Scattering of slow neutrons from long-wavelength magnetic fluctuations in UPt_3

1995 J. Phys.: Condens. Matter 7 7325

(http://iopscience.iop.org/0953-8984/7/37/006)

J. Phys.: Condens. Matter 7 (1995) 7325–7333. Printed in the UK

Scattering of slow neutrons from long-wavelength magnetic fluctuations in UPt_3

N R Bernhoeft†‡‖ and G G Lonzarich§

† Physics Department, University of Durham, Durham, UK

‡ Centre d'Etudes Nucléaires de Grenoble, BP85X, 38041 Grenoble Cédex, France

§ Cavendish Laboratory and The Interdisciplinary Centre on Superconductivity, University of Cambridge, Cambridge CB3 0HE, UK

Received 15 March 1995, in final form 2 June 1995

Abstract. We report the observation of a low-energy contribution to the cross section of scattering of slow neutrons near the forward direction in the normal state of UPt_3 at low temperatures. The results point to the existence of a weak component in the imaginary part of the dynamical wavevector-dependent susceptibility which may be represented approximately in a one-pole model with a wavevector-dependent relaxation frequency $\Gamma(q)$. In the low-q range of the measurements, $\Gamma(q)$ is up to an order of magnitude lower than the nearly q-independent characteristic frequency of a previously studied and well established dominant component of the dynamical susceptibility. The new low-frequency contribution is interpreted in terms of the effect of low-lying fermion quasi-particles in the presence of strong spin–orbit coupling.

1. Introduction

The low-energy and long-wavelength properties of highly degenerate Fermi systems are conventionally represented in terms of fermion quasi-particle and hole excitations. Near the Fermi level and at low temperatures, collisions between these excitations are assumed to be effectively suppressed in the *normal* state by the Pauli principle, so that the effects of residual interactions, although not necessarily weak, may be treated in a self-consistent field manner, i.e. in terms of the Landau Fermi-liquid model [1].

When the quasi-particles carry a charge, the applicability of this microscopic picture may be checked, in minute detail, by the study of phenomena associated with the quantization of cyclotron motion in a magnetic field, under the condition that the cyclotron radius is much greater than the effective quasi-particle 'dimension' but smaller than its mean free path. For either charged or neutral quasi-particles, the microscopic predictions of the Fermi-liquid model may also, in principle, be investigated by means of the scattering of slow neutrons from the local magnetic fields generated by long-wavelength fluctuations in the quasi-particle spin density.

The signature of the Fermi-liquid state in this case may be found in the wavevector dependence of the rate of decay of spontaneous spin fluctuations at long times. In the absence of collisions, the ballistic transport of spin by the quasi-particles may be expected to bring about the decay of a component of a spontaneous fluctuation of wavevector q over a time comparable to that required for a quasi-particle to travel over a wavelength $2\pi/q$,

‖ Present address: European Synchrotron Radiation Facility, BP220, F38043, Grenoble Cédex, France.

0953-8984/95/377325+09$19.50

and hence at a rate of the order of $v_F q$, where v_F is the characteristic *quasi-particle* Fermi velocity.

The interaction between quasi-particles in this collisionless regime does not change the qualitative q dependence of the spin relaxation rate but suppresses the magnitude of the linear coefficient by the reciprocal of the Stoner enhancement factor, defined as the ratio of the overall spin susceptibility to that obtained when the quasi-particle interactions are formally neglected. For an isotropic homogeneous system the final relaxation rate is then of the order of $v_F q(1 + F_0^a)$, where F_0^a is the isotropic and spin antisymmetric Landau parameter.

This elementary picture may be expected to apply for a simple Fermi liquid when the total spin is conserved and collisions can be ignored. Collisions with other excitations or defects lead to diffusive rather than ballistic transport by the quasi-particles and hence in the low-q limit to a quadratic rather than a linear fluctuation relaxation spectrum. Non-conservation of spin in the presence of strong spin–orbit coupling on the other hand can lead to a relaxation of the spin density locally against the lattice without long-range transport at low q. Even if the total spin is *not* conserved, however, some components of a spontaneous fluctuation of the magnetization of wavevector q may, presumably, still be expected to decay slowly in the collisionless Fermi-liquid region as $q \to 0$ if the effective magnetic moment carried by a quasi-particle is *not* vanishingly small.

