 0$, with the width 2Δ. Additionally, a maximum at $V = 0$ appears for $Z \neq 0$, so a double-minimum structure develops in $\mathrm{d}V/\mathrm{d}I(V)$; see figure 4(a). This feature is considered as a signature of an Andreev-reflection current. It smears out with increasing temperature, and vanishes above T_c; see figure 4(b).

In the modified BTK theory (Srikanth and Raychaudhuri 1992, Pleceník *et al* 1994), the broadening of the quasiparticle DOS in the superconductor is introduced following Dynes *et al* (1978):

$$N(E, \Gamma) = \mathrm{Re}\left\{ \frac{E - \mathrm{i}\Gamma}{\sqrt{(E - \mathrm{i}\Gamma)^2 - \Delta^2}} \right\}. \tag{16}$$

Here Γ is the broadening parameter which accounts for the finite quasiparticle lifetime τ; $\Gamma = \hbar/\tau$. The finiteness of Γ results in a broadening of the minimum in $\mathrm{d}V/\mathrm{d}I(V)$ and a decrease of its depth (figure 4(c)).

The I–V characteristic of a S–c–N contact at $eV \gg \Delta$ is

$$I(V) = \frac{V}{R_N} + I_{ex}$$

where I_{ex} is the so-called excess current, equal in the ballistic regime (Zaitsev 1980, 1984) to

$$I_{ex} = \frac{4\Delta}{3eR_N} \tag{17}$$

and R_N is the resistance in the normal state; see figure 5(a).

As was first pointed out by Artemenko *et al* (1979) for dirty contacts, I_{ex} comes from the AR. Its value is approximately half of that for the ballistic regime:

$$I_{ex} = \frac{1}{2}\left(\frac{\pi^2}{4} - 1\right)\frac{\Delta}{eR_N}. \tag{18}$$

The above-mentioned theoretical results are confirmed well by measurements of I–V and $\mathrm{d}V/\mathrm{d}I(V)$ curves for PCs with conventional superconductors, but they can, of course, be modified for HF superconductors. Recently, taking into account intrinsic lifetime effects of HF quasiparticles, Anders and Gloos (1997) demonstrated that the zero-bias minimum in $\mathrm{d}V/\mathrm{d}I(V)$ should have a V-like shape instead of the flat BTK bottom. Again, with decrease of the coupling between two electrodes, a more BTK-like shape of $\mathrm{d}V/\mathrm{d}I(V)$ (figure 4(c)), decreased in the amplitude and width, can be observed. Unfortunately, it is not clear whether it is possible to substitute in equation (15) the strong energy dependence for HF quasiparticles directly.

Let us consider the electron–phonon interaction in S–c–N contacts at an energy well above Δ (Khlus 1983). It leads to a non-linearity of the I–V characteristics identical to that for N–c–N contacts, so the $\mathrm{d}^2V/\mathrm{d}I^2$ curves represent a sum of the electron–phonon interaction functions according to equation (3).

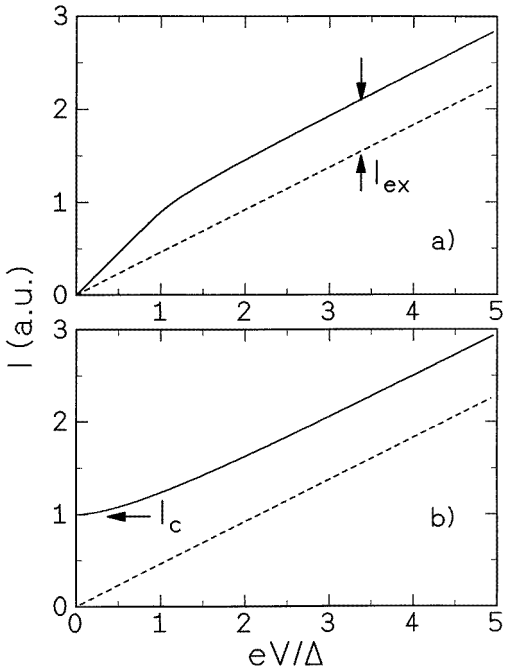

Figure 5. (a) I–V curves calculated according to equation (15) for $Z = 0$, $\Gamma = 0$, and $T/T_c = 0.1$ (solid line) and $T/T_c = 1$ (dashed line). (b) A sketch of the I–V curve of a Josephson contact in the RSJ model according to Licharev (1985) (solid line) and in the normal state (dashed line). The arrows indicate the definition of the excess, I_{ex}, and the critical, I_c, currents.

In a ballistic S–c–S contact, a subharmonic gap structure caused by multiple Andreev reflection appears in $dV/dI(V)$ for $eV \leqslant \Delta$, and becomes progressively more pronounced with decreasing transmission of the interface (Oktavio *et al* 1983, Cuevas *et al* 1996). For $eV \gg \Delta$ the excess current $I_{ex} \propto (\Delta_1 + \Delta_2)$ is present in the I–V curves. The d^2V/dI^2 characteristic of the S–c–S contact at an energy well above Δ should also reflect the electron–phonon interaction function (Khlus and Omel'yanchuk 1983).

The destruction of superconductivity by the critical current and the heating effects in S–c–S contacts were considered by Iwanyshyn and Smith (1972). They lead to a sharp minimum in $dV/dI(V)$ at $V = 0$ in an energy region of order 2Δ, while the excess current vanishes at larger biases. Thus, vanishing of I_{ex} with a bias testifies to the heating of the PC.

A characteristic feature of S–c–S contacts is the occurrence of a Josephson supercurrent (see figure 5(b)) with the following critical value I_c for a ballistic contact (Kulik and Omel'yanchuk 1977, 1978):

$$I_c = \frac{\pi \Delta_0}{eR_N}. \tag{19}$$

The value is lower by a factor of 2/3 for a dirty constriction (Kulik and Omel'yanchuk 1975). An unambiguous indication of the presence of a Josephson current is a Fraunhofer diffraction pattern of the critical current in a magnetic field:

$$I_c(H) = I_c(0) \left| \frac{\sin(\pi \Phi/\Phi_0)}{(\pi \Phi/\Phi_0)} \right| \tag{20}$$

where $\Phi_0 = 2.07 \times 10^{-15}$ T m^2, and Φ is the magnetic flux penetrating the contact. For Josephson contacts between conventional and unconventional superconductors, the critical-current oscillations in a magnetic field should have a minimum at zero magnetic field for some direction instead of a maximum as for conventional superconductors (see, e.g., van Harlingen 1995). This feature can be used to search for non-s-wave symmetry of the order parameter.

4. Heavy-fermion compounds

Heavy-fermion compounds containing certain lanthanide (mainly cerium) or actinide (mainly uranium) elements belong to the class of highly correlated electronic systems which display at low temperatures dramatically different properties as compared to the ordinary metals both in the normal and superconducting states. The crucial role here belongs to the partially filled f-electron shells of cerium or uranium ions, which carry magnetic moments. The large separation (about 3.5 Å) between neighbouring f atoms compared to the f-shell radius (about 0.7 Å) excludes a direct overlap of the f-electron wave functions; however, the correlation of electrons in the f shell favours the formation of a local magnetic moment. The localized f states can hybridize via exchange interaction with the conduction electron states and form a narrow band with the enormous electronic DOS $N(E)$, corresponding to quasiparticles with mass values two to three orders of magnitude higher than the free-electron mass.

The strength of the hybridization is responsible for the different mechanisms of formation of the delocalized f electrons (Steglich 1989, Grewe and Steglich 1991). Strong hybridization results in short-lived magnetic excitations or spin fluctuations in the itinerant f-electron system. Alternatively, for weak hybridization, local magnetic moments are formed at the f-atom sites, and demagnetization of these moments takes place through the 'cloud' of the conduction electrons which is known as Kondo compensation. This results in an effective enhancement of the band-electron mass at $T \ll T_K$, where

$$T_K \sim (E_F/k_B)\exp(-1/N(E_F)J)$$

is the so-called Kondo temperature, and J is the exchange integral of the localized spin and the conduction electron spin. The Kondo temperature demarcates the crossover from high-temperature local magnetic moment behaviour to the low-temperature non-magnetic state. This is easy to see even from simple transport measurements. At $T \gg T_K$ the logarithmic decrease of $\rho(T)$ with increasing T, characteristic of isolated Kondo impurities, is clearly evident for HF samples. On lowering the temperature, a transition from isolated Kondo impurities to many-body coherent dynamical screening of the local magnetic moments occurs. This results in Fermi-liquid behaviour, e.g., $\rho(T) = \rho_0 + AT^2$ at $T \to 0$ with a gigantic value of the A-coefficient. Since A is proportional to $(m^*)^2$, where m^* is the effective mass of the quasiparticles, this indicates narrow-band formation. Correspondingly, the Sommerfeld coefficient in the electronic specific heat γ and the Pauli spin susceptibility χ, both proportional to m^*, have huge values at low temperatures compared to those for room temperature. In addition to Kondo screening at low temperatures, an indirect interaction between localized magnetic moments caused by Ruderman–Kittel–Kasuya–Yosida (RKKY) spin-density oscillations is possible. This interaction can be characterized by the temperature $T_{RKKY} \sim (N(E_F)J)^2$, and it leads to antiferromagnetic ordering of the magnetic moments for $T < T_{RKKY}$. The competition between the RKKY and Kondo processes is governed by the exchange constant J. Following the approach of Doniach (1977), the crossover from the magnetic $T_{RKKY} \gg T_K$ state to the non-magnetic $T_{RKKY} \ll T_K$ ground state occurs at

the critical value J_c, below which a magnetically ordered state is formed; and above J_c, an effective suppression of the magnetic moments takes place before the ordering.

