

Novel neutron resonance mode in $d_{x^2-y^2}$ superconductors

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We show that a new resonant magnetic excitation at incommensurate momenta, observed recently by inelastic neutron scattering experiments on $\text{YBa}_2\text{Cu}_3\text{O}_{6.85}$ and $\text{YBa}_2\text{Cu}_3\text{O}_{6.6}$, is a *spin exciton*. Its location in the magnetic Brillouin zone and its frequency are determined by the momentum dependence of the particle-hole continuum. We identify several features that distinguish this novel mode from the previous resonance mode observed near $\mathbf{Q} = (\pi, \pi)$, such as its intensity maximum which occurs in a different part of the magnetic Brillouin zone.

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While it seems established that the resonance peak is a universal feature of the high-temperature superconductors [1, 2, 3], its origin, its role for the pairing process, and the effects arising from its interactions with electrons are still intensively debated [4]. The peak's intensity is the highest at $\mathbf{Q} = (\pi, \pi)$ and in $\text{YBa}_2\text{Cu}_3\text{O}_{6+x}$ (YBCO), where it was studied in great detail, its frequency $\Omega_{res}(\mathbf{Q})$ follows the same doping dependence as T_c , with $\Omega_{res}(\mathbf{Q}) \approx 41$ meV near optimal doping. As one moves away from \mathbf{Q} , the peak disperses downwards and its intensity decreases rapidly, vanishing around $\mathbf{Q}_0 = (0.8\pi, 0.8\pi)$. The doping dependence of $\Omega_{res}(\mathbf{Q})$, the downward dispersion of the resonance, and the fact that \mathbf{Q}_0 coincides with the distance between nodal (diagonal) points on the Fermi surface are consistent with the theoretical idea that the resonance peak is a particle-hole bound state below the spin gap (a spin exciton) [5] (for a review of other theoretical scenarios, see Ref. [6]).

Recent inelastic neutron scattering (INS) experiments on $\text{YBa}_2\text{Cu}_3\text{O}_{6.85}$ [7] in the superconducting (SC) state detected a new resonant magnetic excitation at incommensurate momenta, but at frequencies *larger* than $\Omega_{res}(\mathbf{Q})$. A similar result was obtained for underdoped $\text{YBa}_2\text{Cu}_3\text{O}_{6.6}$ [8]. This new resonance mode is particularly pronounced along the diagonal of the magnetic Brillouin zone (MBZ) at $\mathbf{q} \lesssim \mathbf{Q}_0$. It was suggested [7] that this new resonance is a particle-hole bound state with an upward dispersion originating at \mathbf{Q} (see Fig. 5a in Ref. [7]).

In this Letter, we show both analytically and numerically that the new resonance mode is indeed a spin exciton that emerges below T_c due to a feedback effect on the collective spin excitations arising from the opening of the SC gap in a $d_{x^2-y^2}$ superconductor. We demonstrate that the new resonance appears only at momenta less than \mathbf{Q}_0 , and is separated from the previous resonance by a region near \mathbf{Q}_0 in which no resonance exists (the “silent band” of Ref. [7]). Thus, the new resonance does *not* form an upward dispersing branch originating

at \mathbf{Q} . We identify several qualitative features that distinguish this new resonance (the Q^* mode) from the old one (the Q mode). In particular, we show that while the intensity of the Q mode is largest along $\mathbf{q} = (\pi, \eta\pi)$ and $\mathbf{q} = (\eta\pi, \pi)$, the Q^* mode has its largest intensity along the MBZ diagonal.

We begin by presenting our numerical analysis of the new Q^* mode. Its emergence can be understood within an RPA approach for which the spin susceptibility is given by

$$\chi(\mathbf{q}, \omega) = \frac{\chi_0(\mathbf{q}, \omega)}{1 - g(\mathbf{q})\chi_0(\mathbf{q}, \omega)} \quad (1)$$

where $g(\mathbf{q})$ is the fermionic four-point vertex, and $\chi_0(\mathbf{q}, \omega)$ is the free-fermion susceptibility, which in the SC state is given by the sum of two single bubble diagrams consisting of either normal or anomalous Greens functions [5, 9]. For our numerical calculation of $\chi_0(\mathbf{q}, \omega)$, we used a SC gap with $d_{x^2-y^2}$ symmetry and a normal state tight binding dispersion

$$\epsilon_{\mathbf{k}} = -2t(\cos k_x + \cos k_y) - 4t' \cos k_x \cos k_y - \mu \quad (2)$$

with $t = 250$ meV, $t'/t = -0.4$, and $\mu/t = -1.083$. The Fermi surface (FS) obtained from Eq. (2) is shown in Fig. 3(a). It describes well the FS measured by photoemission experiments on $\text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_{8+\delta}$ [10].

