Inelastic scattering of light by magnetic excitons in the pseudo Ising antiferromagnets K2CoF4 and Rb2CoF4


To cite this version:


HAL Id: jpa-00209193
https://hal.archives-ouvertes.fr/jpa-00209193
Submitted on 1 Jan 1979

HAL is a multi-disciplinary open access archive for the deposit and dissemination of scientific research documents, whether they are published or not. The documents may come from teaching and research institutions in France or abroad, or from public or private research centers.

L’archive ouverte pluridisciplinaire HAL, est destinée au dépôt et à la diffusion de documents scientifiques de niveau recherche, publiés ou non, émanant des établissements d’enseignement et de recherche français ou étrangers, des laboratoires publics ou privés.
Inelastic scattering of light by magnetic excitons in the pseudo Ising antiferromagnets $K_2CoF_4$ and $Rb_2CoF_4$

J.-P. Gosso and P. Moch
C.N.R.S., Laboratoire P.M.T.M., Université Paris-Nord, avenue J.-B.-Clément, 93430 Villetaneuse, France

M. Quilichini
D.R.P., C.N.R.S., LA 71, Université Pierre-et-Marie-Curie, 4, place Jussieu, 75230 Paris Cedex 05, France

J.-Y. Gesland and J. Nouet
Laboratoire de Physique de l'État Condensé, E.R.A. 682, C.N.R.S., Faculté des Sciences, Route de Laval, 72017 Le Mans Cedex, France

(Reçu le 28 mai 1979, accepté le 12 juillet 1979)

Résumé. — Outre l'effet Raman dû aux phonons, antérieurement signalé, nous observons la diffusion inélastique de la lumière liée aux excitons magnétiques dans $K_2CoF_4$ et $Rb_2CoF_4$ ; les spectres polarisés ont été étudiés à basse température en présence d'un champ magnétique pouvant atteindre 5 teslas environ. Nous interprétons les résultats au moyen d'un modèle multi-excitonique déduit d'un Hamiltonien tenant compte du couplage spin-orbite, du champ cristallin de basse symétrie et d'une interaction d'échange de type Heisenberg, perturbant l'état fondamental $4f^4_4^+(O_h)$ des ions $Co^{2+}$. De l'étude de l'exciton de plus basse énergie, nous déduisons la valeur du champ cristallin quadratique $0_394$ ($425 \text{ cm}^{-1}$ et $635 \text{ cm}^{-1}$, respectivement pour $K_2CoF_4$ et $Rb_2CoF_4$) ainsi que celle du paramètre d'échange isotrope $J$ ($11.55 \text{ cm}^{-1}$ et $9.65 \text{ cm}^{-1}$, respectivement pour $K_2CoF_4$ et $Rb_2CoF_4$). Le reste du spectre, qui contient des contributions à la diffusion par un et par deux excitons, est partiellement interprété, grâce au modèle ci-dessus décrit, pour les valeurs ainsi déterminées des paramètres $A$ et $J$, dont il dépend. Les valeurs de $J$ sont voisines de celles qui ont été antérieurement trouvées pour $KCoF_3$ et $RbCoF_3$, comme prévu. Le comportement de type Ising, qui provient du champ cristallin de basse symétrie, apparaît moins marqué lorsqu'on l'évalue à partir de la courbe de dispersion des magnons que ne le laissaient prévoir d'anciennes mesures magnétiques, résultat confirmé par une récente étude de diffusion inélastique de neutrons dans $Rb_2CoF_4$.

Abstract. — In addition to previously reported phonon Raman scattering, we have observed inelastic light scattering by magnetic excitons in $K_2CoF_4$ and $Rb_2CoF_4$; the polarized spectra have been studied at low temperatures under an applied magnetic field up to about 5 teslas. The results are interpreted on the basis of a multilevel propagating exciton model deduced from a Hamiltonian including the single ion spin-orbit coupling, the low symmetry crystal field and a nearest neighbours Heisenberg exchange interaction, acting in the $4f^4_4^+(O_h)$ ground state of $Co^{2+}$. From the lowest exciton study we derive a value of the tetragonal crystal field $A$ ($425 \text{ cm}^{-1}$ and $635 \text{ cm}^{-1}$, respectively in $K_2CoF_4$ and $Rb_2CoF_4$) and of the isotropic exchange parameter $J$ ($11.55 \text{ cm}^{-1}$ and $9.65 \text{ cm}^{-1}$, respectively in $K_2CoF_4$ and $Rb_2CoF_4$). The remaining features of the spectra, which include one and two exciton scattering, are partly interpreted using the above model with these values. The derived $J$ values are close to previous determinations in $KCoF_3$ and $RbCoF_3$, as expected. The Ising character, which is caused by the low symmetry crystal field, is found to be less marked when evaluated from the magnon dispersion results than it was indicated by previous magnetic measurements, and this has been confirmed by recent inelastic neutron scattering data in $Rb_2CoF_4$.

1. Introduction. — The study of magnetic excitations in cobaltous antiferromagnets has given rise to an abundant but rather controversial literature [1, 11]. The unquenching of the orbital momentum in the $4f^4_4^+$ ground state of $Co^{2+}$ provides a number of magnetic excitons, the description of which depends
upon the spin-orbit coupling, the low symmetry crystal field and the exchange interaction. Due to the orbital degeneracy this last term is expected to show a complex behaviour: various authors [12, 13] have proposed rather general forms for the exchange interactions including orbital effects; a lot of independent constants are necessary, which considerably complicates the calculation of the magnetic excitons and hinders any convincing fit to the experimental data. On the other hand, several authors claim that a rather simple Heisenberg or Ising model provides a good agreement with experimental data [1, 2, 8, 11], while others [3, 10] find that such a model is only applicable to the lowest energy exciton.