The superconducting state of HF systems is also dramatically different from that of ordinary superconductors. A well established power law of the temperature dependence of several transport properties, i.e. the specific heat and the magnetic field penetration depth in the SC state of HF compounds, gives strong evidence of zeros of the SC energy gap along lines or at points on the Fermi surface. Despite the low SC critical temperature of about 1 K, all HF superconductors have a large upper critical field H_{c2}, from a few to above ten teslas, along with a huge initial slope of H_{c2}. The magnitude of the initial slope is consistent with the large effective mass of the heavy quasiparticles, while the upper critical field for $T \to 0$ exceeds the Pauli limit calculated considering one Bohr magneton per quasiparticle. Moreover, the SC states of HF systems are very sensitive to small concentrations of impurities. For example, even a few per cent of non-magnetic impurities can totally suppress the superconductivity. All of the above-mentioned features point to a possible unconventional anisotropic order parameter $\Delta(\boldsymbol{k})$ for HF materials. The occurrence of the multiple superconducting phase transition, identified in UPt_3 and $(U_{1-x}Th_x)Be_{13}$, favours this assumption.

For more details of the HF phenomenon, the reader is referred to the reviews by Brandt

Table 1. Low-temperature properties of HF compounds investigated by means of PCs. The data are taken from the reviews by Fisk *et al* (1988) and Grewe and Steglich (1991), and the papers by Naidyuk *et al* (1988) (U_2PtC_2), Geibel *et al* (1991a, b) (UNi_2Al_3, UPd_2Al_3), Paschke *et al* (1996) ($CeCu_5Au$), Peysson *et al* (1986) (CeB_6), and Kittel (1971) (Al). The residual resistance of UBe_{13} was measured just above T_c. A linear coefficient A in the resistivity (instead of a quadratic one) is manifested for UBe_{13}, U_2PtC_2, and UNi_2Al_3. Note that there is considerable scattering of the data determined by different authors, related to the quality of the samples.

Compound	Structure	T_c (K)	T_N (K)	ρ_0 ($\mu\Omega$ cm)	A ($\mu\Omega$ cm K^{-2})	γ (mJ mol^{-1} K^{-2})	ξ (nm)	λ (nm)	m^*/m
Superconductive									
$CeCu_2Si_2$	Tetragonal	0.65	—	5–10	10	1000	9	450	380
UBe_{13}	Cubic	0.9	—	~ 100	36	1100	9.5	1000	260
U_2PtC_2	—	1.5	—	~ 10	4.5	75	—	—	—
UPt_3	Hexagonal	0.5	5	1	1	450	20	360	180
URu_2Si_2	Tetragonal	1.4	17.5	10	0.18	180	10	700	140
UNi_2Al_3	Hexagonal	1	4.6	6	1.6	120	—	—	70
UPd_2Al_3	Hexagonal	2	14	3.5	~ 0.2	150	8.5	400	65
Magnetic									
$CeAl_2$	Cubic	—	3.9	2.8	0.84	135	—	—	—
CeB_6	Cubic	—	2.4	1	—	260	—	—	—
$CeCu_2Ge_2$	Tetragonal	—	4.1	10	2.7	100	—	—	—
$CeCu_5Au$	Orthorhombic	—	2.3	30	—	650	—	—	—
U_2Zn_{17}	Rhombohedral	—	9.6	10	0.5	200	—	—	—
No ordering									
$CeAl_3$	Hexagonal	—	—	65	7.5	1600	—	—	—
$CeCu_6$	Orthorhombic	—	—	10	—	1300	—	—	—
Simple metal									
Al	Cubic	1.2	—	< 0.01	0	1	2000	30	1

and Moshchalkov (1984), Fisk *et al* (1988), and Grewe and Steglich (1991). In addition, we present in table 1 the characteristic properties of the HF compounds, investigated by means of PC.

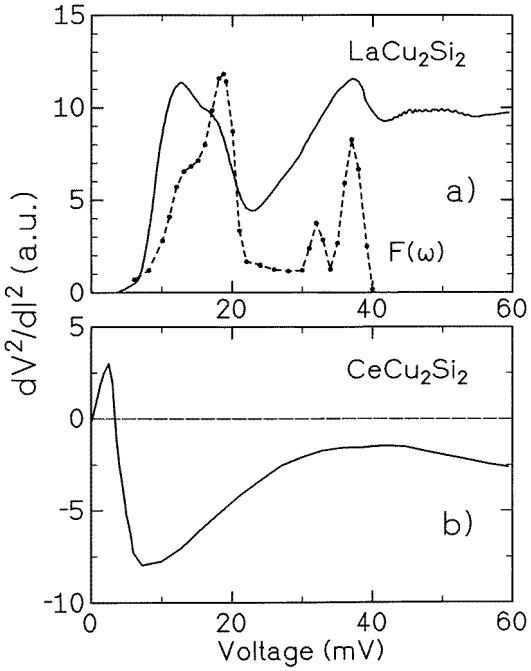

Figure 6. (a) $d^2V/dI^2(V)$ for a $LaCu_2Si_2$ contact ($R(0) = 13.5\ \Omega$) at $T = 4.2$ K (solid curve), and the phonon density of states $F(\omega)$ (dashed curve) at 8 K (Gompf *et al* 1987). (b) The $d^2V/dI^2(V)$ characteristic for a $CeCu_2Si_2$ contact ($R(0) = 30\ \Omega$) at $T = 4.2$ K. After Naidyuk *et al* (1985b).

5. PC study of heavy-fermion systems in the normal state

The first effort to investigate HF systems by means of PCs was undertaken by Wohlleben's group from Cologne, only a few years after the discovery of this method. Along with performing an investigation of magnetically ordered $GdCu_2Si_2$ and the intermediate-valence compound $YbCu_2Si_2$ (Leppin *et al* 1977, 1978), they mentioned in passing a check of the contacts between Mo and the heavy-fermion system $CeCu_2Si_2$. However, no results were presented. Only eight years later, in 1985, an intensive investigation of HF compounds by means of PCs began. Table 2 contains a list of the HF materials on which PC measurements have been performed.

The most dramatic result at that time was the statement by Moser *et al* (1986a, b) concerning the possibility of measuring by means of PCs the energy dependence of the electronic DOS (Moser *et al* 1985a, b). Here a minimum in $dV/dI(V)$ at zero bias was found for PCs on $CeAl_3$, $CeCu_6$, and UPt_3, with a width comparable to the expected narrow maximum of $N(E)$ at E_F for these materials. By analogy with the tunnelling approach (see Wolf 1985), the authors stated that the differential conductance dI/dV of a PC is proportional to the electronic DOS $N(E)$. It would be very tempting to measure the

Table 2. A list of heavy-fermion single crystals (sc) or polycrystals (pc) in the normal (N), superconducting (SC) and antiferromagnetic (AF) states that have been investigated by means of point contacts and mechanically controllable break junctions (MCBJs).

Compound	Sample	State	Counter-electrode	Reference
$CeAl_2$	sc	AF	Mo	Paulus and Voss (1985)
	sc	AF	Mo	Sato *et al* (1985)
	sc	AF	W	Gloos *et al* (1996a)
	sc	AF	$CeAl_2$ (MCBJ)	Gloos (1996)
$CeAl_3$	—	N	Mo	Moser *et al* (1985a, b)
CeB_6	sc	AF	Pt	Paulus and Voss (1985)
	sc	AF	Mo	Sato *et al* (1986)
	sc	AF	Mo, CeB_6	Kunii (1987a, b, 1992, 1996)
	sc	AF	Cu	Naidyuk *et al* (1991b)
$CeCu_6$	sc	N	Mo	Sato *et al* (1985, 1986)
	pc	N	Pt	Brück *et al* (1986)
	sc	N	Mo	Moser *et al* (1986b)
	sc	N	Cu	Naidyuk *et al* (1991b)
$CeCu_5Au$	sc	AF	Pt, $CeCu_5Au$	Paschke *et al* (1996)
$CeCu_2Ge_2$	sc	AF, SC	Cu, Mo, W	Gloos *et al* (1995a, b)
	sc	AF	W	Gloos *et al* (1996a)
$CeCu_2Si_2$	pc	SC	Nb	Han *et al* (1985)
	pc	N	Cu, $CeCu_2Si_2$	Naidyuk *et al* (1985a)
	sc	N	Pt	Paulus and Voss (1985)
	pc	SC	Al	Poppe (1985)
	sc	N	Mo	Sato *et al* (1986)
	sc	SC	Ag	de Wilde *et al* (1994)
	pc	SC	W, Cu, Ag	Gloos *et al* (1995b, 1996a, b)
	sc	SC	$CeCu_2Si_2$ (MCBJ)	Gloos *et al* (1998)
$Ce_xLa_{1-x}Cu_2Si_2$	pc	N	Cu, $Ce_xLa_{1-x}Cu_2Si_2$	Naidyuk *et al* (1985b)
UBe_{13}	pc	SC	Nb, Ta	Han *et al* (1985, 1986)
	pc	SC	W	Nowack *et al* (1987a, b)
	sc	SC	homo	Moreland *et al* (1994)
	pc	SC	W	Gloos *et al* (1996a, c)
	pc	SC	Cu, Pt, Ir, Ta, UBe_{13}	Kvitnitskaya *et al* (1996)
	pc	SC	UBe_{13} (MCBJ)	Gloos *et al* (1998)
$UCu_{4+x}Al_{8-x}$	pc	AF	Cu, Fe, Ni, $UCu_{4+x}Al_{8-x}$	Naidyuk *et al* (1993)
UIr_2Al_2	—	N	Mo	Solanki-Moser *et al* (1987)
UNi_2Al_3	pc	SC	W	Gloos *et al* (1996a, b)
	pc	SC	UNi_2Al_3 (MCBJ)	Gloos *et al* (1997, 1998)
UPd_2Al_3	pc	SC	Nb	He *et al* (1992)
	sc	AF	W	Volodin *et al* (1994)
	sc	SC	W	Gloos *et al* (1996a, b)
	sc	SC	UPd_2Al_3 (MCBJ)	Gloos *et al* (1998)

energy-dependent DOS in a such simple way by means of PCs, but the analysis carried out at the beginning of section 2 shows that the DOS does not enter directly into the current through the PC or its resistance.

