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Second-Order Modulation of a Transversal Light Beam by Magnetic Resonance in Na Oriented Vapour.

G. ALZETTA and E. ARIMONDO

Istituto di Fisica dell'Università - Pisa
Gruppo Nazionale di Struttura della Materia - Pisa

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Summary. — Second-order effects due to magnetic resonance of sodium atoms oriented by optical pumping have been investigated by the transversal-beam monitoring technique. An extension of Cohen-Tannoudji and Laloë theory has been developed to include these effects and its implications compared with the experimentals results.

Introduction.

Several methods have been developed for the study of ground-state magnetic resonance of atoms oriented by optical pumping. In particular, experiments with circularly polarized transversal light beams, monitoring the magnetization in a direction orthogonal to the static magnetic field, have been reported (1,2). In resonance conditions, light absorption at the Larmor frequency is observed with an amplitude dependent on the intensity of the transversal magnetization induced by the radiofrequency field. In the experiments on Hg (3,4) the Faraday effect at the Larmor frequency has been examined on the transversal beam.

A general method has been devised by COHEN-TANNOUDJI and LALOË (5) that thoroughly accounts for these effects. However in evaluating the intensity

(1) H. G. DEHMELT: *Phys. Rev.*, **105**, 1924 (1957).
(2) W. E. BELL and A. L. BLOOM: *Phys. Rev.*, **107**, 1559 (1957).
(3) J. MANUEL and C. COHEN-TANNOUDJI: *Compt. Rend.*, **257**, 413 (1963).
(4) F. LALOË, M. LEDUC and P. MINGUZZI: *Compt. Rend.*, B **266**, 1517 (1968).
(5) C. COHEN-TANNOUDJI and F. LALOË: *Journ. Phys.*, **28**, 505, 722 (1967).

of the light received by the detector the analysis has been limited to the linear terms of the density matrix: a Larmor frequency modulation is obtained in alkali vapours monitoring with either D_1 or D_2 light beams. Nevertheless, neglecting multiple scattering, it is possible to obtain from the theory that a modulation of the transversal light beam at the second harmonic of the Larmor frequency of the spin system has to occur. This case is of importance when observation of the effects of the resonance on the monitoring beam is performed with crossed linear polarizer and analyser. Such an arrangement has been used by CORNEY, KIBBLE and SERIES ([6]) to study the forward scattering in Hg and by NOVIKOV ([7]) to examine the Cotton-Mouton effect in Cs vapour. Modulation at higher frequencies also results in multiple-scattering phenomena where the optical thickness of the sample has to be taken into account ([8]).

In the present investigation, modulation phenomena due to the quadratic terms of the density matrix are examined. An extension of the Cohen-Tannoudji and Laloë theory to this case is given and compared with the experimental results obtained on Na oriented vapours.

1. – Second-order modulation.

Consider an oriented vapour; let $|\varepsilon_\omega^I\rangle$ represent, in the vector space of the polarization states, the vector which describes the component at frequency ω of the monitoring light incident in the x-direction orthogonal to the magnetic field. The vector of the transmitted wave is

$$|\varepsilon_\omega^T(x_0, t)\rangle = M_\omega\left(t-\frac{x_0}{c}\right)|\varepsilon_\omega^I\rangle ,$$

where M_ω is an operator which depends on the properties of the vapour. The explicit form of M_ω can be simplified by the following assumptions:

i) the density matrix of the vapour in the ground state is independent of the x-position, *i.e.* the vapour is supposed to be uniform;

ii) the time taken by the monitoring light to traverse the vapour cell is short compared to the evolution time of the atomic system.

Then one gets (ref. ([5]), eqs. (6.8) and (6.9))

$$M_\omega\left(t-\frac{x_0}{c}\right)=\exp\left[-C_\omega\left(t-\frac{x_0}{c}\right)\right] ,$$

([6]) A. CORNEY, B. P. KIBBLE and W. SERIES: *Proc. Roy. Soc.*, A **293**, 70 (1966).
([7]) L. N. NOVIKOV: *Žurn. Èksp. Teor. Fiz., Pis'ma v. Redak.*, **6**, 11 (1967).
([8]) G. W. SERIES: *Proc. Phys. Soc.*, **88**, 995 (1966).