Our main results are presented in Figs. 1 and 2, in which we plot $\text{Im}\chi(\mathbf{q}, \omega)$, obtained from a numerical evaluation of Eq.(1), along $\mathbf{q} = \eta(\pi, \pi)$ in the SC state. The intensity plot of $\text{Im}\chi(\mathbf{q}, \omega)$ shown in Fig. 1 possesses all salient features observed in the INS experiments. First, we identify a downward dispersion of the Q mode (indicated by a red arrow). Second, the Q^* mode (indicated by a yellow arrow) is located at frequencies larger than the frequency of the Q mode at (π, π) , and is confined to a small region of momentum space near \mathbf{Q}_0 . The momentum position of the Q^* mode is almost independent of energy. Third, there exists a region in momentum space around \mathbf{Q}_0 that separates the Q mode from the Q^*

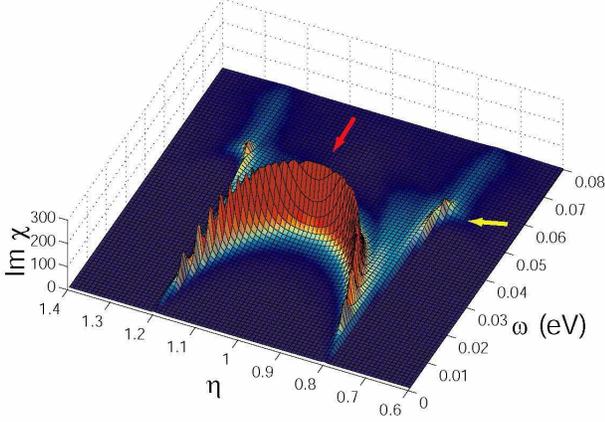


FIG. 1: (color) RPA results for magnetic excitations in a $d_{x^2-y^2}$ superconductor. $\text{Im}\chi$ obtained from Eq. (1) as a function of momentum (along $\mathbf{q} = \eta(\pi, \pi)$) and frequency in the SC state. We used $\Delta_{\mathbf{k}} = \Delta_0(\cos k_x - \cos k_y)/2$ with $\Delta_0 = 42$ meV, and $g(\mathbf{q}) = g_0[1 - 0.1(\cos q_x + \cos q_y)]$, with $g_0 = 0.573$ eV in order to reproduce the correct energy position of the Q^* mode near $0.8(\pi, \pi)$ and the Q mode at (π, π) .

one (the “silent band” of Ref. [7]). Note, that as \mathbf{Q}_0 is approached from (π, π) , the Q mode frequency, $\Omega_{res}(\mathbf{q})$, as well as its intensity, rapidly decreases. In Fig. 2 we present $\text{Im}\chi$ along $\mathbf{q} = \eta(\pi, \pi)$ for several frequencies. We clearly see that the two modes are separated in momentum and frequency space. The shaded area in Fig. 2 represents the silent band, in which $\text{Im}\chi$ is strongly reduced from its resonance values. We also find that the position of the Q^* mode is almost frequency independent, with a maximum intensity at $\omega \approx 54$ meV, in agreement with the experimental observations (see Fig. 2a of Ref. [7]).

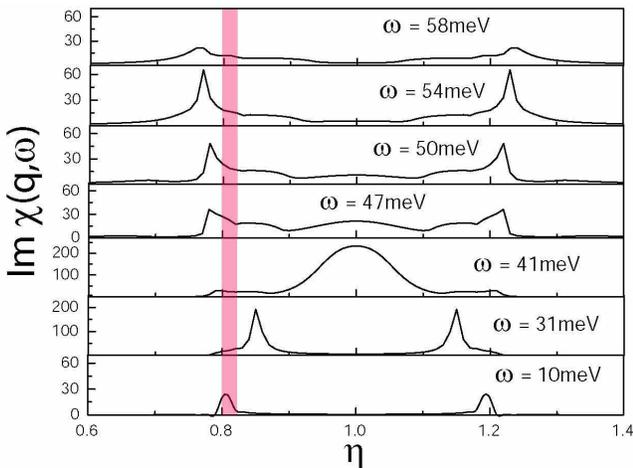


FIG. 2: $\text{Im}\chi(\mathbf{q}, \omega)$ along $\mathbf{q} = \eta(\pi, \pi)$ for several frequencies. The position of the silent band is indicated by the shaded area. These results reproduce the experimental INS data, see Fig. 2(a) in Ref. [7].