Table 2. (Continued)

UPt_3	sc	SC	Al, Nb, UPt_3	Poppe (1985)
	sc	N	Mo, UPt_3	Moser *et al* (1986a, b)
	pc	SC	Pt	Nowack *et al* (1987a)
	pc	N	Cu, UPt_3	Lysykh *et al* (1988, 1989)
	sc	SC	Pt	Goll *et al* (1993, 1994, 1995a, b)
	pc	SC	Ag	de Wilde *et al* (1994, 1996)
	pc	SC	W	Gloos *et al* (1996a, b)
	sc	SC	Zn, Pb, NbTi	Naidyuk *et al* (1996b)
	pc	SC	UPt_3 (MCBJ)	Gloos *et al* (1998)
	sc	SC	Pt	Obermair *et al* (1998)
UPt_x	pc	N	UPt_x	Jansen *et al* (1987)
U_2PtC_2	pc	N	U_2PtC_2	Moser *et al* (1986b)
	pc	N, SC	Ag, Cu, W, Ta, Mo, U_2PtC_2	Naidyuk *et al* (1988)
URu_2Si_2	sc	AF, SC	Ag	Naidyuk *et al* (1991a)
	sc	AF, SC	Ag, URu_2Si_2	Nowack *et al* (1992)
	sc	AF, SC	Mo, W	Hasselbach *et al* (1992)
	sc	SC	Ag	de Wilde *et al* (1994)
	sc	AF	Au	Escudero *et al* (1994)
	sc	AF	Ag, Fe, Cu, URu_2Si_2	Naidyuk *et al* (1995)
	sc	SC	Zn, NbTi	Nowack *et al* (1995a, b)
	pc	SC	Ag	Samuely *et al* (1995)
	pc	SC	W	Gloos *et al* (1996a, b)
	sc	SC	Pt	Naidyuk *et al* (1996c)
	sc	AF	URu_2Si_2	Naidyuk *et al* (1996a)
	sc	SC	URu_2Si_2 (MCBJ)	Naidyuk *et al* (1997)
	sc	SC	Pt, URu_2Si_2	Rodrigo *et al* (1997)
	sc	SC	URu_2Si_2 (MCBJ)	Gloos *et al* (1998)
$URu_{2-x}Re_xSi_2$	pc	AF	W	Steiner *et al* (1996)
U_2Zn_{17}	—	AF	Cu, Zn, U_2Zn_{17}	Samuely *et al* (1987)
	pc	AF	W	Gloos *et al* (1996a)

A different approach in the interpretation of PC measurements on the HF compounds at that time was taken by Naidyuk *et al* (1985a) and Paulus and Voss (1985). They took into account the huge residual resistivity of HF materials at low temperatures, and examined a heating model.

Let us first present the PC measurements on $LaCu_2Si_2$ (Naidyuk *et al* 1985b), a compound that is isostructural to the well known HF system $CeCu_2Si_2$ (figure 6). The $d^2V/dI^2(V)$ curve shows a number of features, which coincide quite well with the phonon density of states, so the PC spectrum is believed to reflect the electron–phonon interaction function according to equation (3). $d^2V/dI^2(V)$ for $CeCu_2Si_2$ differs radically from that for $LaCu_2Si_2$. The lack of features, due to the electron–phonon interaction as well as a weak dependence of the PC spectra on the quality of the contacts, attests to the regime being non-ballistic in this case. In fact, the behaviour of the $dV/dI(V)$ curve is similar to the $\rho(T)$ dependence of the bulk material (see figure 7), as is expected for the thermal regime. This is clearly visible also for the other HF compounds represented in figure 7. Moreover, the asymmetric part of $dV/dI(V)$ is in qualitative agreement with the temperature dependence of the Seebeck coefficient, $S(T)$, in accord with equation (14); see figure 8.

An excellent agreement between the $dV/dI(V)$ and $dV/dI(V = 0, T) \propto \rho(T)$ curves

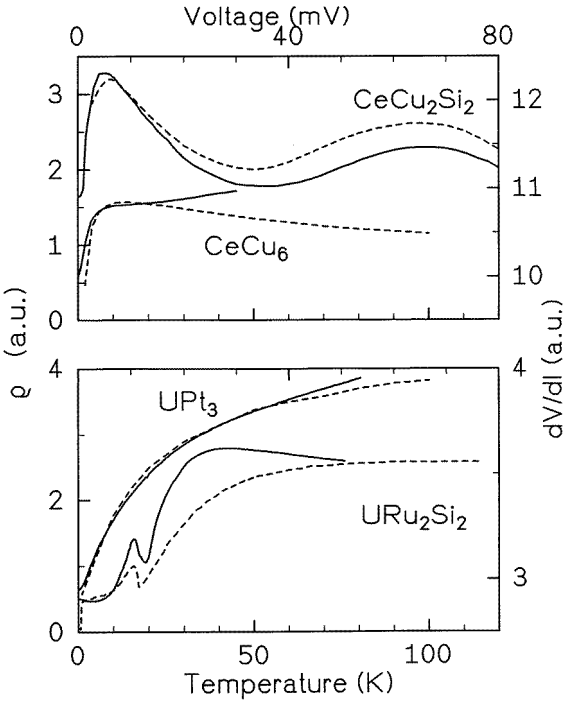

Figure 7. $\mathrm{d}V/\mathrm{d}I(V)$ for a PC (solid curves), and the specific resistivity $\rho(T)$ (dashed curves) for $CeCu_2Si_2$, $CeCu_6$, UPt_3, and URu_2Si_2. The $\rho(T)$ curves are from: Stewart (1984) ($CeCu_6$ and UPt_3), Brandt and Moshchalkov (1984) ($CeCu_2Si_2$), and Mydosh (1987) (URu_2Si_2).

was found for UPt_3 homocontacts (Lysykh *et al* 1988), where the adjustment between the V- and T-scales has been carried out according to equation (5), taking into account the temperature-dependent Lorenz number.

The next step towards proving the thermal model is that of computing the I–V characteristics according to equation (7). This was done for the UPt_x ($x = 1, 2, 3, 5$) compounds by Jansen *et al* (1987), and for $UCu_{4+x}Al_{8-x}$ by Naidyuk *et al* (1993). The calculations were found to describe the experimentally measured curves well, which further confirms both the theoretical model and that a thermal regime occurs, for these high-resistivity compounds.

As is often encountered for HF point contacts in the thermal limit, the relative change in $\mathrm{d}V/\mathrm{d}I(V, T = \text{constant})$ with increase of the voltage is smaller than the change in $\rho(T)$ versus T or that calculated from equation (7), especially with increasing PC resistance. This observation could be explained by including an additional constant contribution to the PC resistance caused by the following factors:

(i) increased residual resistivity of the metal inside the constriction (Naidyuk *et al* 1993);
(ii) a spoiled, poorly conducting interface layer (Nowack *et al* 1992, Gloos *et al* 1998);
(iii) a low transparent barrier at the boundary (Naidyuk *et al* 1991b); and
(iv) strong reflection of the electrons by mismatch of the effective masses or Fermi momenta of the contacting metals (Gloos *et al* 1995a).

Magnetic field has a strong influence on the transport properties of many HF compounds, leading to a substantial decrease of their resistivity, e.g. for $CeCu_6$ (Onuki *et al* 1985) and,

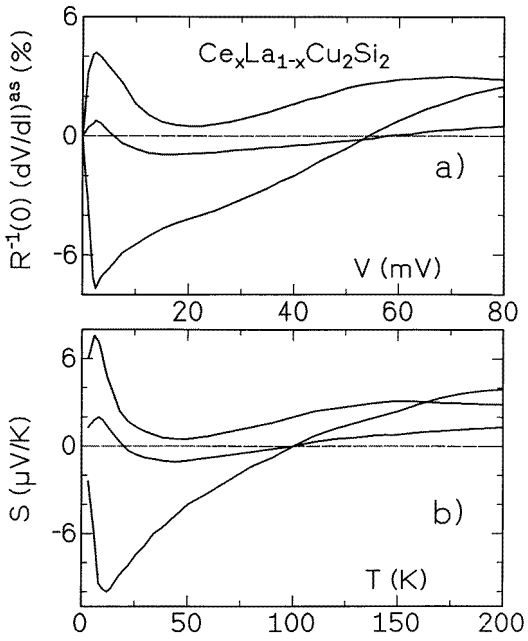

Figure 8. (a) The asymmetric part of $\mathrm{d}V/\mathrm{d}I(V)$ for a heterocontact between $Ce_xLa_{1-x}Cu_2Si_2$ and Cu (from top to bottom, $x = 0.3$, 0.7, and 1). After Naidyuk *et al* (1985a). (b) The thermopower versus T for the same compounds. After Brandt and Moshchalkov (1984).

especially, for CeB_6 (Peysson *et al* 1986). One would then expect a reduction of the inelastic scattering and an overcoming of the heating effects in strong magnetic fields. Nevertheless, the $\mathrm{d}V/\mathrm{d}I(V)$ curves for these compounds in high magnetic fields behave similarly to those for the bulk resistivity (Naidyuk 1991b), which again supports the assertion that the measured non-linearities have resistive natures. While, Kunii *et al* (1987a, 1992, 1996) drew attention to the remarkable difference between $\mathrm{d}V/\mathrm{d}I(V)$ and the bulk specific resistivity for CeB_6, they did not measure the temperature dependence of the PC resistance to ensure that the latter behaves for these contacts as in the bulk.