In order to help to resolve this controversy we have performed electronic Raman scattering on K$_2$CoF$_4$ and Rb$_2$CoF$_4$. These crystals are generally regarded as good examples of two-dimensional Ising antiferromagnets: such an assumption is mainly supported by magnetic susceptibility measurements [14, 15], which are essentially sensitive to the behaviour of the lowest Kramers doublet and do not drastically depend upon the precise form of the exchange interaction within the $^4T_a^+$ sublevels manifold. On the other hand, a direct spectroscopic investigation of several magnetic excitons is a rather severe test of the previously proposed exchange models [14, 16] based on a simple anisotropic Heisenberg Hamiltonian within $^4T_a^+$.

In the next section we briefly discuss the polarization selection rules for Raman scattering, derived from symmetry arguments, and we present the self-consistent excitonic models used. To a large extent, the symmetries of the magnetic excitons are independent of the models, but useful compatibility relations can be made, which help to evaluate their degree of validity. In section 3 we present the experimental data, which are discussed and interpreted in section 4.

2. Symmetries and excitonic models. — 2.1 Symmetries and selection rules. — In the paramagnetic state ($T > T_N$), K$_2$CoF$_4$ and Rb$_2$CoF$_4$ crystallize in the D$_{4h}$ system. Below $T_N$ (respectively 107 and 101 K in K$_2$CoF$_4$ and Rb$_2$CoF$_4$), they antiferromagnetically order with a sublattice magnetization parallel to Z (fig. 1), their crystal structure [15, 17, 18] being described by a [D$_{2h}$ + (θ, t$_0$) D$_{2h}$] type IV Shubnikov space group: C$_h$ ccm (in the preceding notation, θ is the time inversion operator and t$_0$ a non primitive translation). Strictly speaking the fourfold symmetry around Z is destroyed; however the crystal consists of very weakly magnetically coupled plane layers perpendicular to Z, which means that, practically, the fourfold symmetry does not disappear and that the crystals can be viewed as planar antiferromagnets described by a [C$_{4v}$ + (θ, t$_0$) C$_{4v}$] type IV Shubnikov space group.

The Co$^{2+}$ site unitary symmetry group changes from D$_{4h}$ to C$_{2h}$ when the crystals order antiferromagnetically but, in the above more realistic planar approximation, C$_{2h}$ has to be replaced by C$_4$. It is useful to introduce the C$_{4h}$ site symmetry which would be involved in a mean field model (mean field parallel to Z) since our effectively used single ion Hamiltonian remains invariant under the C$_{4h}$ symmetry operations (see below).

The $^4T_a^+$ ground state (in the cubic O$_h$ symmetry approximation) is split into 6 Kramers doublets under the combined effect of the low symmetry tetragonal crystal field and spin-orbit coupling, in the paramagnetic phase where the Co$^{2+}$ site symmetry is D$_{4h}$: the Kramers doublets belong to [19] $^4T_a^+$ (D$_{4h}$) or $^4T_b^+$ (D$_{4h}$). The mean field then splits each doublet into either $^4T_a^+$ (C$_{4h}$) and $^4T_b^+$ (C$_{4h}$) or $^4T_b^+$ (C$_{4h}$) and $^4T_a^+$ (C$_{4h}$) states. The lowering of the symmetry to C$_{2h}$ or, more realistically, to C$_4$ allows a final classification of the single ion levels. The crystal magnetic excitons in D$_{2h}$ or C$_4$, are then easily found using the Loudon method [20]. In the planar fourfold symmetry approximation each single ion transition induces a doubly degenerate exciton or a Davydov split pair, depending upon the representation of the...
transition. The Raman polarization selection rules are ZX, XZ, ZY, YZ for the degenerate excitons while such polarizations are forbidden for Davydov pairs. The use of the exact in place of the approximate space group removes the equivalence of X and Y: the above mentioned degenerate excitons are split and the ZX and ZY polarizations are no longer equivalent. In the following we shall neglect this last effect which is probably very weak. On the other hand, we shall refer to the representations of D_{4h} rather than C_{4v}, in conformity with the effectively used Hamiltonian (see below): the forms of the Raman tensors are not modified but, in the labelling of the exciton representations, the even parity is explicitly indicated.

2.2. Excitons. — The model uses a simple anisotropic Heisenberg Hamiltonian in order to describe the exchange interaction within 4F^+.

2.2.1 The mean field model. — The cubic crystal field splits the 4F free Co^{2+} ion ground state into three components 4T_{1g}, 4T_{2g} and 4T_{1u}, the lowest being 4T_{1g}. The mean field substates are obtained by diagonalizing the following Hamiltonian within 4T_{1g}:

$$\mathcal{H}_0 = \mathcal{H}_1 + \mathcal{H}_{\text{MF}}$$

(1)

where \(\mathcal{H}_1\) is a single ion Hamiltonian describing the spin orbit coupling and the tetragonal crystal field:

$$\mathcal{H}_1 = A \mathbf{L} \cdot \mathbf{S} + A \left( L_z^2 - \frac{1}{3} L(L+1) \right) - \frac{15}{4} \frac{z^2}{U} \left[ 2 (L_x^2 S_x^2 + L_y^2 S_y^2 + L_z^2 S_z^2) - (\mathbf{L} \cdot \mathbf{S})^2 \right] .$$

(2)

Within 4T_{1g} the standard basis is defined by a set of twelve \(|L, S, Lz, Ms\rangle\) functions where \(S = 3/2\) and where the orbital part is generated by a pseudo \(L=1\) angular momentum [21]. The first term in \(\mathcal{H}_1\) is related to the first order spin-orbit coupling; \(A\) can be written as [21]:

$$A = \alpha \lambda'$$

(3)

where the \(\alpha'\) coefficient takes into account the mixing of the 4F and 4P free ion states; in the weak field limit (no mixing) it would be equal to \(-1.5\), while it is evaluated to be \(-1.42\) from spectroscopic data in Co^{2+} complexes [21]. \(\lambda'\) is an effective spin-orbit coupling constant which is slightly smaller than the free ion value. 4T_{1g} interacts with 4T_{1u} through the spin-orbit coupling and the tetragonal crystal field. The resulting second order spin-orbit coupling has been shown by Kanamori [22] to take the effective form defined by the third term in \(\mathcal{H}_1\), where \(U\) is the energy of 4T_{1u} above 4T_{1g} (\(U = 8 D_0\) in the weak field limit). Finally the second term in \(\mathcal{H}_1\) simultaneously takes into account the first order tetragonal crystal field effect and the remaining second order terms [23] (which act by renormalizing \(A\)).