In what follows, we discuss the physical origin of our results. We begin by reviewing the emergence of the Q mode. Its origin lies in a feedback effect on the spin excitation spectrum (i.e., on χ_0) due to the opening of the SC gap. This feedback effect is universal in that it only requires that the FS possesses points, \mathbf{k} and \mathbf{k}' , that can be connected by \mathbf{Q} (so-called “hot spots”) and that $\text{sgn}(\Delta_{\mathbf{k}}) = -\text{sgn}(\Delta_{\mathbf{k}'})$. To demonstrate the universality of the resonance mode, we rewrite the denominator of Eq. (1) as

$$\xi_{\mathbf{Q}}^{-2} - g(\mathbf{Q})\Delta\chi_0(\mathbf{Q}, \Omega) \quad (3)$$

where $\xi_{\mathbf{Q}}^{-2} = 1 - g(\mathbf{Q})\chi_0^{ns}(\mathbf{Q}, 0)$ and $\Delta\chi_0(\mathbf{Q}, \Omega) = \chi_0(\mathbf{Q}, \Omega) - \chi_0^{ns}(\mathbf{Q}, 0)$. The main contribution to the static normal state susceptibility $\chi_0^{ns}(\mathbf{Q}, 0)$ comes from high internal frequencies in the fermionic bubble. This implies that χ_0^{ns} and thus $\xi_{\mathbf{Q}}^{-2}$ are non-universal quantities that depend on the details of the band structure. However, the appearance of the resonance mode only requires $\xi_{\mathbf{Q}}^{-2} > 0$, which is satisfied in the paramagnetic state. In contrast, $\Delta\chi_0(\mathbf{Q}, \Omega)$ is *universal* since its main contribution comes from small internal frequencies of order Δ . $\Delta\chi_0$ can therefore be evaluated by simply linearizing the fermionic dispersion near \mathbf{k} and \mathbf{k}' . The integration over momentum [11] yields at $T = 0$

$$\Delta\chi_0(\mathbf{Q}, \Omega) = -i\frac{\gamma_{\mathbf{Q}}^{NS}}{16} \sum_{\{\mathbf{k}, \mathbf{k}'\}} \int_{-\infty}^{\infty} d\omega \times \left(1 - \frac{\omega_+ \omega_- + \Delta_{\mathbf{k}} \Delta_{\mathbf{k}'}}{\sqrt{\omega_+^2 - \Delta_{\mathbf{k}}^2} \sqrt{\omega_-^2 - \Delta_{\mathbf{k}'}^2}} \right) \quad (4)$$

where $\omega_{\pm} = \omega \pm \Omega/2$, and the summation runs over all pairs of FS points \mathbf{k} and \mathbf{k}' separated by \mathbf{Q} . There are eight fermionic scattering processes from \mathbf{k} to \mathbf{k}' in the first zone. The momenta \mathbf{k} involved are those in which the boundary of the magnetic zone (defined by $\cos k_x = -\cos k_y$) crosses the FS. Two of these scattering processes are direct with $\mathbf{k} - \mathbf{k}' = \mathbf{Q}$, two involve umklapp scattering with $\mathbf{k} - \mathbf{k}' = \mathbf{Q} - (2\pi, 2\pi)$, and four involve umklapp scattering with $\mathbf{k} - \mathbf{k}' = \mathbf{Q} - (2\pi, 0)$ and $\mathbf{k} - \mathbf{k}' = \mathbf{Q} - (0, 2\pi)$. A pair of FS points that are connected by direct scattering via \mathbf{Q} is shown in Fig. 3(a) (\mathbf{Q} is represented by a dashed arrow).