To our knowledge, there is only one example of a PC on HF compounds for which the behaviour of $\mathrm{d}V/\mathrm{d}I(V)$ really contradicts (even qualitatively) the $\rho(T)$ dependence (Naidyuk *et al* 1995). $\mathrm{d}V/\mathrm{d}I(V)$ for a URu_2Si_2 PC shows a maximum at $V = 0$; that is, $\mathrm{d}V/\mathrm{d}I$ decreases with the voltage around zero bias (figure 9(b)), whereas the zero-bias differential resistance $\mathrm{d}V/\mathrm{d}T(V = 0, T)$ increases with temperature. It also has a maximum at T_N that corresponds well to $\rho(T)$ (figure 9(a)). This correspondence is very important, since it testifies to the consistency of the transport properties of a PC with those of the bulk. Consequently, the PC presented is not in the thermal limit, at least at small biases, and $\mathrm{d}V/\mathrm{d}I(V)$ reflects some spectroscopic features. A gap opening on the Fermi surface in the AF state below the Néel temperature T_N due to spin-density-wave formation was considered as a possible reason (Naidyuk *et al* 1995, 1996a). The Sharvin resistance (equation (2)) can be written as $R_{Sh} \propto p_F/ne^2 \propto 1/S_F$, and should be sensitive to the destruction of a part of the Fermi surface S_F, because of the formation of a gap Δ_{AF}. Hence R_{Sh} increases at voltages $eV < \Delta_{AF}$ due to the removal of part of the Fermi surface, which leads to a maximum at zero bias. Note that this maximum was never observed by partially replacing Ru by Re in the $URu_{2-x}Re_xSi_2$ compounds in paramagnetic or ferromagnetic states (Steiner

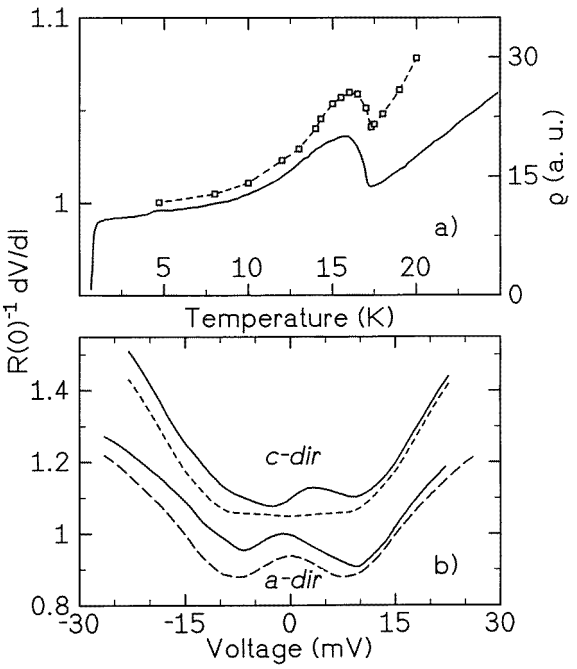

Figure 9. (a) The temperature dependence of the normalized PC resistance at $V = 0$ of a URu_2Si_2 homocontact ($R(0) = 3.2\ \Omega$) in the *ab*-plane (symbols) compared with the bulk resistivity according to Mydosh (1987) (solid curve). (b) Normalized $dV/dI(V)$ curves (solid curves) and their symmetric parts (dashed curves) for PCs established along the *c*-direction (the upper two curves) and in the *ab*-plane (the bottom two curves). The PC resistances are 3.3 and 3.2 Ω, respectively; $T = 4.2$ K. After Naidyuk *et al* (1995).

et al 1996) in the absence of spin-density-wave formation.

The most intriguing observation by Naidyuk *et al* (1995, 1996a) is that the zero-bias maximum in the symmetric part

$$(dV/dI)^{sym} = [dV/dI(+|V|) + dV/dI(-|V|)]/2$$

is well pronounced for PCs in the *ab*-plane, but it is usually suppressed or absent for PCs along the *c*-direction, whereas the asymmetric maximum is present for both curves (figure 9(b)). This observation is consistent with the STM experiments of Aliev *et al* (1991) and Aarts *et al* (1994), showing that the gap at the Fermi surface in URu_2Si_2 is strongly anisotropic and opens in the *ab*-plane.

Figure 10 shows the behaviour of the dV/dI maximum (i.e. related to the gap) at zero bias, extracted from the experimental curves at different temperatures, as well as the width of the maximum for three PCs. The value of the gap estimated from the width of the zero-bias maxima in the symmetric part of dV/dI is of the order of 10 meV, which corresponds well to the gap values obtained from the tunnelling, resistivity, specific heat, and neutron scattering measurements cited by Aarts *et al* (1994). The zero-bias maximum decreases rapidly with increasing temperature, while its width remains practically constant up to temperatures close to $T_N = 17.5$ K. This behaviour (figure 10(b)) looks rather like a first-order transition near T_N and deviates from that expected from the mean-field theory. Recall that in the weak-coupling limit the thermodynamics of the transition to the spin-density-wave state is described in the same way as that for a condensate in the BCS superconducting

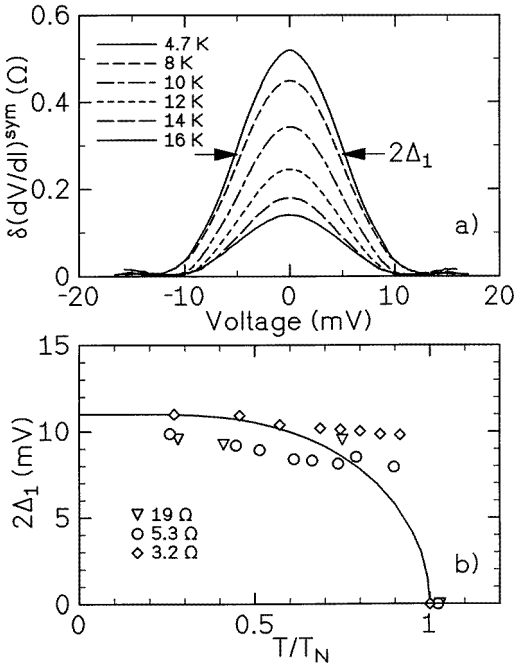

Figure 10. (a) The temperature behaviour of the zero-bias maximum $\delta(\mathrm{d}V/\mathrm{d}I)^{\mathrm{sym}}$ of the symmetric part of $\mathrm{d}V/\mathrm{d}I(V)$ for a URu_2Si_2 homocontact established in the *ab*-plane. (b) The temperature dependence of the width of the zero-bias maximum for three PCs (different symbols) established in the *ab*-plane. The solid curve indicates a mean-field temperature dependence. After Naidyuk *et al* (1995, 1996a).

state. In both cases a gap develops in the single-particle excitation spectrum with the well known BCS temperature dependence.

Elsewhere, the distance between the minima in the $\mathrm{d}V/\mathrm{d}I$ characteristics was taken as a gap width from a study of URu_2Si_2 PCs by Escudero *et al* (1994). They explain the $\mathrm{d}V/\mathrm{d}I(V)$ characteristics—that is, the whole structure with a double minimum near zero bias—by Andreev-like scattering of electrons on the spin-density-wave condensate in the AF state below T_N, in analogy with the case of a superconductor; see section 3. However, for $T > T_N$ only the $\mathrm{d}V/\mathrm{d}I$ maximum at zero bias disappears, while the overall $\mathrm{d}V/\mathrm{d}I$ minimum still does not vanish. Hence, the latter has no relation to the AF state and to the spin-density-wave gap. More probably, the double minimum arises due to a combination of the zero-bias maximum and a Maxwell contribution to the PC resistance increasing with voltage (temperature) like $\rho(T)$ (see equation (8)).

Alternatively, Hasselbach *et al* (1992) have proposed the change in the scattering amplitude of hybridized heavy quasiparticles, as they undergo the magnetic transition, as a possible mechanism for the formation of the 'gap-like' feature in $\mathrm{d}V/\mathrm{d}I(V)$. Unfortunately, the mechanism is explained there without details. On the other hand, Rodrigo *et al* (1997) analysed the $\mathrm{d}I/\mathrm{d}V(V)$ curves of URu_2Si_2–Pt contacts or homocontacts in terms of conductance spectra, considering the effect of resonance in the DOS at the Fermi level, as is appropriate to the Kondo system. They argued that the results are consistent with the existence of a Kondo-like bound state of the U^{4+}-ionic configuration and conduction electrons. However, this interpretation is questionable, since, in spite of the STM technique

used, the authors deal with true metallic PCs. Their resistances were between 15 and 100 Ω, and their $\mathrm{d}I/\mathrm{d}V(V)$ curves were similar to the $(\mathrm{d}V/\mathrm{d}I(V))^{-1}$ curves from the above-discussed papers, e.g. those of Nowack *et al* (1992), Hasselbach *et al* (1992), and Escudero *et al* (1994) (see table 2). In this case, for a metallic conducting PC, as we mentioned at the beginning of section 2, a direct observation of the DOS is not possible.

A pronounced feature in the $\mathrm{d}V/\mathrm{d}I(V)$ characteristics of heterocontacts with HF compounds is an asymmetry, which is mainly due to the large thermopower of the HF materials. In the case of URu_2Si_2 PCs, the asymmetric part of $\mathrm{d}V/\mathrm{d}I(V)$ has the same sign as $S(T)$; however, as mentioned by Naidyuk *et al* (1995), its temperature dependence contradicts the heating model at low biases. This is an additional confirmation of the spectroscopic origin of the above-discussed zero-bias 'gap' structure. As regards the possible non-thermal nature of the asymmetry in this case, in the framework of the ballistic PC theory of Nowack and Klug (1992), the $\mathrm{d}V/\mathrm{d}I$ characteristics might still image the electronic DOS, if energy-dependent backscattering processes are considered. In this case, the $\mathrm{d}V/\mathrm{d}I$ asymmetry in a ballistic heterocontact between a HF system and a simple metal could be due to the asymmetric peak in $N(E)$ at E_F for HF compounds.