where

$$G_\omega(t) = \frac{Nq^2}{2\varepsilon_0\hbar}\frac{\omega l}{c}\,\mathrm{Tr}_f\,\{\sigma^f(t)\,K_\omega(t)\}\,.$$

The refraction indices of the vapour principal polarizations are proportional to the eigenvalues of the operator G_ω and are proportional to the ground-state density matrix. The nondiagonal elements in this matrix oscillate at different frequencies Ω and contribute to an oscillating part in the refraction index of the vapour. The transmitted light beam will contain sidebands at the frequencies $n\Omega$ (n is a positive or negative integer) whose intensity depends on $(\sigma^f)^{|n|}$ (*). The intensity of the transmitted light depends on the interference among the different sidebands so that light beats at all frequencies $n\Omega$ are to be expected.

The « weak scattering » approximation implies that only the first sidebands at the frequencies $\pm\Omega$ are to be retained in the transmitted wave [6]. COHEN-TANNOUDJI and LALOË have made this approximation in their work, but in the calculation of the transmitted intensity they have neglected the interference between the sidebands $+\Omega$ and $-\Omega$. These interference terms depend on the square of the density matrix and are the ones experimentally observable with crossed linear polarizers. This can be shown with a simple model [3].

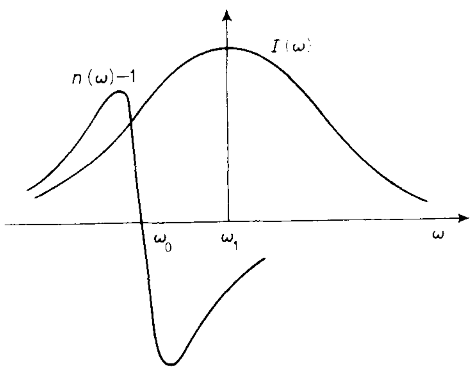

Fig. 1. – Schematic diagram of frequency dependence of the monitoring light and vapour anomalous dispersion.

Let $I(\omega)$ be the spectral distribution of the monitoring light beam: neglecting the diamagnetic effect, consider for the vapour the anomalous dispersion curve $n(\omega)-1$ of width Δ' centered at the frequency ω_0 (Fig. 1). The Faraday rotation of the polarization plane of the wave at a frequency ω will be $M_x[n(\omega)-1]$, where M_x is the vapour magnetization in the direction of the monitoring beam. If the analyser is at an angle θ to the polarizer, the transmitted light is

$$\int I(\omega)\cos^2\{\theta + M_x[n(\omega)-1]\}\,\mathrm{d}\omega$$

(*) The sidebands are the eigenwaves of the Maxwell equation for the light propagation in the medium [6].

and with a series expansion in $M_x[n(\omega)-1]$

$$\cos^2\theta\int I(\omega)\,\mathrm{d}\omega-\sin 2\theta M_x\int I(\omega)[n(\omega)-1]\,\mathrm{d}\omega-\\ -\cos 2\theta M_x^2\int I(\omega)[n(\omega)-1]^2\,\mathrm{d}\omega\,. \tag{1}$$

The first-order term describes the transversal Faraday effect at the Larmor precession frequency. The second-order term can be written $\cos 2\theta M_x^2\Phi$, where

$$\Phi=\int I(\omega)[n(\omega)-1]^2\,\mathrm{d}\omega\propto\int I(\omega)\frac{(\omega-\omega_0)^2}{[(\omega-\omega_0)^2+\Delta'^2/4]^2}\,\mathrm{d}\omega\,, \tag{2}$$

so that, to the second order, symmetric frequencies on both sides of ω_0 equally contribute to the rotation. Owing to M_x^2, this term contains both static and double frequency effects for the Faraday rotation and is the only one different from zero when $\theta=\pi/2$.

As already stated, modulation at harmonics of Ω in the transmitted light is also due to multiple scattering.