In the normal state, all eight processes equally contribute to $\Delta\chi_0(\mathbf{Q}, \Omega)$, and one obtains from Eq. (4) $\Delta\chi_0(\mathbf{Q}, \Omega) = -i\gamma_{\mathbf{Q}}^{NS}\Omega$ which identifies $\gamma_{\mathbf{Q}}^{NS}$ with the Landau damping rate [12]. In the superconducting state, $\Delta\chi_0$ becomes a complex function. According to Eq. (4), its imaginary part vanishes below a critical frequency $\Omega_c(\mathbf{Q}) = |\Delta_{\mathbf{k}}| + |\Delta_{\mathbf{k}'}|$, which is the same for all eight scattering channels. The $d_{x^2-y^2}$ symmetry of the SC gap implies $\Delta_{\mathbf{k}} = -\Delta_{\mathbf{k}'}$, resulting in a discontinuous jump of $\text{Im}\chi_0$ at $\Omega_c(\mathbf{Q})$ from zero to $\pi\gamma_{\mathbf{Q}}^{NS}|\Delta_{\mathbf{k}}|$ [5]. Simultaneously, $\text{Re}\Delta\chi_0(\mathbf{Q}, \Omega) > 0$ is non-zero, diverges logarithmically at $\Omega_c(\mathbf{Q})$, and scales as Ω^2 at small frequencies.

$\text{Re}\Delta\chi_0$ therefore varies between 0 at $\Omega = 0$ and ∞ at $\Omega = \Omega_c$. Since $\xi_{\mathbf{Q}}^{-2} > 0$, one finds that for any positive $g_{\mathbf{Q}}$, $\chi(\mathbf{Q}, \Omega)$ [Eq. (1)] acquires a pole at a frequency $\Omega_{res} < \Omega_c(\mathbf{Q})$ where $\text{Re}\Delta\chi_0(\mathbf{Q}, \Omega_{res}) = \xi_{\mathbf{Q}}^{-2}/g_{\mathbf{Q}}$ and $\text{Im}\Delta\chi_0(\mathbf{Q}, \Omega_{res}) = 0$. $\text{Im}\chi$ thus exhibits a δ -function at Ω_{res} , representing a spin exciton below the particle-hole (p-h) continuum.

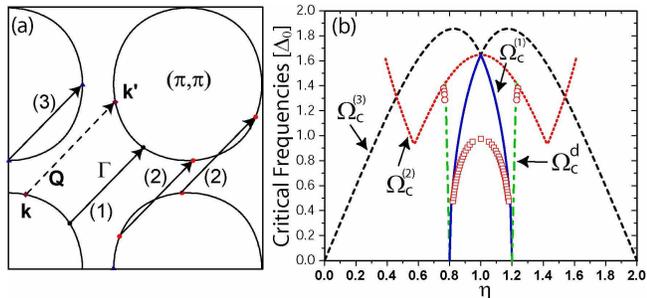


FIG. 3: (a) FS of Eq. (2) and magnetic scattering vectors. (b) Momentum dependence of the critical frequencies $\Omega_c^{(i)}$ ($i = 1, 2, 3$) and the direct gap Ω_c^d along $\mathbf{q} = \eta(\pi, \pi)$ (see text). The open squares (circles) represent the position of the Q (Q^*) resonance.

We next consider the dispersion of the exciton. For $\mathbf{q} \neq \mathbf{Q}$, the degeneracy of the scattering channels is lifted. In particular, for momenta along $\mathbf{q} = \eta(\pi, \pi)$, one now has three different critical frequencies, $\Omega_c^{(i)}(\mathbf{q})$ ($i = 1, 2, 3$) [5]. $\Omega_c^{(1)}$ is the critical frequency associated with direct scattering, $\Omega_c^{(2)}$ with umklapp scattering involving $\mathbf{q} - (2\pi, 0)$ and $\mathbf{q} - (0, 2\pi)$, and $\Omega_c^{(3)}$ with umklapp scattering by $\mathbf{q} - (2\pi, 2\pi)$. The scattering momenta for \mathbf{q} close to \mathbf{Q}_0 are shown in Fig. 3(a), and $\Omega_c^{(i)}(\mathbf{q})$ for $\mathbf{q} = \eta(\pi, \pi)$ are presented in Fig. 3(b). For all of these scattering processes, $\Delta_{\mathbf{k}}$ and $\Delta_{\mathbf{k}'}$ still have opposite signs [13]. As a result, $\text{Im}\Delta\chi_0$ exhibits three discontinuous jumps at $\Omega_c^{(i)}(\mathbf{q})$, and $\text{Re}\Delta\chi_0$ diverges logarithmically at $\Omega_c^{(i)}(\mathbf{q})$ [5]. However, $\text{Im}\Delta\chi_0$ is zero only below the smallest $\Omega_c^{(i)}(\mathbf{q})$, and hence a true resonance is only possible below the smallest critical frequency. The splitting of the critical frequencies for $\mathbf{q} \neq \mathbf{Q}$ was also discussed in Ref. [14].