\(\mathcal{J}_{\text{MF}}\) is written as:

$$\mathcal{J}_{\text{MF}} = 2 z J_{\parallel} \langle S_Z \rangle_0 S_Z$$

(4)

where \(2 J_{\parallel}\) is the (negative) Ising part of the exchange between two neighbouring antiferromagnetic ions in a plane layer and where \(\langle S_Z \rangle_0\) stands for the mean value of \(S_Z\) in the ground state; \(z = 4\) is the number of antiferromagnetic neighbours.

For a given set of the constants involved in the expression of \(\mathcal{J}_{\text{MF}}\), a self-consistent diagonalization has to be performed: it gives rise to twelve sublevels (labelled \(p\), from 0 to 11), whose representations in \(C_{4h}\) are easily found.

A magnetic field along Z induces a perturbation which is approximately written as:

$$\mathcal{J}_{\text{H}} = (1/2) \langle p \mid L_z + 2 S_Z \rangle \mu_B H$$

(5)

where \(r\) is an orbital reduction factor taking into account the admixture of the 2p and 2s fluorine ligand orbitals in the Co^{2+} d orbitals [21]: its measured value, from spin resonance experiments in KMgF_3: Co is 0.93. Under \(\mathcal{J}_{\text{H}}\) the two sublattices no longer remain equivalent and, as a result, a splitting \(\delta W_p\) of each sublevel, is easily calculated to first order in \(H\):

$$\delta W_p = \left( 2 - 8 \frac{z}{2} - 8 \right) \frac{\langle p \mid L_z + 2 S_Z \rangle \mu_B H}{\langle p \mid L_z + 2 S_Z \rangle}$$

(6)

with:

$$\mathcal{K}_{\text{sp}} = \langle p \mid S_Z \mid p \rangle , \quad \mathcal{K}_{\text{lp}} = \langle p \mid L_z \mid p \rangle ,$$

$$\mathcal{K}_S = -2 \sum_{1 \leq q \leq 11} \frac{\langle q \mid S_Z \rangle \langle 0 \rangle}{E_q} = \langle q \mid S_Z \rangle \langle 0 \rangle,$$

$$\mathcal{K}_{LS} = -2 \sum_{1 \leq q \leq 11} \frac{\langle 0 \mid S_Z \rangle \langle q \mid L_z \rangle \langle 0 \rangle}{E_q} = \langle q \mid L_z \rangle \langle 0 \rangle.$$

Expression (6) allows an effective \(g_{||p}\) factor to be defined for each \(|p \rangle \rightarrow |q \rangle \) transition:

$$g_{||p} = \frac{(\delta W_p - \delta W_0)/2 \mu_B H}{(\delta W_p - \delta W_0)/2 \mu_B H}.$$
and
\[ \mathcal{H}_\text{opt} = \mathcal{H}_1 + J_\parallel [2(z - 1) \langle S_z \rangle + \langle S_z \rangle] S_z. \]  
(10)

In (10), \( \langle S_z \rangle \) is the mean value of \( S_z \) in the first excited state.

2.2.2 Propagating excitons. — Starting from the single ion mean field levels, the propagating excitons are found using the following Hamiltonian:

\[ \mathcal{H}_\text{T} = \sum_{i} \mathcal{H}_{1i} + \sum_{i} \mathcal{H}_{1i} - \sum_{i, \bar{i}} 2(J_\parallel S_{xi} S_{\bar{x}\bar{i}} + J_\perp (S_{xi} S_{\bar{x}\bar{i}} + S_{yi} S_{\bar{y}\bar{i}})) \]  
(11)

(next neighbours)

where \( i \) and \( \bar{i} \) respectively run over up and down sublattices and where the last summation is restricted to antiferromagnetic nearest neighbours. Each mean field excited state \( |p\rangle \) gives rise to a degenerate exciton if \( |p\rangle \) belongs to \( C_4h \) or \( C_4v \) and gives rise to a Davydov (eventually split) pair if \( |p\rangle \) belongs to \( C_4d(C_{4v}) \) or \( C_4d(C_{4v}) \) : the perturbation to the mean field model is only due to \( J_\perp \) for the degenerate excitons and to \( J_\parallel \) for the Davydov pairs. As previously noticed, some physical properties have been successfully interpreted using a Heisenberg exchange interaction (which is described by eq. (11)) while a pure Ising exchange within \( 4\Gamma^1 \) implies \( J_\parallel = 0 \) : the Ising-like properties are explained by the large value of \( A \) which leads to a large gap and a rather small dispersion of the lowest excitonic branch, even supposing an isotropic exchange. The question then arises whether a Heisenberg exchange still can account for the new experimental data in \( \text{K}_3\text{CoF}_4 \) and \( \text{Rb}_2\text{CoF}_4 \).