This slow quasi-particle component of the spin relaxation process has indeed been identified experimentally in liquid ^{3}He [2], in paramagnetic d metals with strongly exchange-enhanced susceptibilities Ni_3Ga [3] and $TiBe_2$ [4] and in the d metals with small spin polarizations such as MnSi [5, 6], Ni_3Al [7] and $ZrZn_2$ [8], systems whose behaviour in the normal state at low temperatures, long wavelengths and low energies seem to be in general accord with the Fermi-liquid model. In f-electron heavy-fermion compounds, in which the spin–orbit coupling is strong and the f electrons are in a highly correlated, nearly localized state, however, the magnetic fluctuation spectrum appears to be dominated by the effects of local non-Fermi-liquid relaxation processes (see [9] for a recent review of the field), and the very existence of a slow quasi-particle component might be called into question.

In this paper we present experimental evidence that a finite part of the overall spectrum of magnetic fluctuations in one of these materials, UPt_3, at long wavelengths and low frequencies in the normal state may indeed exhibit properties not unlike that anticipated by the Fermi-liquid model, but as modified by the coupling of the fermions to the dominant non-Fermi-liquid component of fluctuations in the local magnetization. These findings complement the earlier work in this system [10, 11] and are qualitatively consistent with the observation, via quantum oscillatory phenomena, of a well defined Fermi surface of heavy excitations which account, within the framework of the Fermi-liquid theory, for the linear coefficient of the specific heat in the normal state [12].

2. Experimental details

The sample of UPt_3 was selected from ingots originally purified for the study of the de Haas–van Alphen effect [12] and subsequently used in a series of thermodynamic and transport measurements [13]. The mean free path of charged quasi-particles, as inferred from the magnetic field dependence of the amplitude of the de Haas–van Alphen oscillations, is in excess of 1000 Å in the low-temperature limit of the normal state. As attention was focused on forward scattering averaged over spin and crystallographic orientations, a polycrystalline sample (more readily obtained in a highly ordered and pure form than a single crystal) could be employed. It was in the shape of a cylinder 40 mm long and 10 mm in diameter, composed of grains typically 1–2 mm in dimension.

The energy- and momentum-resolved neutron scattering cross-section was measured by means of a time-of-flight technique on the IN5 spectrometer at the European High-Flux Reactor, Institut Laue–Langevin, Grenoble. The high available flux at long wavelengths, the low background at small scattering angles near the forward direction, the simultaneous detection over wide angle and energy transfer, and the high resolution made this spectrometer ideally suited to the aims of this study. At selected temperatures, the cross section $d^2\sigma/(d\Omega dE)$ was measured for unpolarized 3.3 ± 0.1 meV incident neutrons scattered at a polar angle θ (the angle between the incident and scattered beams) in the range $4° < \theta < 15°$, averaged over the azimuthal angle in the range $0° < \phi < 360°$, with energy transfer -1.5 meV$< E < 4$ meV. The scattering wavevector $q(\theta, E)$ may be evaluated for each θ and E via the simultaneous energy–momentum conservation condition

$$q(\theta, E)^2 = k_N^2[2 + E/E_N - 2\cos\theta(1 + E/E_N)^{1/2}] \tag{1}$$

where E_N and $\hbar k_N$ are the energy and momentum, respectively, of the incident neutrons. The scattering wavevector $q_0 = q(\theta, 0)$ at zero energy transfer was thus in the range $0.09\ \text{Å}^{-1} < q_0 < 0.33\ \text{Å}^{-1}$, and well below the typical dimensions of the Brillouin zone (the average radius of which is $0.75\ \text{Å}^{-1}$). The resolutions of the elastic scattering wavevectors and of the energy transfers are in the ranges 0.01–0.05 Å^{-1} FWHM and 0.05–0.2 meV FWHM, respectively.