6. PC study of heavy-fermion systems in the superconducting state

Direct information about the SC order parameter is accessible via PCs by means of study of the Andreev reflection at a normal-metal–superconductor (N–S) interface and/or by means of the study of Josephson effects as shown in section 3. A short but comprehensive review of the probing of the energy gap in HF superconductors, focused on PC and tunnelling spectroscopy, was given recently by von Löhneysen (1996). Here, more emphasis is placed on the recent progress in the field of PC measurements on HF superconductors, and all available papers will be mentioned. The most recent papers will be reviewed more thoroughly, especially those concerning experiments based on mechanically controllable break junctions.

6.1. S–c–N contacts

The first PC measurements on HF superconductors, by Nowack *et al* (1987a), showed a pronounced minimum at $V = 0$ below T_c with the width roughly corresponding to the expected SC gap value. The magnitude of the minimum was up to 50% of R_N for UBe_{13}, in accord with the BTK theory (Blonder *et al* 1982). In contrast, $\mathrm{d}V/\mathrm{d}I(V)$ for PCs on UPt_3 showed much smaller non-linearities, which constitute only a few per cent. The ratio $2\Delta(0)/k_\mathrm{B}T_\mathrm{c}$ was estimated to be close to the BCS value; see table 3.

Weak SC features in $\mathrm{d}V/\mathrm{d}I$ for UPt_3 contacts were found in further work by Goll *et al* (1993), de Wilde *et al* (1994), and Naidyuk *et al* (1996b). Obermair *et al* (1998) reported that the weak SC signal for UPt_3, presumably caused by the suppression of the order parameter at the surface, cannot be enhanced even by different surface treatments. An anisotropic structure in $\mathrm{d}V/\mathrm{d}I(V)$ for UPt_3 below T_c has been found by Goll *et al* (1993); see figure 11. While for the *c*-direction of the hexagonal UPt_3 crystal, a pronounced minimum (quite often even a double minimum) in $\mathrm{d}V/\mathrm{d}I$ was observed, the absence of such features for the perpendicular direction was the rule. Thus, the authors suggest (see also the paper by von Löhneysen and Goll 1994) an anisotropic order parameter for UPt_3; that is, the gap vanishes in the basal plane as predicted by various theoretical scenarios of unconventional superconductivity in UPt_3 (Sigrist and Rice 1992). Furthermore, as mentioned by Goll *et al* (1995a, b), $\mathrm{d}V/\mathrm{d}I(V)$ cannot be fitted well by the BTK model

Table 3. Gap values of HF superconductors as determined from the $dV/dI(V)$ curves of PCs using the BTK or modified BTK fit, with both isotropic and anisotropic order parameters, from the full width (FW) of the minimum at zero bias, from the peak structure (PS), or simply from estimation. Here T_c is the temperature at which SC features vanish in dV/dI.

Compound	$2\Delta(0)$ (μeV)	$2\Delta(0)/kT_c$	Method	Reference
$CeCu_2Si_2$	$\geqslant 160$	$\geqslant 2$	Estimated	de Wilde *et al* (1994)
UBe_{13}	260	3.4	FW	Nowack *et al* (1987a)
	290	4.2	From I–V	Moreland *et al* (1994)
UPt_3	140	3.5	FW	Nowack *et al* (1987a)
	58–78	1.3–1.8	Isotropic BTK	Goll *et al* (1993)
	150	3.96	Anisotropic BTK	de Wilde *et al* (1994)
	51–107	2.6–4.4	Anisotropic BTK	Goll *et al* (1995a, b)
	100	~ 2.6	Modified BTK	Naidyuk *et al* (1996b)
U_2PtC_2	440	3.4	Isotropic BTK	Naidyuk *et al* (1988)
URu_2Si_2	~ 1000	~ 9	PS	Nowack *et al* (1992)
	420	~ 4	Anisotropic BTK	Hasselbach *et al* (1992)
	700	~ 6	Anisotropic BTK	Samuely *et al* (1995)
	$\geqslant 340$	$\geqslant 3$	Estimated	de Wilde *et al* (1994)
a-direction	~ 500	~ 5	Modified BTK	Naidyuk *et al* (1996c)
c-direction	~ 1400	~ 12	Modified BTK	Naidyuk *et al* (1996c)
BCS superconductor				
Zn	250	3.45	Modified BTK	Naidyuk *et al* (1996d)

with an isotropic gap. Their analysis favours two-dimensional representation of the order parameter with the orbit part having a line of nodes in the basal plane and point nodes along the *c*-axis.

de Wilde *et al* (1994) also claimed that the $dV/dI(V)$ behaviour of UPt_3 PCs is attributable to the presence of nodes in the SC order parameter $\Delta(\boldsymbol{k})$, whereas data for $CeCu_2Si_2$ and URu_2Si_2 point to the absence of zeros in $\Delta(\boldsymbol{k})$. This is in contrast to the statement of Hasselbach *et al* (1992) and Samuely *et al* (1995) that an anisotropic d-wave order parameter describes the $dV/dI(V)$ characteristics of URu_2Si_2 better. However, it is worth noting that according to Hasselbach *et al* (1992) the lowest temperature was only about $T_c/2$ and the SC features were very weak, while Samuely *et al* (1995) used a polycrystalline sample, which does not allow an unequivocal conclusion to be reached concerning the gap anisotropy.

Naidyuk *et al* (1996c) have found a clear double-minimum structure in the $dV/dI(V)$ curves for some (about 5% of the total amount) PCs between single-crystal URu_2Si_2 and Pt, as expected in the case of Andreev reflection with a non-zero barrier parameter Z. The structure vanishes on reaching T_c or the upper critical magnetic field B_{c2} of URu_2Si_2. A fitting of the dV/dI curves (see figure 12) according to the modified BTK model yields a SC gap for the *c*-axis that is a few times the size of that for the *a*-axis. This points to a strong anisotropy of the SC gap in URu_2Si_2, probably because of the suppression of superconductivity in the *ab*-plane due to the exchange field of the ferromagnetically ordered layers of the U magnetic moments in the basal plane according to the model of Kulic *et al* (1991). They also predict $2\Delta(0)/k_BT_c$ to be up to a few times the size of the standard BCS value, 3.5, which is indeed observed; see table 3. Another point in support of the model is that a double-minimum structure was found for PCs with suppressed AF features (i.e., the

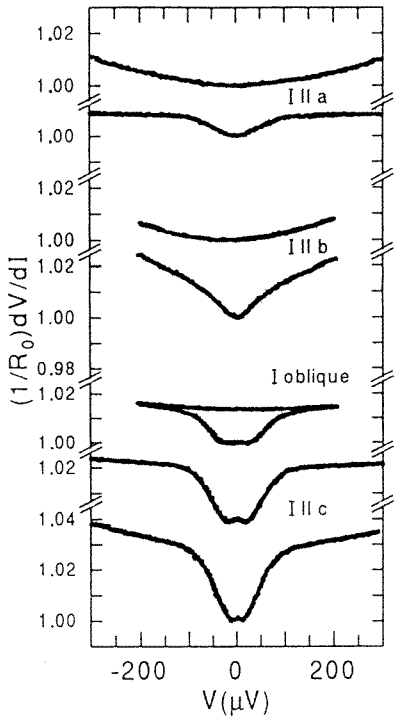

Figure 11. $\mathrm{d}V/\mathrm{d}I(V)$ curves for PCs on UPt_3 single crystal measured at $T < 50$ mK for preferred current directions as indicated. After Goll *et al* (1993).

broad $\mathrm{d}V/\mathrm{d}I$ maximum at $V = 0$ for the normal state was absent; see section 5).

Gloos *et al* (1996a, b, c) were the first to critically analyse the contribution to $\mathrm{d}V/\mathrm{d}I(V)$ for S–c–N contacts arising from the Maxwell term of the PC resistance; see equation (13). It was found that the decrease of the PC zero-bias resistance δR below T_c scales inversely with the radius of the PC (figure 13), determined by equation (9), instead of scaling inversely with the PC area as expected for AR. Additionally, the magnitude of δR corresponds quite well to the Maxwell contribution of the complete PC resistance with ρ, which is as for the bulk. This gives strong evidence that the SC minimum of the $\mathrm{d}V/\mathrm{d}I$ curves at $V = 0$ is mainly due to the Maxwell resistance vanishing in the SC state. On the other hand, the zero-bias PC resistance at low temperatures, $R(V = 0, T \simeq 0)$, in the SC state scales fairly well with the inverse contact area, albeit being approximately two orders of magnitude higher than the value estimated from the Sharvin formula, equation (2). This means that the Andreev-reflected hole current leading normally to a decrease of the PC zero-bias resistance is suppressed in these contacts. We recall that, according to the theory of Artemenko *et al* (1979), $R(V = 0, T \simeq 0)$ can still increase, but not to more than the R_N-value for diffusive S–c–N junctions with $d \ll \xi$.

The question arises of why superconductivity is destroyed in PCs with increasing voltage. The simplest cause would be the local heating of PCs at the rate $eV \simeq 3.63k_BT$, taking into account the large specific resistivity of HF compounds and the prevalence of the Maxwell resistance for typical PC resistances of up to a few tens of ohms; see figure 3. The heating in SC homocontacts is unambiguously demonstrated in the next section. Nevertheless, there are several experimental contradictions to the heating model.