The excitonic states are obtained using the pseudo-boson operator method described by Walker [24] and previously used by Buyers et al. for the study of \( \text{KCoF}_3 \) [1] and \( \text{CoF}_2 \) [2]. Introducing the operators:

\[ a^+_p = |p\rangle \langle o| \quad \text{and} \quad a_{pl} = |o\rangle \langle p| \quad 1 \leq p \leq 11 \]  
(12)

which act on the single ion mean field states and satisfy \( [a_{pl}, a^+_{pl}] = [a_{pl}, a^+_{pl}] = 1 \) (the other commutators being 0), the Hamiltonian can be written, neglecting the 3rd and 4th order terms, as:

\[ \mathcal{H}_\text{T} \simeq \sum_{p \in \Gamma^1} E_p(a^+_p a_{pl} + a^+_p a_{pl}) \]
\[ + \sum_{p \in \Gamma^1} E_p(a^+_p a_{pl} + a^+_p a_{pl}) + 2J_{\parallel} \sum_{p \in \Gamma^1} S^{\text{op}}(a^+_p a^+_p a_{pl} + a_{pl} a_{pl} + a^+_p a_{pl} + a_{pl} a_{pl}) \]
\[ + \sum_{p \in \Gamma^1} E_p(a^+_p a_{pl} + a^+_p a_{pl}) - 2J_{\perp} \sum_{p \in \Gamma^1} S^{\text{op}}[(1 + \epsilon_p \epsilon_q)(a^+_p a_{pl} + a_{pl} a^+_p + a^+_p a_{pl} + a_{pl} a^+_p)] \]
\[ + \sum_{p \in \Gamma^1} E_p(a^+_p a_{pl} + a^+_p a_{pl}) \]
(13)

where:

\[ S^{\text{op}} = \langle o| S_z | pi \rangle, \]
\[ S^{\text{op}} = \langle o| S_x | pi \rangle \]

and \( \epsilon_p = \pm 1 \) (\( S_{\text{op}} = \langle o| S_x | pi \rangle \) and \( S_{\text{op}} \) have the same modulus but, depending upon the state \( |p\rangle \), their arguments differ by \( \pm \pi/2 \) or \( -\pi/2 \), which leads to \( \epsilon_p = \pm 1 \).

The Hamiltonian (13) is diagonalized with the help of a Fourier development followed by a linear transformation, as described in reference [1]. It is easily seen from eq. (13) that the diagonalization can be performed separately for the degenerate excitons and for the Davydov pairs. Moreover the specific form of \( \mathcal{H}_\text{T} \) does not induce any shift or splitting for the excitons related to the \( p \in \Gamma^1 \) mean field states : actually they have not to be regarded as forming a Davydov pair and they behave in the same way as mean field states.

For a further discussion it is of interest to compare the above model with the simplified one [5, 7] where the interaction between excited levels is neglected, i.e.:
The dispersion formula resulting from eq. (14) are given by:

\[ E_p(k) = E_p \]

\[ E_p(k) = E_p \left( 1 \pm \left( 4 \frac{S_2^{\pm} r_2}{E_p} \right)^{1/2} \right) \text{ if } |p| \in \Gamma_5^+ \] (15b)

\[ E_p(k) = E_p \left( 1 - \left( 4 \frac{S_2^{\pm} r_2}{E_p} \right)^{1/2} \right) \text{ if } |p| \in \Gamma_5^+ \text{ or } \Gamma_8^+ \] (15c)

with \( r_2 = \cos (ak_x) \cos (ak_y) \) (\( a \) : distance between 2 antiferromagnetic neighbours).

As shown in section 4, for the used values of the constants, expressions (15b) and (15c) do not much differ from the results that we obtained from an exact diagonalization of eq. (13). On the other hand, for these values, the lowest exciton is degenerate and an immediate comparison can be made between its dispersion given by eq. (15c) (for \( p = 1 \) in this case) and the corresponding expression using a spin 1/2 model: Breed et al. [15] interpret the magnetic properties of \( \text{K}_2\text{CoF}_4 \) and \( \text{Rb}_2\text{CoF}_4 \) using the following effective spin 1/2 Hamiltonian:

\[ h = - \sum_{n} \left( 2j_{n} s_{x n} s_{y n} + 2j_{n} (s_{x 1} s_{y 1} + s_{y 2} s_{y 2}) \right). \] (16)

The spin-wave spectrum derived from (16) is given by:

\[ e(k) = - \tilde{z} \gamma_2 \left( 1 - \left( \frac{j_{\parallel}}{j_{\perp}} \gamma_2 \right) \right)^{1/2} \] (17)

\( e(k) \) has the same form as \( E_q(k) \) if one chooses:

\[ f_{\parallel} = - E_q/z \] (18a)

\[ j_{\perp} = 4 J_{\perp} S_2^{012}. \] (18b)

Moreover, we found that the exact dispersion law for the lowest degenerate exciton obtained by diagonalizing eq. (13) can be written in the form of eq. (17) but the derivation of \( j_{\parallel} \) is no longer given by eq. (18b). Finally the effect of a magnetic field applied along the \( Z \) axis can be calculated using the self-consistently modified mean field states which form a new basis where the two sublattices are no longer equivalent. The diagonalization is substantially simplified if one neglects the interaction between excited mean field levels of different \( p \) : to first order in \( H \), it is easily shown that the degenerate excitons are split exactly in the way predicted by the mean field model (expressions (6) and (7)) ; indeed, this result is no longer true for the Davydrop pairs which show a shift quadratic in \( H \), at least when \( \mu_B H \) is small compared to the \( H = 0 \) Davydrop splitting. As discussed in the following sections, our experimental data in a magnetic field only concern the Davydrop excitons; as long as, for \( H = 0 \), the simplified excitonic model of eq. (14) and the more exact one of eq. (13) do not too markedly differ, we believe that the mean field \( g \) values provide a sufficiently good approximation, and we used these values to fit our results. Notice that a more precise evaluation involving a complete diagonalization in presence of \( H \) is slightly tedious but not difficult, but due to the limited experimental precision we assumed, at least as a first step, that it was not necessary.