The focus on scattering near the forward direction rather than in the neighbourhood near a reciprocal-lattice vector not only permits the use of polycrystalline rather than single-crystal samples but also tends to maximize scattering from fluctuations associated with the spin and orbital moments of electrons relative to that of the density of nuclei. Further, to highlight contributions of thermal excitations in the electron system, e.g. of activated quasi-particle and hole pairs, we consider the difference scattering relative to that observed in the low-temperature limit. The results discussed below represent the scattering relative to the background at 1.3 K at a temperature of 10 K which is high enough to have a significant thermal activation of low-lying magnetic fluctuations and lies in the heavy-fermion 'coherent' state below the maximum of the temperature-dependent magnetic susceptibility.

To minimize noise a carefully shielded cryostat with low background scattering was employed and, in some data sets, incident neutrons with wavelengths above the Bragg cut-off have been used to suppress elastic nuclear multiple-scattering processes. The results of these tests at long wavelengths are consistent with those presented here with $\lambda_N = 5.0 \pm 0.1$ Å. The absolute calibration of the cross section was performed by monitoring the elastic nuclear incoherent scattering from a sample of purified vanadium outgassed under ultra-high vacuum.

3. The cross section

The scattering of neutrons through the interaction of their magnetic moment with a space- and time-varying magnetic field $\boldsymbol{B}(\boldsymbol{r}, t)$ may be usefully expressed in terms of the power spectrum of such fluctuations [14]. For an unpolarized plane-wave beam and an isotropic body in the Schrödinger–Pauli model, the probability per atom, incident flux, solid angle Ω and energy E of neutrons scattered from initial momentum $\hbar\boldsymbol{k}$ to a final value of

$\hbar k' = \hbar(k + q)$ with increase in energy $E = \hbar\omega$ can be expressed as

$$\frac{d^2\sigma}{d\Omega\, dE} = \frac{\gamma_N^2 \alpha}{48\pi^3} \frac{1}{\hbar^2 cn} \frac{|k'|}{|k|} \langle |B(q, \omega)|^2 \rangle \tag{2}$$

where γ_N is the neutron gyroscopic factor (−1.913), α is the fine-structure constant (1/137), c is the speed of light, n is the number of atoms per unit volume, and $\langle |B(q, \omega)|^2 \rangle = \langle B(-q, -\omega) \cdot B(q, \omega) \rangle$ is the appropriate power spectrum of field fluctuations defined as the Fourier transform in r and t of the *asymmetrized* autocorrelation function $\langle B(r', 0) \cdot B(r' + r, t) \rangle$ in which r' is averaged over the volume of the system.

Contributions to $B(r, t)$ arise from fluctuations in the electron spin density, atomic orbital currents and spatially extended transverse currents. In our chosen conditions with low incident neutron energy and small scattering angle to the forward direction the dominant contribution may be expected to be due to spin and associated orbital moment fluctuations for which one may write in an isotropic model

$$\langle |B(q, \omega)|^2 \rangle = (4\pi)^2 \frac{6\hbar\omega}{\exp(\hbar\omega/k_B T) - 1} \frac{\mathrm{Im}[\chi(q, \omega)]}{\omega} \tag{3}$$

where $q = |q|$ and $\mathrm{Im}[\chi(q, \omega)]$ is the imaginary (absorptive) part of the generalized magnetic susceptibility.

When the low-frequency properties of the generalized susceptibility may be modelled approximately by a single imaginary pole we may write for sufficiently small ω and q

$$\mathrm{Im}[\chi(q, \omega)] = \omega z \chi \frac{\Gamma(q)}{\omega^2 + \Gamma(q)^2} \tag{4}$$

where χ is the static susceptibility, z measures the relative weight of the low-frequency pole and $\Gamma(q)$ is the pole frequency or effective relaxation rate [15, 16]. When (4) is employed to model the quasi-particle contribution in the collisionless regime of the Fermi-liquid theory ($\omega \ll v_F q$), the relaxation rate may be expressed in leading order in q as

$$\Gamma(q) = \gamma \chi^{-1} q \tag{5}$$

where γ is a q-independent parameter and $q \ll k_F$, where k_F is the characteristic Fermi wavevector. For a homogeneous isotropic Fermi liquid in which spin is conserved overall, $z = 1$ and $\gamma = 2v_F\chi(1 + F_0^a)/\pi = 2\mu^2 k_F^2/\pi^3$, where μ is the magnetic moment of the quasi-particles. Since k_F and μ are not renormalized by particle and spin-conserving interactions, γ is here an invariant depending only on bare parameters. As discussed in section 5, this conclusion no longer holds when spin conservation breaks down.