It follows from Fig. 3(b) that for $0.8 < \eta < 1.2$, the smallest critical frequency, $\Omega_c^{(1)}$, corresponds to direct (i.e., non-umklapp) scattering. $\Omega_c^{(1)}$ decreases away from $\eta = 1$ and eventually vanishes at $\eta = 0.8$, when direct scattering occurs between nodal points at the Fermi surface [5]. Since the exciton is necessarily located below $\Omega_c^{(1)}$, its frequency also decreases and eventually vanishes at $\eta = 0.8$. Moreover, upon approaching $\eta = 0.8$ the jump in $\text{Im}\Delta\chi_0(\mathbf{q}, \Omega_c^{(1)})$ decreases as \mathbf{k} and \mathbf{k}' approach the nodal points. Accordingly, the resonance frequency Ω_{res} moves closer to $\Omega_c^{(1)}$, and the intensity of the resonance decreases [5].

At $\mathbf{Q}_0 = 0.8(\pi, \pi)$, $\text{Im}\chi_0(\mathbf{Q}_0, \Omega)$ is non-zero for $\Omega > 0$, and one finds $\text{Im}\chi(\mathbf{Q}_0, \Omega) = \gamma_{\mathbf{Q}_0}^{SC}\Omega$ at small frequencies, where $\gamma_{\mathbf{Q}_0}^{SC} = \gamma_{\mathbf{Q}}^{NS} \frac{1}{8} \frac{\pi v_F}{4v_{\Delta}}$ and v_{Δ} is the gap velocity at the nodal points. The factor $1/8$ arises since only a single (direct) scattering channel contributes to $\gamma_{\mathbf{Q}_0}^{SC}$, while eight channels contribute to $\gamma_{\mathbf{Q}}^{NS}$. However, since the Fermi velocities at the nodal points are antiparallel, $\gamma_{\mathbf{Q}_0}^{SC}$ depends on v_F only through the ratio $v_F/v_{\Delta} \sim 20$ which compensates the small prefactor [15]. As a result, $\gamma_{\mathbf{Q}_0}^{SC}$ is comparable to $\gamma_{\mathbf{Q}}^{NS}$, thus giving rise to a weak and featureless frequency dependence of $\text{Im}\chi(\mathbf{Q}_0, \Omega)$ similar to that of $\text{Im}\chi_0(\mathbf{Q}, \Omega)$ in the normal state. The vanishing of the gap in the p-h continuum at \mathbf{Q}_0 together with the large value of $\gamma_{\mathbf{Q}_0}^{SC}$ explains the experimental observation of a “silent band” in Ref. [7] (the position of which is indicated by the shaded area in Fig. 2).

For momenta $\mathbf{q} < \mathbf{Q}_0$, (i.e., $\eta < 0.8$) the nodal points cannot be connected, and a *direct* gap opens for excitations into the p-h continuum. This gap is independent of the SC gap, and given by $\Omega_c^d = v_F \cdot (\mathbf{Q}_0 - \mathbf{q})$ (see dashed-dotted line in Fig. 3(b)). Due to a large $|v_F|$, Ω_c^d becomes equal to $\Omega_c^{(2)}$ close to \mathbf{Q}_0 at $\mathbf{q}_d = \eta_d \mathbf{Q}$ with $\eta_d = 0.773$ (i.e., $\mathbf{Q}_0 - \mathbf{q}_d = 0.027\mathbf{Q}$ for the dispersion of Eq. (2)). For $\eta_d < \eta < 0.8$, one finds that $\text{Im}\chi_0$ vanishes below Ω_c^d , and $\text{Im}\chi_0 \sim \sqrt{\Omega - \Omega_c^d}$ for $\Omega > \Omega_c^d$. Hence, $\text{Re}\chi_0$ does not diverge at Ω_c^d , and no resonance peak exists in this region, extending the silent band. However, $\text{Re}\chi_0$ possesses a logarithmic divergence at $\Omega_c^{(2)}$, and hence it satisfies the resonance condition $\text{Re}\Delta\chi_0(\mathbf{q}, \bar{\Omega}_{res}) = \xi_{\mathbf{q}}^{-2}/g_{\mathbf{q}}$ at some frequency $\bar{\Omega}_{res}$ below $\Omega_c^{(2)}$. Once Ω_c^d crosses $\bar{\Omega}_{res}$, the damping at Ω_{res} vanishes and a true pole in $\text{Im}\chi$ occurs, leading to the appearance of the Q^* mode (open circles in Fig. 3(b)). As one moves further away from \mathbf{Q}_0 , one finds that the Q^* resonance is rapidly suppressed. This suppression arises from the rapid decrease of $\Omega_c^{(2)}$ as well as the decrease of the bare static spin susceptibility, $\chi_0(\mathbf{q}, 0)$. This behavior of $\chi_0(\mathbf{q}, 0)$ is similar to that in a system of free fermions with a circular Fermi surface, if we identify \mathbf{Q}_0 with $2\mathbf{p}_F$. Both effects lead to a rapid shift of the Q^* mode towards the edge of the p-h continuum, and to a decrease in its intensity. As a result, this resonance is only visible near \mathbf{Q}_0 .