### Table I.

<p>| Observed polarization and symmetry representations in ( \text{D}_{4h} ) (( Z//C_4 ); ( X, Y//C_2 )). |
|---|---|---|---|---|</p>
<table>
<thead>
<tr>
<th><strong>Spectrum</strong></th>
<th><strong>1st setting</strong></th>
<th><strong>Symmetry</strong></th>
<th><strong>2nd setting</strong></th>
<th><strong>Symmetry</strong></th>
</tr>
</thead>
<tbody>
<tr>
<td>( X(ZX) )</td>
<td>---</td>
<td>( \Gamma_5^+ )</td>
<td>( X(ZY) )</td>
<td>( Z )</td>
</tr>
<tr>
<td>( X(YX) )</td>
<td>( \Gamma_5^+ ), ( \Gamma_3^+ )</td>
<td>( X(XY) )</td>
<td>( \Gamma_5^+ ), ( \Gamma_3^+ )</td>
<td></td>
</tr>
<tr>
<td>( X(YZ) )</td>
<td>( \Gamma_5^+ )</td>
<td>( X(XZ) )</td>
<td>( \Gamma_5^+ )</td>
<td></td>
</tr>
<tr>
<td>( X(ZZ) )</td>
<td>( \Gamma_5^+ )</td>
<td>( X(ZX) )</td>
<td>( \Gamma_5^+ )</td>
<td></td>
</tr>
</tbody>
</table>

Raman tensors in \( \text{D}_{4h} \), with the same reference axes as above:

\[
\begin{bmatrix}
0 & c & 0 \\
0 & a & 0 \\
0 & b & 0
\end{bmatrix}
\begin{bmatrix}
d & 0 & 0 \\
0 & 0 & 0 \\
0 & 0 & 0
\end{bmatrix}
\]

\[
\begin{bmatrix}
0 & e & 0 \\
e & 0 & 0 \\
0 & 0 & 0
\end{bmatrix}
\begin{bmatrix}
0 & 0 & f \\
0 & 0 & 0 \\
0 & 0 & 0
\end{bmatrix}
\]

\[
\begin{bmatrix}
0 & 0 & 0 \\
0 & g & 0 \\
0 & 0 & 0
\end{bmatrix}
\begin{bmatrix}
0 & 0 & 0 \\
0 & 0 & 0 \\
0 & 0 & 0
\end{bmatrix}
\]

\( \Gamma_1^+ \) (antisymmetric) \( \Gamma_3^+ \)

\( \Gamma_5^+ \) (symmetric)
Table II. — Raman lines in K₂CoF₄ near liquid Helium temperature. H is the applied magnetic field, γ is the width at half-maximum (vw : very weak ; w : weak ; ms : medium strong ; s : strong ; vs : very strong. There are roughly three orders of magnitude between vw and vs).

<table>
<thead>
<tr>
<th>Position (cm⁻¹)</th>
<th>Intensity, width : γ (cm⁻¹)</th>
<th>Polarization</th>
<th>Observations</th>
<th>Symmetry</th>
</tr>
</thead>
<tbody>
<tr>
<td>108</td>
<td>s, γ = 3</td>
<td>ZZ, ZY</td>
<td>phonon</td>
<td>Γ₅⁺</td>
</tr>
<tr>
<td>154</td>
<td>s, γ = 3</td>
<td>ZZ, ZY</td>
<td>phonon</td>
<td>Γ₅⁺</td>
</tr>
<tr>
<td>163.5</td>
<td>vw, γ = 2</td>
<td>ZZ, ZY</td>
<td>split by H</td>
<td>Γ₅⁺</td>
</tr>
<tr>
<td>184</td>
<td>s, γ = 3</td>
<td>ZZ</td>
<td>phonon</td>
<td>Γ₁⁺</td>
</tr>
<tr>
<td>207</td>
<td>ms, γ = 3</td>
<td>ZZ, ZY</td>
<td>split by H</td>
<td>Γ₅⁺</td>
</tr>
<tr>
<td>314</td>
<td>vw</td>
<td>ZZ, ZY</td>
<td>phonon</td>
<td>Γ₅⁺</td>
</tr>
<tr>
<td>355</td>
<td>s, γ = 3</td>
<td>ZZ</td>
<td></td>
<td>Γ₁⁺</td>
</tr>
<tr>
<td>374</td>
<td>vs</td>
<td>ZZ</td>
<td>phonon</td>
<td>Γ₁⁺</td>
</tr>
<tr>
<td>407</td>
<td>broad band</td>
<td>major component ZZ</td>
<td></td>
<td>Γ₁⁺</td>
</tr>
<tr>
<td>426</td>
<td>vs</td>
<td>ZZ</td>
<td></td>
<td>Γ₁⁺</td>
</tr>
<tr>
<td>465</td>
<td>vw</td>
<td>ZZ</td>
<td></td>
<td>Γ₁⁺</td>
</tr>
<tr>
<td>555</td>
<td>vw</td>
<td>ZZ</td>
<td></td>
<td>Γ₁⁺</td>
</tr>
<tr>
<td>565</td>
<td>w</td>
<td>ZZ</td>
<td></td>
<td>Γ₁⁺</td>
</tr>
<tr>
<td>716</td>
<td>ms</td>
<td>ZZ</td>
<td></td>
<td>Γ₁⁺</td>
</tr>
<tr>
<td>729</td>
<td>w</td>
<td>XY</td>
<td></td>
<td>Γ₁⁺</td>
</tr>
<tr>
<td>1180</td>
<td>vs, γ = 40</td>
<td>ZZ</td>
<td></td>
<td>Γ₁⁺</td>
</tr>
<tr>
<td>1417</td>
<td>w</td>
<td>ZZ</td>
<td></td>
<td>Γ₁⁺</td>
</tr>
</tbody>
</table>

Temperature of the scattering volume was estimated to lie around a few K. Some studies of their temperature dependence were performed using an exchange gas cryostat in the 5-300 K temperature range. A magnetic field up to 4.8 teslas, parallel to the studied scattering direction, was available for the low temperature measurements.