For an anisotropic system with strong spin–orbit coupling, such as UPt_3 which crystallizes in a hexagonal structure, the generalized susceptibility is in general a tensor with an anisotropic wavevector dependence. In the long-wavelength and low-frequency limit, the effect of anisotropy on the cross section, averaged over spin and crystallographic orientations, may be taken into account in UPt_3 by replacing the scalar susceptibility χ in (4) by a suitable average over the tensor components.

4. Results

The central results of this study are shown in figure 1 in which is plotted the difference cross section $I(\theta, E) = \Delta[d^2\sigma/(d\Omega\, dE)]$ at 10 K relative to the background at the reference temperature $T_0 = 1.3$ K as a function of neutron energy transfer at elastic wavevectors 0.11, 0.18, 0.25 and 0.33 Å^{-1}. The cross section is averaged over crystallographic orientations,

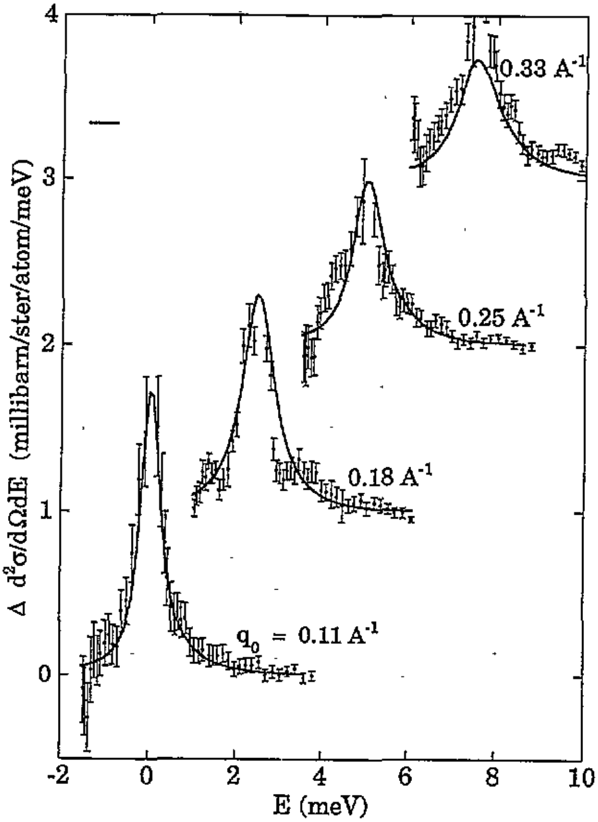

Figure 1. Energy- and momentum-resolved scattering cross-section $\Delta[\mathrm{d}^2\sigma/(\mathrm{d}\Omega\,\mathrm{d}E)]$ for 3.3 meV incident neutrons measured by the time-of-flight method in the forward direction from UPt_3 where E is the energy transferred to the neutron. The data are the differences between the cross section, averaged in the manner described in the text, at 10 K and that at 1.3 K. In frames left to right the average elastic scattering wavevector increases as 0.11 Å^{-1}, 0.18 Å^{-1}, 0.25 Å^{-1} and 0.33 Å^{-1}, and the origins have been shifted by 2.5 meV in the abscissa and 1 mb/sr/meV/average atom in the ordinate. The solid lines are fits using the cross section convolved with the instrumental resolution and the single-pole model for $\chi(q,\omega)$ given by (4) and (5) with pole weight $z \simeq 0.18$, the bulk average susceptibility $\chi \simeq 2.0 \times 10^{-4}$ cgs [17] and $\hbar\gamma \simeq 0.5$ μeV Å^{-1}.