Expanding M_ω as a function of G_ω one gets

$$|\varepsilon^T_\omega(x_0,t)\rangle=\left[1-G_\omega\left(t-\frac{x_0}{c}\right)+\frac{1}{2}G^2_\omega\left(t-\frac{x_0}{c}\right)\right]|\varepsilon^I_\omega\rangle\,. \tag{3}$$

Let $\pi_T=I|e_{\lambda_0}\rangle\langle e_{\lambda_0}|$ be the polarization matrix of the incident beam. Then, to the second order, the matrix for the transmitted light is

$$\pi_T(t)=\pi_I-\int I(\omega)\left[G_\omega\left(t-\frac{x_0}{c}\right)\pi_I(\omega)+\pi_I(\omega)G^\dagger_\omega\left(t-\frac{x_0}{c}\right)\right]\mathrm{d}\omega+\\ +\int I(\omega)\left[G_\omega\left(t-\frac{x_0}{c}\right)\pi_I(\omega)G^\dagger_\omega\left(t-\frac{x_0}{c}\right)+\right.\\ \left.+\frac{1}{4}G^2_\omega\left(t-\frac{x_0}{c}\right)\pi_I(\omega)+\frac{1}{4}\pi_I(\omega)\,G^{\dagger 2}_\omega\left(t-\frac{x_0}{c}\right)\right]\mathrm{d}\omega\,. \tag{4}$$

The first-order terms have been considered in detail in [5]. The second-order terms $G^2_\omega\pi_I$ and $\pi_I G^{\dagger 2}_\omega$ are due to multiple scattering. The term $G_\omega\pi_I G^\dagger_\omega$, obtained also in the weak-scattering approximation, gives the Faraday effect with crossed polarizer and analyser. In the following Section all these terms will be considered in detail in the case of alkali metals.

2. – Alkali metals.

Let us consider a monitoring beam containing either D_1 or D_2 light and suppose that the hyperfine structure of the excited level is unresolved so that $\Delta + \Delta' \gg \Delta W$, Δ being the line width of the incident light, Δ' the Doppler width in the vapour, and ΔW the hyperfine splitting in the excited level $^2P_{\frac{1}{2}}$ or $^2P_{\frac{3}{2}}$. No particular hypothesis is made about the relation between $\Delta + \Delta'$ and the hyperfine splitting of the ground state: for Na these quantities are of the same order. Moreover, since only the transitions between Zeeman sublevels have been experimentally examined, the hyperfine coherence will be ignored.

With the previous hypothesis the theory gives to first order the following results (ref. ([5]) eq. (5.14)):

$$\pi_T - \pi_I = -\frac{\alpha}{2}\Gamma_1'\,[B_1(t), \pi_I]_+ - i\alpha\Delta E_1'[B_1(t), \pi_I] - \\ -\frac{\alpha}{2}\Gamma_2'[B_2(t), \pi_I]_+ - i\alpha\Delta E_2'[B_2(t), \pi_I]\,, \tag{5}$$

where the indices 1 and 2 refer to the hyperfine levels of the ground state. Right and left circular polarizations are the principal polarizations of the system and in the representation with this basis the eigenvalues of the matrix $B_i(t)$ turn out to be

$$\begin{cases} \langle e_1|B_i(t)|e_1\rangle = -\dfrac{1}{\sqrt{3}}\langle T_0^{(0)}\rangle_i + \dfrac{1}{\sqrt{2}}\langle T_0^{(1)}\rangle_i\,, \\[2ex] \langle e_2|B_i(t)|e_2\rangle = -\dfrac{1}{\sqrt{3}}\langle T_0^{(0)}\rangle_i - \dfrac{1}{\sqrt{2}}\langle T_0^{(1)}\rangle_i\,, \end{cases} \tag{6}$$

where $\langle T_0^{(q)}\rangle_i$ means $\mathrm{Tr}_f\{P_{F_i}\sigma^f P_{F_i} T_0^{(q)}\}$. Because the hyperfine structure of the excited level is supposed unresolved, the $B_i(t)$ do not depend on $\langle T_0^{(2)}\rangle_i$. Hence by monitoring the vapour with a transversal beam, the oscillating magnetization along the x-direction can be detected observing at the Larmor frequency either the absorption of circular light or the Faraday effect.