Absorption is severe in the visible and useful spectra could be obtained principally with the 4 579 Å and 4 765 Å Argon laser lines (~ 500 mW) and the 6 471 Å Krypton laser line (~ 300 mW). For all the lines used the absorption is stronger when the polarization is perpendicular to Z : the experimental intensity is then significantly lower when the incident light is polarized along X and Y. A Coderg T800 monochromator with a photon counting detection was used to analyse the scattered light. For some lines the weak intensities prevented polarization selection rules to be studied.

3.2 RESULTS. — Besides the 4 known phonon lines [25], which are slightly shifted and sharpened, we observed new features at low temperatures. The complete results are listed in table II and table III, respectively for K₂CoF₄ and Rb₂CoF₄.

Figure 2 displays a representative spectrum, for a wide wave-number range scan in K₂CoF₄. The major new features are three sharp and rather intense lines at 163.5 cm⁻¹ (Γ₅⁺ i.e. : ZZ, ZY, ZX, XY) and 355 cm⁻¹ (Γ₁⁺) and two broad bands around 460 and 1 180 cm⁻¹ ; in addition we observed two sharp Γ₁⁺ lines in the high frequency range (716 cm⁻¹, 1 417 cm⁻¹) and a weak XY line at 729 cm⁻¹. A magnetic field H parallel to Z linear-
Fig. 2. — A typical scan of the Raman spectrum of K$_2$CoF$_4$ on a wide wave-number range ($T \sim 2$ K).

Ly splits the 163.5 and 207 cm$^{-1}$ peaks, allowing an effective $g_\parallel$ factor to be defined (table IV), while it does not affect the remaining spectrum in a measurable way. Figure 3 shows the splitting of the 207 cm$^{-1}$ line by $H$. As expected from the high anisotropy along $Z$, which predicts very small quadratic shifts when $H$ is perpendicular to $Z$, we observed that $H$ has no effect in this case.

Table IV. — Experimental $g_\parallel$ values for the lines split by a magnetic field.

<table>
<thead>
<tr>
<th>Crystal</th>
<th>$\omega$ (cm$^{-1}$)</th>
<th>$g_\parallel$</th>
</tr>
</thead>
<tbody>
<tr>
<td>K$_2$CoF$_4$</td>
<td>163.5</td>
<td>3.9</td>
</tr>
<tr>
<td></td>
<td>207</td>
<td>3.7</td>
</tr>
<tr>
<td>Rb$_2$CoF$_4$</td>
<td>155</td>
<td>5</td>
</tr>
<tr>
<td></td>
<td>179</td>
<td>1.4</td>
</tr>
<tr>
<td></td>
<td>209</td>
<td></td>
</tr>
<tr>
<td></td>
<td>213</td>
<td>4.45</td>
</tr>
</tbody>
</table>

Fig. 3. — Splitting of the 207 cm$^{-1} \Gamma'_z$ line in K$_2$CoF$_4$ in a magnetic field parallel to Z.

The 213 cm$^{-1}$ line presents a complex shape, with a shoulder at 209 cm$^{-1}$, which may indicate the existence of two neighbouring peaks; however this shape is apparently unmodified under the magnetic field splitting (fig. 4) and, due to this apparent rigidity, we analyzed it as consisting of a single line. In contrast with the case of K$_2$CoF$_4$, the precision for the $g_\parallel$ factors is poor in Rb$_2$CoF$_4$ (table IV): the 155 and 178 cm$^{-1}$ peaks lie very near to an intense phonon line, and precise position measurements cannot be made; the complex shape of the 213 cm$^{-1}$ also prevents precise evaluations being made.

In Rb$_2$CoF$_4$ the low frequency spectrum shows four new lines, three of them (155, 178, 213 cm$^{-1}$) being linearly split by a magnetic field parallel to $Z$ and showing $\Gamma'_z$ symmetry, while the fourth (322 cm$^{-1}$) is unsplit and has a $\Gamma'_z$ behaviour. Here again a high frequency broad spectrum also exists.

When increasing the temperature up to around $T_N$, the low frequency spectrum (below about 350 cm$^{-1}$) broadens and disappears in both compounds, which proves that it is related to the magnetic ordering. In the high frequency range some broad band scattering persists, but it is difficult to analyse quantitatively.

4. Interpretation. — Due to their general properties it is clear that most of the new reported features are related to electronic Raman scattering by magnetic excitons. An attempt to give a very precise quantitative account of the high frequency spectrum (500 cm$^{-1}$ and above) is known to be unrealistic in Co fluorides [7, 9, 10] partly because most of the lines are rather broad, which prevents precise measurements, especially when neighbouring features overlap, partly because it is difficult to separate the pair-excitation contributions (two-phonon and two-magnon spectra) from the single excitation ones. On the other hand, in the upper levels the frequencies are mainly monitored
by the spin-orbit and the low symmetry crystal field constants and an attempt to derive a reasonably precise determination of the exchange from their study would then be hopeless. We found it more efficient to focus on the low frequency spectrum, which is generally better understood, and to derive adequate constants in order to incorporate them in the Hamiltonian of the system; in a further step it would be of interest to compare our high frequency spectrum with the predictions resulting from the diagonalization of this Hamiltonian.