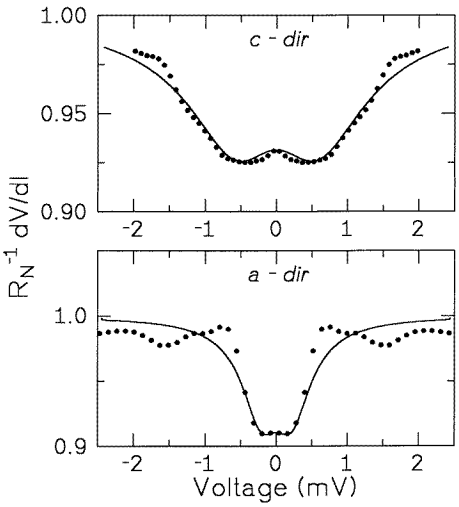

Figure 12. The symmetric part of $\mathrm{d}V/\mathrm{d}I(V)$ obtained for a URu_2Si_2–Pt contact (data points) and that calculated according to equations (15) and (16) (solid curves) for two directions at $T \simeq 50$–80 mK. The fitted parameters for the curves along the *c*- and *a*-axes are: $\Delta = 0.6$ meV, $Z = 0.43$, $\Gamma = 0.65$ meV; and $\Delta = 0.25$ meV, $Z = 0.33$, $\Gamma = 0.25$ meV, respectively. After Naidyuk *et al* (1996c).

As was shown for URu_2Si_2 in section 5, for some contacts non-heating effects may determine I–V characteristics at low biases. Another example of a non-heating process is the AR for UPt_3–Zn PCs observed by Naidyuk *et al* (1996b). Here the AR features of Zn were fitted quite well at different temperatures on the basis of the supposition that $T = T_{\mathrm{bath}}$; that is, no noticeable increase of the temperature with the bias in the PC occurs. Nevertheless, the trivial heating should be taken into consideration in each case, before dealing with more exotic phenomena in PCs.

Another cause of the suppression of the superconductivity might be the current density exceeding the critical value, or the influence of a self-magnetic field. For all of these causes, as well as for heating, the excess current should disappear at large biases according to Iwanyshyn and Smith (1972). Experimentally, an excess current is often observed up to very high voltage biases, especially for highly Ohmic junctions.

Finally, if one assumes that the SC state in a HF electrode remains intact at high voltage biases, then the appearance of a Maxwell term poses no problem, since quasiparticles with energies larger than Δ can propagate freely into a superconductor, involving the penetrating longitudinal electric field and the full normal differential resistance. Since in a normal-metal electrode the charge-carrier movement is conservative, it is still possible to obtain spectroscopic information about the energy gap in superconductors, without there being an AR hole current.

6.2. S–c–S contacts

Poppe (1985) and Han *et al* (1985) were among the first to perform measurements of PCs between ordinary singlet and HF superconductors. Figure 14 shows the effect of temperature and a field on the I–V characteristic of a $CeCu_2Si_2$–Al contact. The I–V curves in the SC state display a supercurrent at $V = 0$ and an excess current at a finite

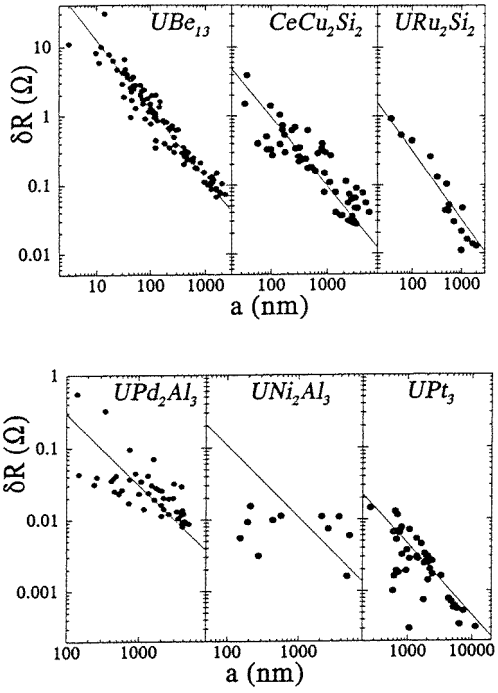

Figure 13. The resistance drop δR at zero bias and low temperatures versus the PC radius for contacts between tungsten and various HF compounds. The solid lines show $\delta R \propto d^{-1}$. After Gloos *et al* (1996a).

voltage. From the absolute magnitude of the maximal supercurrent I_c, which reaches 15–80% of the theoretical value for ordinary superconductors (see Ambegaokar and Baratoff (1963a, b) and equation (19)), Poppe (1985) concludes that $CeCu_2Si_2$ behaves like an s-wave superconductor. It is worth noting that the excess current vanishes at the temperature 0.7 K (figure 14), which is noticeably lower than T_c for Al, or in a magnetic field that is negligible with respect to the critical field of $CeCu_2Si_2$, i.e. when one of the electrodes is still superconducting. In this case (i.e., for a S–c–N contact), I_{ex} should still exist and be only a factor of two smaller than for a S–c–S contact (Zaitsev 1984). This observation raises questions as regards the formation of a weak link between two contacted superconductors; indeed it forces us to assume that a weak link is located inside the same material—probably in Al.

In the case of UPt_3, Poppe (1985) found neither Josephson coupling nor any other SC gap features in vacuum tunnelling experiments below T_c using both Pt and UPt_3 counter-electrodes.

Han *et al* (1985) observed an ac Josephson effect and a Fraunhofer-type pattern of the supercurrent in a magnetic field in PCs between $CeCu_2Si_2$, UBe_{13}, $LaBe_{13}$ and Nb, surprisingly well above the critical temperature of the HF samples. The authors explained the results as an anomalous proximity induced by Nb s-wave superconductivity in HF compounds. The further observation of I_c decreasing below T_c for UBe_{13} in a UBe_{13}–Nb PC, in contrast to the increase in I_c for a Mo–Nb contact, was interpreted by Han *et al* (1986) as indicating the occurrence of a negative s-wave proximity effect. They claim that these observations favour an unconventional SC ground state in HF. However, the model

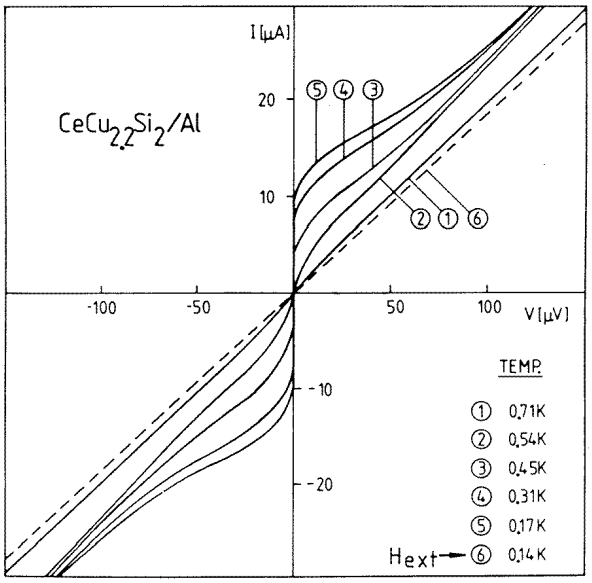

Figure 14. I–V characteristics of a $CeCu_2Si_2$–Al Josephson contact at different temperatures and under an external magnetic field. After Poppe (1985).

of a proximity-induced Josephson effect has been questioned by Kadin (1990). He found it fundamentally flawed, and proposed a possible alternative explanation based on the model of a phase-slip centre near the tip of a superconducting electrode in a S–c–N contact.

A dc Josephson effect between UPd_2Al_3 and Nb was reported by He *et al* (1992). They observed a persistent current in a composite Nb ring, bridged by a small pointed rod of UPd_2Al_3, and trapped flux, which was in discrete quantum states separated by the flux quantum $h/2e$. A supercurrent was observed in the I–V characteristics, but with a very small product I_cR_N: only a few tenths of a per cent of the theoretical value given by equation (19). The measurements confirmed a bulk superconducting state in UPd_2Al_3, but did not constrain the nature of the SC ground state in this material.

Nowack *et al* (1995a, b) have measured the $dV/dI(V)$ characteristics of PCs between URu_2Si_2 and the conventional superconductors Zn and NbTi. Contacts between URu_2Si_2 and Zn behave like a superconductor–normal-metal–superconductor junction with a thick normal layer at the HF part of the contact, and show Andreev-reflection structures typical for S–c–N contacts, related both to the heavy-fermion and the conventional superconductors. In contrast, contacts between URu_2Si_2 and NbTi become superconducting at low bias currents (i.e. show a Josephson dc current), though the corresponding critical voltage $V_c \approx 10\ \mu$V is about one order of magnitude below the theoretical value given by equation (19). Nowack *et al* (1995a, b) succeeded in fabricating a closed-loop set-up with two NbTi contacts on a URu_2Si_2 sample, which even showed oscillations in weak magnetic fields vanishing above the critical temperature of the HF superconductor; see figure 15. The last observation proves that the superconductivity of HF materials is involved. It is interesting that the authors detected a finite phase difference between two contacts at zero field (i.e. a minimum rather than a maximum of critical current), as expected for a SQUID between a conventional and a d-wave superconductor. The trapping of vortices near the SQUID and the problem that the residual magnetic field of the small SC coil used was unknown were overcome

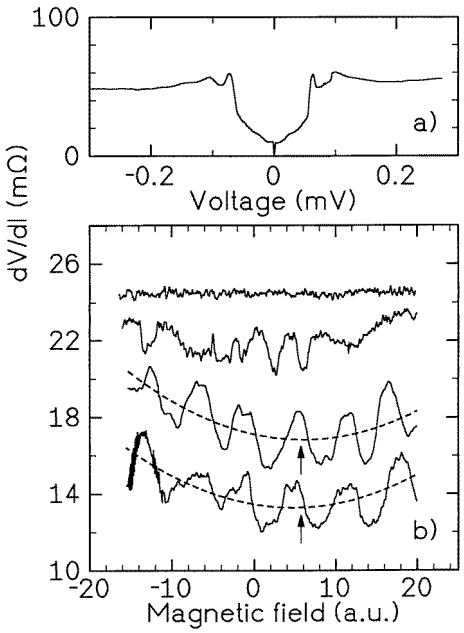

Figure 15. (a) $\mathrm{d}V/\mathrm{d}I(V)$ for a URu_2Si_2–NbTi contact with zero resistance at $V = 0$. (b) The oscillations of the differential resistance $\mathrm{d}V/\mathrm{d}I$ for this contact at $V_{\mathrm{bias}} \neq 0$ versus the magnetic field. The temperatures from top to bottom are: 1.3, 1.0, 0.6, and 0.2 K. The arrows show the position of the minimum of the background (dashed lines). After Nowack *et al* (1995a, b).

by appropriate location of zero field, at the minimum of the parabolic background, which position is known with a precision greater than the distance between the extrema of the oscillating curve. Unfortunately, it was not proved whether the phase difference equals zero when two contacts of the SQUID are established on the same crystal plane. Nevertheless, the Josephson coupling between a conventional superconductor and a HF superconductor URu_2Si_2 was successful, in spite of the fact that the coupling between the HF and the other superconductors is hindered by the formation of the above-mentioned normal layer at the surface of the HF material. Then Josephson effects are possible, assuming that the proximity-induced superconductivity in this thick normal-conduction zone is strong enough. This makes the effective layer of the normal zone thinner, leading to an increase of the coupling. The authors suggest that strong-coupling conventional superconductors with high values of T_c should be used to obtain Josephson contacts, and, in particular, to carry out Sigrist/Rice-type experiments (van Harlingen 1995) designed to measure the phase difference of the order parameter in unconventional superconductors.