the azimuthal angle of scattering and over ranges in wavevectors ± 0.05 Å^{-1} about the central values, just given. To interpret these findings we consider first an approximate form of the difference cross section in a one-pole model for $\chi(q,\omega)$ at low q and ω. From (2)–(4) and the assumption that the dominant temperature dependence arises from the Bose factor in (3), we find, for E well below $E_N \simeq 3.3$ meV (so that $k'/k \simeq 1$ and $q \simeq q_0$) and well above $k_B T_0$,

$$I(\theta, E) \simeq I(\theta, 0)\frac{E/k_BT}{\exp(E/k_BT)-1}\frac{1}{1+(E/\hbar\Gamma(q_0))^2} \tag{6}$$

$$I(\theta, 0) = \sigma_0 k_B T z\chi/\hbar\Gamma(q_0) \tag{7}$$

where σ_0 is a constant given by (2) and (3). Thus, $I(\theta, E)$ has essentially the form of a Lorentzian with a HWHM of $\Gamma(q_0)$ times the normalized Bose factor of characteristic energy $k_BT \simeq 1$ meV. Furthermore, the central peak intensity $I(\theta, 0)$ varies inversely as

$\Gamma(q_0)$. Hence, a wavevector-independent relaxation rate would produce an identical set of spectra, while a rate decreasing with decreasing q_0 should exhibit spectra which narrows in energy and intensifies at the origin as q_0 falls.

The spectra in figure 1 display the latter type of behaviour. The possibility that $\chi(q,\omega)$ in UPt_3 may be modelled at low q and ω by a *single* pole with a *q-independent* linewidth is hence ruled out. This is the key conclusion of this study.

At the lowest value of q_0 the HWHM of the spectrum is 0.5 meV and hence is dominated by the Lorentzian in (6) rather than by the Bose function; at the highest q_0, however, both factors are important. The behaviour of the spectra and variation in $\Gamma(q_0)$, by a factor of between 2 and 3 over the experimental q_0 range, are not inconsistent with that expected qualitatively from a contribution from quasi-particle excitations within a Fermi-liquid model (5). A fit of the spectra to this model with χ set equal to 2×10^{-4} cgs [17] yields estimates for z and $\hbar\gamma$ of 0.18 and 0.50 μeV Å, respectively to a precision of 15%. The fits, shown by solid lines in figure 1, are based on the full expressions for the cross section defined by (2)–(4) and convolved with the spectrometer resolution function, the energy width of which is only 0.07 meV at the elastic position and hence is small compared with $\Gamma(q_0)$ in the experimental range.

These studies are complementary to and supplement previous neutron scattering measurements in this compound [10, 11]. In particular, they pertain to a previously uncovered region in energy–momentum space in which the weight z for the low-energy Lorentzian is only approximately 20% of the total to be expected from an integration of the full $\mathrm{Im}[\chi(q,\omega)/(\pi\omega)]$ over all ω. *The existence of this weak component is not inconsistent* with the presence of a higher-energy part of the spectrum measured in earlier experiments having a characteristic energy, weakly q dependent at small q, of the order of 5–8 meV [9]. This part of the spectrum lies mostly outside the energy window in our difference cross section in figure 1 which is determined by the normalized Bose factor in (6).

5. Discussion

The imaginary part of the generalized magnetic susceptibility at low q and ω in UPt_3 appears to be characterized by a very-low-frequency component of small weight ($z \simeq 0.18$) with a dispersive relaxation rate (roughly represented by $\hbar\Gamma(q) \simeq (2.5$ meV Å$)q$ in the experimental range), and a more dominant higher-frequency part [10, 11] with a weakly q-dependent characteristic energy at low q (of the order of 5 meV). The impossibility of understanding our low-frequency data in terms of the latter *higher-frequency component alone* is graphically illustrated in figure 2 (see caption).