At first glance, the 163.5 and 207 cm\(^{-1}\) lines in K\(_2\)CoF\(_4\) would have to be identified with the two lowest \(I_s^+\) degenerate excitons: however it is not possible to find any reasonable set of values for \(A, \lambda, \lambda, \lambda, \ldots\) which provides good or even approximate fits to our experimental frequency and \(g_{||}\) values. The situation is worse in Rb\(_2\)CoF\(_4\) where, in the same spirit, one could be tempted to attribute the three lowest lines (155, 178 and 213 cm\(^{-1}\)) to the three lowest \(I_s^+\) degenerate excitons. From previously published data concerning magnetic properties [11, 12], we expected that, in both crystals, the lowest \(I_s^+\) branch lies around 200 cm\(^{-1}\), rather well below the other excitons. On the other hand, in K\(_2\)CoF\(_4\), the 207 cm\(^{-1}\) line shows an intensity stronger than that at 163.5 cm\(^{-1}\), and in Rb\(_2\)CoF\(_4\) the 213 cm\(^{-1}\) line is the most intense among the three \(I_s^+\) peaks: we then assumed that the 207 cm\(^{-1}\) and 213 cm\(^{-1}\) lines correspond to the lowest \(k = 0 I_s^+\) exciton respectively in K\(_2\)CoF\(_4\) and Rb\(_2\)CoF\(_4\). We were definitely convinced of the adequacy of this choice by the recently published inelastic neutron scattering measurements in Rb\(_2\)CoF\(_4\) [4] which provide for the \(k = 0\) lowest exciton frequency a 206 ± 6 cm\(^{-1}\) value, which is near our 213 cm\(^{-1}\) measured value. Notice that, due to the better precision of optical spectroscopy measurements (± 0.3 cm\(^{-1}\)), the 213 cm\(^{-1}\) determination appears to be the best one; however, as mentioned above, the complex shape of the line, does not enable the exact value of the frequency to be established: in any case it is not lower than 209 cm\(^{-1}\); such a 4 cm\(^{-1}\) shift would not significantly change the following conclusions.

The game to play is now to find the parameters which provide a 207 cm\(^{-1}\) frequency with a \(g_{||} = 3.7\) value in K\(_2\)CoF\(_4\) and a 213 cm\(^{-1}\) frequency with a \(g_{||} = 4.45\) value in Rb\(_2\)CoF\(_4\) for the lowest \(k = 0 I_s^+\) exciton. The \(A, \alpha, \lambda\) and \(\lambda\) constants appearing in eqs. (2), (3) and (5) show nearly identical values in different cobaltous compounds, provided that the Co\(^{2+}\) environment is a pseudo-cubic fluorine octahedron : following Thornley [21], Buyers [11] and others, we then took: \(A = 228\) cm\(^{-1}\), \(\alpha = -1.42\) (and \(\lambda = -160\) cm\(^{-1}\)), \(r = 0.93, \lambda = 7500\) cm\(^{-1}\). Since an Heisenberg interaction is often expected to provide a good approximate description of the exchange, we chose a fit with \(J_\perp = J_\parallel = J\).

For each compound we are then left with two parameters \(A\) (or, most conveniently, for further comparisons, \(A/A\)) and \(J\) in order to fit two experimental results: the frequency and the \(g_{||}\) value of the \(k = 0\) lowest exciton. This was done using the model developed in section 2.

The method is shown on figure 5 in the case of K\(_2\)CoF\(_4\). Fixing the lowest exciton frequency at 207 cm\(^{-1}\) induces a relation between \(J\) and \(A\) : we have then reported graphically \(J\) as a function of \(A/A\). Due to this relation between \(J\) and \(A/A\), \(g_{||}\) can now be evaluated as a function of \(A/A\) only: the value of \(A/A\) corresponds to \(g_{||} = 3.7\). From figure 5, one immediately derives \(J\) and \(A/A\). We have also reported on figure 5 the dispersion \(\delta\) of the exciton, defined as the ratio of the difference between the Brillouin zone boundary and the \(k = 0\) frequencies to the Brillouin zone boundary frequency (which, in our model, is identified with the mean field value).

The results are reported in table V and are compared with other determinations. For K\(_2\)CoF\(_4\) we find

\[
A/A = -1.86 \quad \text{and} \quad J = -11.55 \text{ cm}^{-1},
\]

which lead to \(\delta = 0.175\). In Rb\(_2\)CoF\(_4\) we obtain:

\[
A/A = -2.78 \quad \text{and} \quad J = -9.65 \text{ cm}^{-1},
\]

which give \(\delta = 0.105\). In the case of Rb\(_2\)CoF\(_4\) our calculated \(\delta\) value allows a direct comparison to be made with the inelastic neutron scattering measurements of Ikeda and Hutchings [4]: we predict a 239 cm\(^{-1}\) zone boundary frequency in good agreement with their 250 ± 8 cm\(^{-1}\) direct experimental determination. Using eq. (17) we find that the \(J_\parallel\) and \(j_\perp\) parameters adequate for a spin 1/2 description are:

\[
j_\parallel = -E_1(0)/4(1 - \delta) \quad (19a)
\]

\[
j_\perp = -\sqrt{\frac{\delta(2 - \delta)}{4(1 - \delta)}} E_1(0) \quad (19b)
\]

Fig. 5. — Variations of \(J, j_\parallel, j_\perp, \lambda\) and \(\delta\), as functions of \(A/A\) when assuming a 207 cm\(^{-1}\) value for the lowest \(k = 0\) exciton frequency. From \(g_{||} = 3.7\) one immediately graphically derives \(A/A, J\) and \(\delta\).
where $E_1(0)$ is the exciton frequency at $k = 0$ (respectively 207 and 213 cm$^{-1}$ for K$_2$CoF$_4$ and Rb$_2$CoF$_4$). Our $J$ determination agrees rather well with the Breed et al. value [15], based on susceptibility measurements but our $J_\parallel$ determination is significantly higher, which means that the Ising behaviour is less marked for the excitation spectrum than for the susceptibility measurements or for the magnetic transition temperature. It is to be noticed that eq. (13) does not take into account the Oguchi's quantum corrections [26], which appear in the normal ordering of the neglected four operator interactions: generally their main effect is to renormalize the exchange interaction by about 5%. So, we do not think that neglecting these corrections could have major effects upon the Ising character. Anyway our predictions concerning the dispersion agree well with the inelastic scattering measurements in the case of Rb$_2$CoF$_4$, which is a sufficient experimental justification.