One reason for the suppression of the superconductivity at the contact interface could be simply a mechanical stress in the contact region in connection with the sensitivity of the HF superconducting state to impurities or distortions. Nowack *et al* (1995a, b) speculate that the development of a normal layer is connected with the magnetic properties of URu_2Si_2, which orders antiferromagnetically.

An analogous experiment was performed by Naidyuk *et al* (1996b) with UPt_3. $\mathrm{d}V/\mathrm{d}I(V)$ for a UPt_3–Zn PC well below T_c showed a double-minimum structure at around $V = 0$, characteristic for Andreev reflection, due to the superconducting Zn. In magnetic fields above the critical field B_c of Zn, the $\mathrm{d}V/\mathrm{d}I$ curves exhibit only a shallow minimum

at $V = 0$, which is attributed to the superconducting UPt_3. A fit of the dV/dI curves on the basis of modified BTK theory yields for Zn an energy gap $\Delta(T)$ perfectly matching the prediction of BCS theory. For UPt_3, the fit gives a value of $\Delta(0)$ (see table 3) which is comparable with the values obtained by Goll *et al* (1993) using the ordinary BTK fit. It should be pointed out that the barrier strength parameter Z was found to lie in the range 0.45 ± 0.02—that is, to be the same as for Zn–Ag contacts (Naidyuk *et al* 1996d). This supports a theory of Deutscher and Nozières (1994) that the boundary condition at the interface with HF compounds involves Fermi velocities without a large mass-enhancement factor. The large Fermi-velocity mismatch between HF samples and ordinary metal would lead to a very large value of Z, characteristic for a tunnel junction. On the other hand, the obtaining of the same Z-parameter for quite different systems (including the contacts with high-T_c superconductors), with a value close to 0.5, hints that all of these contacts might be in the dirty limit, according to figure 9 given by Blonder *et al* (1982).

For UPt_3–NbTi contacts, Naidyuk *et al* (1996b) were unsuccessful in attempts to observe oscillations of the supercurrent in small magnetic fields, as had been achieved for URu_2Si_2 by Nowack *et al* (1995a, b). It looks like the surface layer of UPt_3 shows a stronger pair-breaking effect, taking into account the high Γ value (see equation (16)), which increases with the PC resistance (Naidyuk *et al* 1996b). The Γ parameter is about one order of magnitude larger than that for a Zn–Ag PC of the same size (Naidyuk *et al* 1996d). This gives clear evidence that the interface with a HF compound, namely UPt_3, produces a much stronger pair-breaking effect than one with a normal metal like Ag.

6.3. Mechanically controllable break junctions with HF compounds

The quality of the interface is certainly of great importance in PC experiments. Notable problems are contamination of the PC, for example by oxide and/or other surface dirt, and crystal lattice distortion due to mechanical stress. The technique of using mechanically controllable break junctions (MCBJ), developed by Muller *et al* (1992), seems to be the most effective method for getting clean, non-degraded interfaces, prepared '*in situ*', at low temperatures and under ultrahigh-vacuum conditions. Using identical metals on both sides of the junction permits one to avoid normal quasiparticle reflection and discard a contribution from the second dissimilar electrode, both in the normal-conductivity and in the SC state, neglecting in the latter case the undesirable proximity effects. Additionally, the MCBJ technique enables one to finely control the contact resistance over many orders of magnitude, and to measure I–V characteristics from low-Ohmic metallic contacts to high-resistance tunnel junctions.

The break-junction experiments on HF systems by Gloos *et al* (1997, 1998) confirmed, in general, previous findings obtained by Gloos *et al* (1996a) using HF heterocontacts. R_M was clearly resolved, while no clear-cut signatures of Andreev reflection or Josephson effects was found. Gloos *et al* (1998) claimed that the contribution from R_M of the metallic HF contacts would make it difficult to ever reliably extract any other intrinsic HF properties by using the concepts of Andreev reflection and the Josephson effect.

Naidyuk *et al* (1997) investigated in detail MCBJ with URu_2Si_2 single crystal. Here, the structure in $dV/dI(V)$ was carefully analysed to reveal its origin. A peaked structure in $dV/dI(V)$ often observed for HF compounds at some critical voltage V_c (figure 16) cannot be due to AR in the framework of BTK theory, and it is tentatively attributed to the destruction of superconductivity in PCs. It was concluded that, in the constriction, superconductivity is more probably suppressed by the local heating. Here a correspondence of V_c with T_c according to equation (5) was found, taking into account the fact that

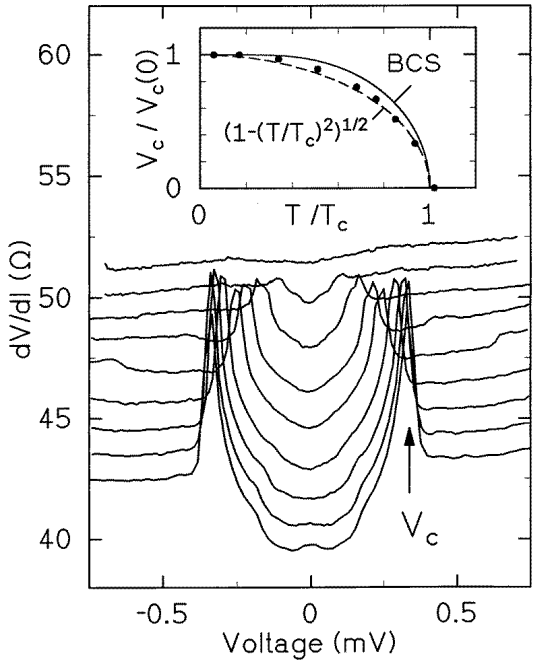

Figure 16. $\mathrm{d}V/\mathrm{d}I(V)$ along the a-direction for a URu_2Si_2 MCBJ with $R(0) \simeq 40\ \Omega$ at the following temperatures: 1.2, 1.1, 1.0, 0.9, 0.8, 0.6, 0.4, 0.2 K (from top to bottom). Inset: the position of the $\mathrm{d}V/\mathrm{d}I$ maximum V_c versus the reduced temperature T/T_c for the same PC ($T_c \simeq 1.1$ K). After Naidyuk *et al* (1997).

$L \simeq L_0$ below 1 K. Additionally, the temperature dependence $V_c(T)$ obeys equation (5) well (figure 16, inset); that is, the contact is actually in the thermal regime.

The coinciding of relation (6), $eV \simeq 3.63k_BT$, for the thermal regime with the well known BCS equality $2\Delta_0 = 3.52k_BT_c$, as well as the similarity of the temperature dependencies of the critical voltage $V_c(T)$ according to equation (5) and the BCS behaviour of $\Delta(T)$ (see figure 16, inset), can lead to misinterpretation of the thermal regime features in $\mathrm{d}V/\mathrm{d}I$ as spectroscopic ones, and to an erroneous determination of the spatial and temperature dependencies of the SC order parameter.

Naidyuk *et al* (1997) have found some $\mathrm{d}V/\mathrm{d}I(V)$ characteristics with a double-minimum zero-bias structure at voltages smaller than V_c (figure 16), which is usually viewed as clear-cut evidence of AR at a N–S interface. They also analysed Kondo scattering by localized magnetic moments (probably of uranium) in the contact region (see equation (11) and Jansen *et al* (1980), Yanson *et al* (1995)), which can also yield a maximum at zero bias and create rather similar features. The authors claimed that it is not simple to distinguish between AR and a Kondo contribution to a zero-bias maximum. It seems to be reasonable to include the Kondo scattering in the BTK theory. Conceivably, elastic spin-flip scattering can be formally described using the Z-parameter if AR and the Kondo effect do not block each other completely.

Gloos *et al* (1997, 1998) and Naidyuk *et al* (1997) investigated MCBJs with zero resistance, i.e., with Josephson-like I–V characteristics. Usually these contacts had a low normal resistance, less than 0.2 Ω in the case of URu_2Si_2. No oscillatory pattern of the critical current in a magnetic field was found. This may be due to the lack of a clear-cut

normal area penetrated by magnetic flux, like, for example, in a Josephson junction with a circular interface.

Alternatively, there is another explanation. Note that the ballistic contribution to the resistance of these contacts is extremely small. For $R_{PC} \leqslant 0.2\ \Omega$ or $d \geqslant 1\ \mu$m (see figure 3), the Sharvin resistance is less than 1/100 of the total normal-state resistance, assuming the bulk value of the residual resistivity and the typical metallic electron density of URu_2Si_2 (see Grewe and Steglich 1991). This means that the vanishing of Maxwell's resistance in the SC state can simulate the appearance of a supercurrent, while $R_{PC} \simeq R_M \approx 0$ because $R_M \gg R_{Sh}$.