These findings differ from those observed in a number of d transition metals close to ferromagnetic instabilities at low temperatures in which the low-q, ω form of $\mathrm{Im}[\chi(q,\omega)]$ is well represented by the model (4) and (5) but with a strongly wavevector-dependent static susceptibility $\chi(q)$, z of order unity and $\hbar\gamma$ in all cases in a narrow range between 2 and 3 μeV Å. The ratio γ/z is comparable to that obtained for UPt_3, but the individual values are about five times greater. The dominance of the quasi-particle component of $\mathrm{Im}[\chi(q,\omega)]$ at low q,ω in the former materials is consistent with approximate conservation of the total spin, i.e. the presence of relatively weak spin–orbit coupling. The value of γ in this limit and in the simplest model given in section 3 depends on only μ_B and k_F and hence may be expected to be weakly varying among similar d systems. Indeed, this model yields estimates of γ close to that actually observed, for values of Fermi wavevectors (about 0.8 Å^{-1}) of the order of that calculated for the dominant d sheets of the Fermi surfaces in these materials.

The importance of spin–orbit coupling in the f systems leads us to consider a minimal

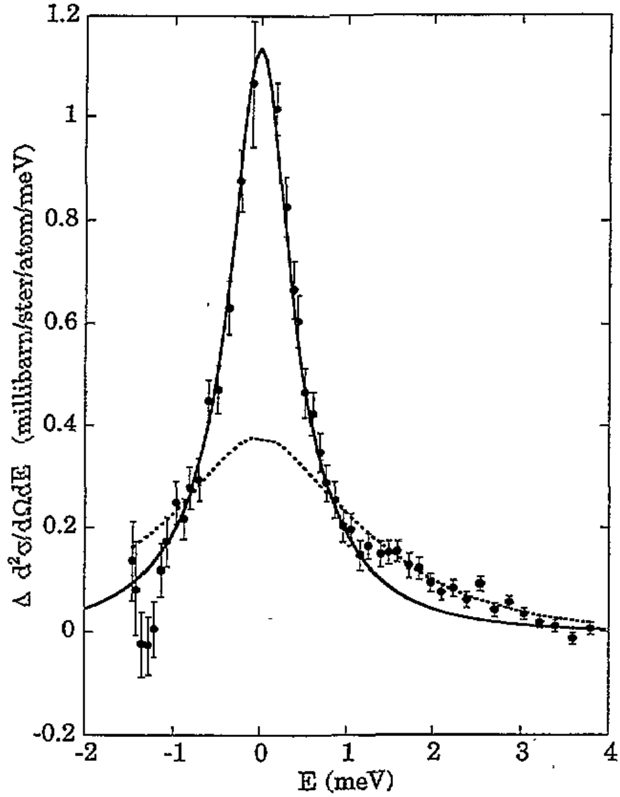

Figure 2. The difference cross section of figure 1 averaged over the experimental wavevector range $0.10\ \text{Å}^{-1} < q < 0.33\ \text{Å}^{-1}$. A single-pole model for $\text{Im}[\chi(q,\omega)]$ with unit weight ($z = 1$) and a q-independent relaxation frequency $\hbar\Gamma = 6$ meV (broken line), employed to describe the previous intermediate-frequency data, fails to account for the low-frequency part of the cross section. The solid line shows for contrast the prediction for $z \simeq 0.18$ and $\hbar\Gamma(q) \simeq (2.5\ \text{meV Å})q$. (Note that in both cases the cross section includes the thermal attenuation factor as in (3).)

description of $\text{Im}[\chi(q,\omega)]$ at low q, ω in terms of a low-frequency Fermi-liquid or 'slow' component plus an intermediate-frequency non-Fermi-liquid or 'fast' component. To arrive at this picture we begin with two initially independent elements of the generalized susceptibility, $\alpha(q,\omega)$ and $\beta(\omega)$ as the *starting* slow q-dependent part and fast q-independent part, respectively. To obtain the final susceptibility, these are then coupled in a mean field manner via a molecular field parameter λ taken to be independent of q and ω. This intermediate-level description might be inferred formally from an examination of the low-frequency action [18] resulting from a successive integration of the high-frequency degrees of freedom down to the frequency scale of interest (i.e. in UPt_3, down to a frequency characteristic of the weakly q-dependent dominant component). The parameters defining $\alpha(q,\omega)$ and $\beta(\omega)$ at this level of description may be already strongly renormalized relative to those of the starting non-interacting system.