Our calculations show that in order to relate $J_\parallel$ and $J_\perp$ to $J$ and $A/A$ it is important to exactly diagonalize eq. (13): the use of the simplified model of eq. (14) would lead to significantly different determinations for $J_\parallel$ and $J_\perp$. In the $J$ and $A/A$ studied range the excited levels interaction, included in eq. (13), reduces the effective Ising character (as measured by $J_\parallel/J_\perp$).

Coming back to our determination of $J$ we find values which do not much differ from the observed ones in KCoF$_3$ [1, 9] and RbCoF$_3$ [8] (respectively $-10.6$ cm$^{-1}$ and $-8.9$ cm$^{-1}$): since the superexchange is mediated in the same way through a $180^\circ$ fluorine bridge, the $J$ values are expected to be about the same.

In table VI we have listed all the excitonic levels, as calculated from the model using the above discussed parameters; the useful representations as well as the polarization selection rules are indicated. The dispersion of the excitons against $\gamma(k)$ is indicated: indeed, the frequencies for $\gamma(k) = 1$ are those of the mean field model. We have also reported the single ion energy levels resulting from $\mathfrak{H}_I(D_{4h})$. As usual [1, 2], only the low frequency excitons show a non negligible dispersion and the mean field model can be used without loss of precision for the high energy excitons.

An attempt to account for the experimental high frequency spectrum using our model is only moderately satisfactory: in K$_2$CoF$_4$ the four components broad band around 460 cm$^{-1}$ (table II) is related to the excitons arising from the $p = 2$, 3 and 4 mean field levels (table VI), while the 1180 cm$^{-1}$ wide band is probably connected with the $p = 8$ and 9 mean field levels. In addition, the two-magnon peak maximum is predicted from eq. (8) to lie at 432 cm$^{-1}$, while its cut-off frequency is 500 cm$^{-1}$: the 460 cm$^{-1}$ band, which, in fact, extends from 405 to 500 cm$^{-1}$, then certainly contains a two-magnon contribution, which is always present in such antiferromagnets [7, 9, 11]. While on the basis of frequency considerations the 716 and 1417 cm$^{-1}$ lines could be attributed respectively to $p = 6$ and $p = 11$, we do not retain such an assignment since it would disagree with the observed polarization selection rules. In Rb$_2$CoF$_4$, due to the poor quality of the high frequency observed spectra, we prefer to delay any attribution.

The above general analysis would finally be considered as successful if it completely explains the low frequency spectrum. Unfortunately in K$_2$CoF$_4$ our model is unable to give account of the $I_5^-$ 163.5 line and of the $I_5^+$ 335 cm$^{-1}$ one; in Rb$_2$CoF$_4$, the 155 cm$^{-1}$ and 178 cm$^{-1}$ $I_5^+$ lines remain to be interpreted as well as the rather intense 332 cm$^{-1}$ $I_5^+$ peak. The attribution of the $I_5^+$ lines to two-
magnon scattering has probably to be rejected, due
to the sharpness of the corresponding peaks which
would not occur in a two-magnon process; on the
other hand, an approximate evaluation of the
two-magnon peak maximum position provides a
value of 432 cm\(^{-1}\) and of 403 cm\(^{-1}\), respectively
for K\(_2\)CoF\(_4\) and Rb\(_2\)CoF\(_4\), which is too high for
such an attribution. The unexplained low frequency
lines, could also derive from the doubling of the
nuclear cell below the magnetic ordering temperature:
strictly speaking, the nuclear Brillouin zone boundary
phonons along \(k = (1, 1, 0)\) are \(k = 0\) phonons
when using the magnetic unit cell, and they could
eventually give rise to Raman features. Such pro-
cesses have been observed especially in Co com-
ounds [1, 27]. However the intensity is non negli-
gible only in the case of a strong exciton-phonon
coupling, which supposes that the exciton and the
phonon frequencies lie near to each other. Even if
we do not disregard such an hypothesis, we would
be surprised if it occurred for two or even three dif-
ferent phonon branches. Especially, the spectacular
electronic character of the \(T_5\) lines, attested by
their large \(g_{ll}\) values, implies a coupling to excitons
much more efficient than the previously observed
ones, in CoF\(_2\) [7, 28] for instance.

The most probable hypothesis is that the unex-
plained lines are related to impurity modes: it is
clear that such a vague attribution needs further
analysis but we are not surprised of such occurrences
and we have some experiment of impurity modes
showing a unusually (and unexplained) high inten-
sity in RbMnF\(_3\) and KMnF\(_3\).

In view of the remaining difficulties a pessimistic
attitude would lead to reject completely the model,
taking argument of the occurrence of complex
exchange, which is expected in the difficult case of
an unquenched orbital momentum [12, 13] in the
\(4T_2^+\) ground state of Co\(^{2+}\). We disagree with such
an extreme position: our determinations on the
basis of one exchange constant within \(4T_2^+\) are well
consistent with all the other data available for the
observed crystals. It has been shown for many other
cobaltous compounds that, at least if one does not
try to give account of the high energy excitons very
precisely, such a model provides a good description
of the magnetic properties. Finally our \(J\) determi-
nations are consistent with the values reported for
crystals where the exchange is expected to show
approximately the same value.

NOTE ADDED IN PROOF. — In a recent letter [29],
F. Macco et al. reported new Raman data in K\(_2\)CoF\(_4\).
Except for a few minor features, we agree with their results and with their assignment of the 206.5 cm\(^{-1}\) line as the one-magnon peak. We think that our experimental study, which deals also with Rb\(_2\)CoF\(_4\) and which includes the dependence on magnetic field, as well as our analysis, confirms and supplements these preliminary measurements.

Acknowledgments. — We thank E. Guiot and N. Lenain for polishing the crystals.

References