The suppression of the Josephson effect could partly arise from a normal interface layer, as is also indicated by the enhanced residual PC resistance R_0 (Gloos *et al* 1997). For S–c–S contacts, the Josephson effect is expected to survive as long as the width of the region with the reduced order parameter (e.g. the normal layer) is less than 3ξ (Licharev 1985). Therefore, even a thin (about 30 nm) normal layer may suppress the Josephson coupling. This layer, which might be of magnetic origin, probably also hinders the observation of SC features by vacuum tunnelling (Poppe 1985, Goschke *et al* 1996), which 'senses' the topmost atomic layer.

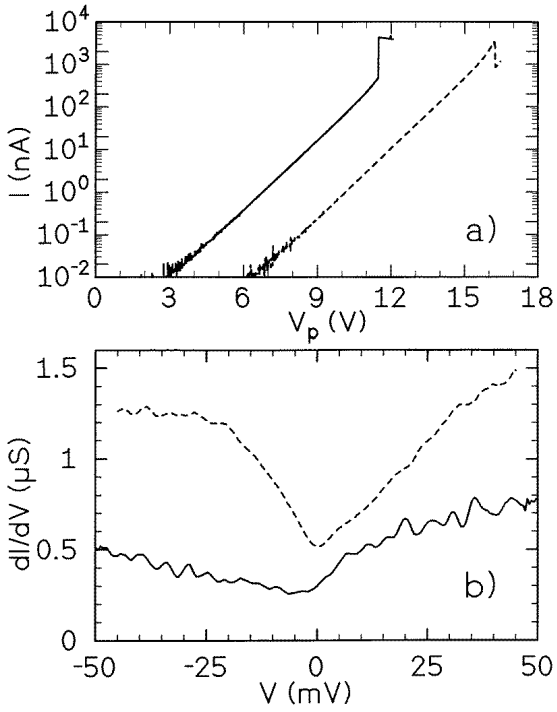

Figure 17. (a) The tunnelling current versus the change of piezo-voltage V_p at $T \leqslant 0.1$ K for successive piezo-voltage sweeps along the *a*-direction (solid line) and the *c*-direction (dashed line). The bias voltage is 0.1 V in both cases. (b) Numerical derivatives of I–V curves in the tunnel regime at $T \leqslant 0.1$ K for the *a*-direction (solid curve) and the *c*-direction (dashed curve). The PC resistance is a few MΩ for both cases. After Naidyuk *et al* (1997).

Recently, true vacuum tunnelling in MCBJs has been realized by Naidyuk *et al* (1997) with a perfect exponential dependence of the current on the width of the vacuum gap (figure 17). The authors report a lack of clear SC features as well as a supercurrent in the

tunnel regime, which demonstrates that this takes place also for atomically clean surfaces. Here, only a broad asymmetric minimum in $\mathrm{d}I/\mathrm{d}V(V)$ is obtained, with a width more than one order of magnitude larger than the expected SC energy gap. The more pronounced minimum in $\mathrm{d}I/\mathrm{d}V(V)$ is for the *c*-direction, contrary to the observations of Aliev *et al* (1991) and Aarts *et al* (1994), which questioned the formation of this feature in view of the gapping of the Fermi surface in the AF state. It should be mentioned that, in consequence of the small capacitance C of MCBJs, the Coulomb-blockade energy $e^2/2C$ can be larger than the thermal energy $k_\mathrm{B}T$, which can produce a similar zero-bias minimum in $\mathrm{d}I/\mathrm{d}V$ (van Bentum *et al* 1988). On the other hand, Gloos *et al* (1998) pointed out that the small capacitance (~ 1 f F) of MCBJs also results in a zero-point energy

$$E_0 = \sqrt{e\bar{h}I_\mathrm{c}/2C}$$

of the Josephson plasma, which can easily exceed the Josephson coupling energy $E_\mathrm{J} = hI_\mathrm{c}/2e$ and suppress the supercurrent I_c.

The SC features in the I–V characteristics of MCBJs on HF superconductors vanish, as mentioned by Naidyuk *et al* (1997) and Gloos *et al* (1998), as soon as the contact resistance is above 1 kΩ—that is, well below the transition from the metallic to the tunnelling regime, which occurs at $R_\mathrm{q} = h/2e^2 = 12.9$ kΩ. This fact points additionally to a non-superconducting heavy-fermion surface.

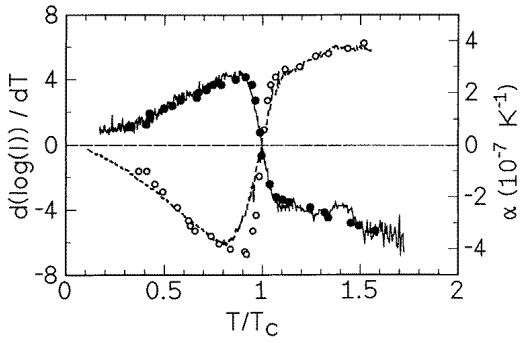

Figure 18. The logarithmic derivative of the tunnel current versus the reduced temperature T/T_c for a URu_2Si_2 MCBJ along the *c*-direction (solid curve) and the *a*-direction (dashed curve) at a bias of 0.1 V and for a constant piezo-voltage. After Naidyuk *et al* (1997). Open and closed circles show the temperature dependence of the thermal expansion coefficients for two corresponding directions. After van Dijk *et al* (1995).

Yet, systematic study of the tunnelling regime with MCBJs yields additional rather interesting information (Naidyuk *et al* 1997). It indicates a nearly free-electron mass of the tunnelling charge carriers, as well as a typical metallic carrier density or an ordinary Fermi wavenumber for URu_2Si_2. The curious thing is that the logarithmic derivative of the temperature dependence of the tunnelling current shows in a direct way the behaviour of the thermal expansion coefficient and its anisotropy in URu_2Si_2 (figure 18).

7. Conclusion

As has been shown in this review, after more than a decade of investigation of HF systems by means of PCs, the studies have yielded considerable knowledge about the physical processes therein and the HF properties themselves, both in normal and SC states. Here we

would like to emphasize mainly the problems which should be accounted for and overcome in further PC research into HF materials.

Still the basic challenge is that of how to separate the Joule-heating effects from the spectroscopic ones or to discriminate a huge thermal contribution to the PC resistance from the ballistic and/or diffusive transport of eV-energized electrons. It is common knowledge that even for some clean metals a thermal regime occurs with increasing voltage above 100 mV (Verkin *et al* 1979). Suppose that the short inelastic mean free path of electrons $l_{\rm in}(eV)$ results in heating of the PC—one expected to overcome heating at low bias voltages—where inelastic scattering is negligible. Then the question is at what voltage does this heating occur for PCs with HF compounds. Certainly, this voltage is different both for different HF materials and for different contacts. It is clear that the smaller the resistivity of the bulk samples, or the smaller the PC size, the less the contribution from the Maxwell resistance and the greater the damping of the heating. Since it is impossible to estimate $l_{\rm in}$ for a concrete contact and definite bias voltage *a priori*, the regime of current flow may be determined only by analysing the PC characteristic. For example, for a normal-conduction PC, the thermal regime was irrefutably proved to be dominant by Naidyuk *et al* (1985a) and Paulus and Voss (1985). Conversely, Naidyuk *et al* (1995) established that non-thermal transport predominates for some URu_2Si_2 PCs, taking into account the qualitative difference between the voltage $\mathrm{d}V/\mathrm{d}I(V, T = \text{constant})$ and the temperature dependence $\mathrm{d}V/\mathrm{d}I(V = 0, T)$ of the PC resistance. Therefore it is very desirable in all cases to measure $\mathrm{d}V/\mathrm{d}I(V = 0, T)$, which should behave as $\rho(T)$. This proves additionally that the properties of the material in the constriction correspond to those of the bulk.

The freezing out of normal scattering processes contributes considerably to the SC signal $\mathrm{d}V/\mathrm{d}I(V)$ of PCs of HF compounds as was shown by Gloos *et al* (1996a, 1998). They claimed that one should not obtain spectral information from such contacts, unless one takes into account $R_{\rm M}$. In spite of the pronounced AR-like feature in $\mathrm{d}V/\mathrm{d}I(V)$ for S–N contacts given by Goll *et al* (1993) and Naidyuk *et al* (1996c), more attempts should be made to clarify the small AR signal, and the development of microscopic models of AR in HF compounds is very necessary. Additionally, for SC contacts, one should distinguish between a number of other processes involved, mainly local heating at the rate of approximately $k_{\rm B}T = eV/3.63$, and destruction of the superconductivity by a high current density, or by the self-magnetic field of the current through the junction. After this has been done, the spatial as well as the temperature dependencies of the SC order parameter (the energy gap) may be recovered unambiguously from the PC spectra.

Undoubtedly, one would expect to obtain the most spectacular results as regards a SC ground state of a HF system from Josephson or SQUID measurements. Here, progress is rather slow in comparison with that for high-$T_{\rm c}$ materials (see van Harlingen 1995). Nevertheless, the first success in realizing Josephson coupling between HF compounds and 'strong-coupling' conventional superconductors achieved by He *et al* (1992) and Nowack *et al* (1995a, b) gives hope for further progress, probably with microstructural contacts for SQUIDs or other artificial weak links.

It should be noted that traditional tunnelling spectroscopy (Wolf 1985) measurements on the HF systems are almost entirely lacking. The absence of SC features in the vacuum-tunnelling regime of MCBJs on HF systems has been reported by Moreland *et al* (1994) and Naidyuk *et al* (1997), as well as by Poppe (1985) and Goschke *et al* (1996), on the basis of STM experiments. This shows that the properties of the HF interface matter, and the question is whether it is actually possible to observe the intrinsic properties of HF compounds by means of tunnelling experiments.

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