The susceptibility of the composite system in the simplest mean field approximation is then

$$\chi(q,\omega) = \frac{\alpha(q,\omega) + \beta(\omega) + 2\lambda\alpha(q,\omega)\beta(\omega)}{1 - \lambda^2\alpha(q,\omega)\beta(\omega)}. \tag{8}$$

For sufficiently low q the low-ω part of (8) (i.e. the renormalized slow part for an assumed hypothetical homogeneous system) may be obtained by replacing $\beta(\omega)$ by $\beta = \beta(\omega = 0)$ and employing the Fermi-liquid expansion

$$\alpha(q,\omega)^{-1} = \alpha^{-1} - \mathrm{i}\omega/(\gamma_0 q) \tag{9}$$

where $\alpha = \alpha(0,0)$ and γ_0 is the coefficient defining the initial relaxation rate $\gamma_0\alpha^{-1}q$. Then $\mathrm{Im}[\chi(q,\omega)]$ reduces to (4) and (5) with χ, z and γ defined as

$$\chi = \frac{\alpha + \beta + 2\lambda\alpha\beta}{1 - \lambda^2\alpha\beta} \tag{10}$$

$$z = 1 - \beta/\chi \tag{11}$$

and

$$\gamma = \gamma_0(\alpha + \beta + 2\lambda\alpha\beta)/\alpha. \tag{12}$$

The renormalized coefficient γ can be smaller than γ_0 if the coupling parameter λ is negative (antiferromagnetic). Taking $\hbar\gamma_0 \simeq 2.5$ μeV Å as in the d systems, for example, a good description of the UPt_3 spectra is obtained for $\lambda < 0$ with the starting parameters defined by $\alpha \simeq \chi/2$, $\beta/\alpha \simeq 1.6$ and $(1 - \lambda^2\alpha\beta)^{-1} \simeq 10$. For an overall description of both the slow and the fast components the ω dependence of $\beta(\omega)$ must be included. It may be modelled by a function $\beta(\omega)$ whose imaginary part has a broad peak at a characteristic frequency only slightly higher than the renormalized value of the order of 5 meV. To avoid overcounting the slow Fermi-liquid contribution represented by $\alpha(q,\omega)$, we require $\mathrm{Im}[\beta(\omega)/\omega] \to 0$ as $\omega \to 0$ [19].

A two-component model with some of the above features has been presented and described in microscopic terms in [18]. In the latter work the slow (intermediate 'fermion') component is small due to the assumed weak overlap between the 'fermion' and bare electron states and, before mean field coupling is carried out, the susceptibility is dominated by the fast (intermediate 'spin fluctuation' in the language of [18]) contribution. Therefore, terms other than $\beta(\omega)$ in the numerator of the right-hand side of (8) could in this case be neglected. Self-interactions within the 'fermion' system itself may in practice, however, lead to a smaller difference between α and β than this model assumed. We also note that in a more realistic treatment of a two-component intermediate model for a lattice it is also necessary to take account of a mean field feedback from components of the magnetization induced at q' separated from q by a reciprocal-lattice wavevector. This effectively alters the form of $\alpha(q,\omega)$ to be used particularly in the term $\lambda^2\alpha(q,\omega)\beta(\omega)$ from that defined in (8). A more consistent treatment appropriate for a lattice requires the introduction of yet more parameters which cannot at this time be independently determined.