The Q and Q^* modes are not only separated in frequency, as discussed above, but their intensity maxima are also located in different parts of the MBZ; this represents a major qualitative distinction between the two modes. In Fig. 4 we present intensity plots of $\text{Im}\chi$ as a function of momentum for $\Omega > \Omega_{res}(\mathbf{Q})$, probing the Q^* mode [Fig. 4(b)], and for $\Omega < \Omega_{res}(\mathbf{Q})$, probing the Q mode [Fig. 4(a)]. The difference is striking. While the intensity of the Q mode is largest along $\mathbf{q} = (\pi, \eta\pi)$ and $\mathbf{q} = (\eta\pi, \pi)$, the Q^* mode has its largest intensity along the diagonal direction, i.e., along $\mathbf{q} = \eta(\pi, \pi)$ and $\mathbf{q} = [(2 - \eta)\pi, \eta\pi]$. This rotation of the intensity pattern

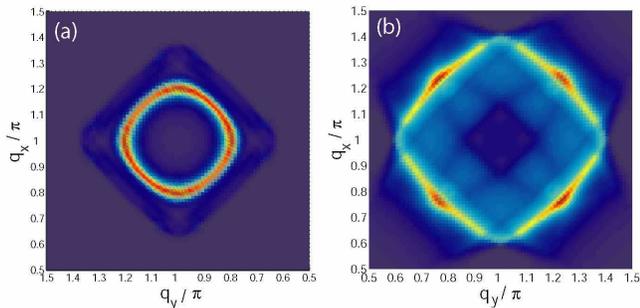


FIG. 4: (color) Intensity plot of $\text{Im}\chi$, Eq. (1), as a function of momentum (blue (red) implies a small (large) value of $\text{Im}\chi$) at: (a) 33 meV (Q mode), and (b) 57 meV (Q^* mode). The Q mode forms a distorted ring in momentum space, with intensity maxima along (π, q) . In contrast, the intensity of the Q^* mode is largest along the zone diagonal. In (b), the periodic structure of $\text{Im}\chi$ around \mathbf{Q} reflects the momentum dependence of the p-h continuum at this energy.

by 45° reflects the qualitative difference in the origin of the two modes. The intensity of the Q mode is at a maximum along $\mathbf{q} = (\pi, \eta\pi)$ and $\mathbf{q} = (\eta\pi, \pi)$, since in this case the fermions that are scattered by \mathbf{q} are located farther from the nodes than for diagonal scattering. In contrast, the Q^* mode arises from the rapid opening of a gap in the p-h continuum below \mathbf{Q}_0 , which is most pronounced along the diagonal directions of the zone. The rotation of the intensity pattern by 45° , which is the most significant distinction between the Q and Q^* modes, was observed in recent INS experiments [8, 16].

In summary, we showed that a new resonant magnetic excitation at incommensurate momenta, observed recently by inelastic neutron scattering experiments on $\text{YBa}_2\text{Cu}_3\text{O}_{6.85}$ [7] and $\text{YBa}_2\text{Cu}_3\text{O}_{6.6}$ [8], is a spin excitation arising from umklapp scattering. Its location in the zone and its frequency are determined by the momentum dependence of the p-h continuum. It is confined to a small region in \mathbf{q} just below \mathbf{Q}_0 , and is separated from the Q resonance by a “silent band” where $\text{Im}\chi$ is strongly suppressed. We also found that the intensity maxima of the two modes are rotated by 45° relative to each other.

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 - [16] The underdoped INS data reveal a pattern in momentum space which is more spot like than the optimal doped data, both for the Q and Q^* modes. This effect can be reproduced in the RPA calculations if the Fermi surface is flattened near the node.