6. Conclusions

Our key finding is that the low-q, ω form of $\mathrm{Im}[\chi(q,\omega)]$ in UPt_3 cannot be understood in *terms of a single-pole model with a wavevector-independent relaxation spectrum*. A minimal description consistent with firstly the present measurement at low energies (determined by the thermal cut-off of about 1 meV) and secondly previous studies [10, 11] at intermediate energies (below about 10 meV) may be based on a 'slow' component with a dispersive relaxation rate and a 'fast' contribution with a weakly wavevector-dependent characteristic frequency. The 'slow' component in $\mathrm{Im}[\chi(q,\omega)]$ accounts, via the Kramers–Kronig relation, for approximately 20% of the total static susceptibility and has a relaxation spectrum not inconsistent with that expected within an elementary Fermi-liquid model. This interpretation is not as unambiguous as that obtained in d transition metals in which a well defined quasi-particle component with a strongly dispersive relaxation spectrum consistent with microscopic Fermi-liquid theory accounts for essentially all the static susceptibility as $q \to 0$. The difference between the behaviours of the f systems may be attributed in part to the strength of the spin–orbit interaction which leads to the existence of a strong and dominant non-Fermi-liquid component even in the limit of long wavelengths.

Acknowledgments

This work was carried out on the IN5 spectrometer at the Institut Laue–Langevin in Grenoble with the support of G Kearley. We also thank D McK Paul, S M Hayden and E J Lindley for their assistance, and G Aeppli, P Coleman, J Flouquet and G A Gehring for valuable discussions. One of us (NRB) is also grateful to J Flouquet for his hospitality and support during the course of this work which has stimulated further measurements jointly with G Aeppli, S M Hayden and A Huxley. The sample was produced in Cambridge in collaboration with L Taillefer with the support of the Science and Engineering Research Council of the UK and has also been employed in the studies described under [12, 13].

References

[1] See, e.g. Baym G and Pethick C 1991 *Landau Fermi-Liquid Theory* (New York: Wiley)
[2] Scherm R, Guckelsberger K, Fak B, Skold K, Dianoux A J, Godfrin H and Stirling W G 1987 *Phys. Rev. Lett.* **59** 217
[3] Bernhoeft N R, Hayden S M, Lonzarich G G, Paul D and Lindley E J 1989 *Phys. Rev. Lett.* **62** 657
[4] Bernhoeft N R *et al* 1995 to be published
[5] Ishikawa Y, Noda Y, Fincher C and Shirane G 1982 *Phys. Rev.* B **25** 25
[6] Ishikawa Y, Noda Y, Uemura Y J, Majkrzak C F and Shirane G 1985 *Phys. Rev.* B **31** 5884
[7] Bernhoeft N R, Lonzarich G G, Mitchell P W and McK Paul D 1983 *Phys. Rev.* B **28** 422
[8] Bernhoeft N R, Law S A, Lonzarich G G and McK Paul D 1988 *Phys. Scr.* **38** 191
[9] Holland-Moritz E and Lander G H, Aeppli G and Broholm C 1994 *Handbook on the Physics and Chemistry of Rare Earths* vol 19, ed K A Gschneider Jr, G H Lander and G R Choppin (Amsterdam: Elsevier)
[10] Aeppli G, Bucher E and Shirane G 1985 *Phys. Rev.* B **32** 7579
[11] Aeppli G, Goldman A I, Shirane G, Bucher E and Lux-Steiner M 1987 *Phys. Rev. Lett.* **58** 808
[12] Taillefer L and Lonzarich G G 1988 *Phys. Rev. Lett.* **60** 1570
[13] See, e.g., Taillefer L, Flouquet J and Lonzarich G G 1991 *Physica* B **169** 257
[14] See, e.g., Lovesey S W 1984 *Theory of Neutron Scattering from Condensed Matter* (Oxford: Clarendon)
[15] Moriya T 1985 *Spin Fluctuations in Itinerant Electron Magnetism* (Berlin: Springer) and references therein
[16] Lonzarich G G and Taillefer L 1985 *J. Phys. C: Solid State Phys.* **18** 4339
[17] Franse J J M, Frings P H, de Visser A, Menovsky A, Palstra T T M, Kes P H and Mydosh J A 1984 *Physica* B **126** 116
[18] Kuramoto Y and Miyake K 1990 *J. Phys. Soc. Japan* **59** 2831
[19] Formally we note that the interpretation of the *high*-energy fluctuations observed in various heavy-electron systems [9] may also be included in our analysis on setting $\alpha \ll \beta$ and by ignoring the frequency dependence of α and incorporating a lowest-order local relaxation approximation to $\beta(\omega)$