If the second-order terms are calculated in the same previous hypothesis, one gets

$$\pi_T - \pi_I = -\alpha^2 \sum_{i,j=1,2}\left\{\frac{\Gamma_{ij}''}{2}[B_i(t)B_j(t), \pi_I]_+ + \right. \\ \left. + i\Delta E_{ij}''[B_i(t)B_j(t), \pi_I] - \Phi_{ij}B_i(t)\pi_I B_j(t)\right\}, \tag{7}$$

where

$$
(8)\quad \left\{
\begin{aligned}
&\frac{\Gamma''_{ij}}{2} + i\,\Delta E''_{ij} = \left[\frac{|\langle J_e\|S\|J_f\rangle|^2}{2J_e+1}\,\frac{q^2}{2\hbar^2}\right]^2 \int f(\boldsymbol{v})\,\mathrm{d}\boldsymbol{v}\int f(\boldsymbol{v}')\,\mathrm{d}\boldsymbol{v}'\cdot\\
&\qquad\cdot\int I(\omega)\,\mathrm{d}\omega\,\frac{1}{\left(\omega-\omega_{ei}-\boldsymbol{K}\cdot\boldsymbol{v}+i(\Gamma/2)\right)\left(\omega-\omega_{ej}-\boldsymbol{K}\cdot\boldsymbol{v}'+i(\Gamma/2)\right)}\,,\\
&\Phi_{ij} = \Phi^r_{ij} + i\Phi^i_{ij} = \left[\frac{|\langle J_e\|S\|J_f\rangle|^2}{2J_e+1}\,\frac{q^2}{2\hbar^2}\right]^2 \int f(\boldsymbol{v})\,\mathrm{d}\boldsymbol{v}\int f(\boldsymbol{v}')\,\mathrm{d}\boldsymbol{v}'\cdot\\
&\qquad\cdot I(\omega)\,\mathrm{d}\omega\,\frac{1}{\left(\omega-\omega_{ei}-\boldsymbol{K}\cdot\boldsymbol{v}+i(\Gamma/2)\right)\left(\omega-\omega_{ej}-\boldsymbol{K}\cdot\boldsymbol{v}'-i(\Gamma/2)\right)}\,,
\end{aligned}\right.
$$

ω_{ei} being the energy separation between the excited state and the hyperfine sublevel F_i. It is worth-while to emphasize that the expression for the second-order variation of the polarization matrix implies a behaviour of the multiple-scattering terms very similar to the first-order ones, with a commutator and an anticommutator. On the contrary the interference term for weak scattering waves shows a completely different behaviour. We find to the second order the same operators $B_i(t)$ already examined up to the first order and from (7) we obtain that the principal polarizations are always the left and right circular ones.

For circularly polarized light, the second-order contribution to the absorption of the monitoring beam is

$$
(9)\qquad \Delta I = \alpha^2 I \sum_{i,j=1,2} (\Gamma''_{ij} - \Phi^r_{ij})\left\{\frac{1}{3}\langle T^{(0)}_0\rangle_i\langle T^{(0)}_0\rangle_j \pm \pm\frac{1}{\sqrt{6}}\left[\langle T^{(0)}_0\rangle_i\langle T^{(1)}_0\rangle_j + \langle T^{(0)}_0\rangle_j\langle T^{(1)}_0\rangle_i\right] + \frac{1}{2}\langle T^{(1)}_0\rangle_i\langle T^{(1)}_0\rangle_j\right\},
$$

where the $\pm$ signs account for the two circular polarizations. The first term depends only on the hyperfine-level populations and does not change in the magnetic resonance. The second term shows a behaviour equal to the first-order absorption: it gives an absorption at the Larmor precession frequency for the two circular polarizations. The third term is independent of the polarization and contains the square of the oscillating magnetization.

Let us now examine the Faraday rotation of a linearly polarized monitoring light by using an analyser which is at an arbitrary angle θ to the polarizer.

For the transmitted light we have

$$
(10)\qquad \Delta I = \alpha^2 I \sum_{i,j=1,2} \{(\Gamma''_{ij} - \Phi^r_{ij})\cos^2\theta\,\tfrac{1}{3}\langle T^{(0)}_0\rangle_i\langle T^{(0)}_0\rangle_j + + \Delta E''_{ij}\cos\theta\sin\theta\sqrt{\tfrac{2}{3}}\left(\langle T^{(0)}_0\rangle_i\langle T^{(1)}_0\rangle_j + \langle T^{(0)}_0\rangle_j\langle T^{(1)}_0\rangle_i\right) + + (\Gamma''_{ij}\cos^2\theta - \Phi^r_{ij}\sin^2\theta)\langle T^{(1)}_0\rangle_i\langle T^{(1)}_0\rangle_j\}\,.
$$

The first term depends only on the spin populations. The second term is the analogue of the first-order rotation and gives a Faraday effect at the Larmor precession frequency. The third term, dependent on M_x^2, has two parts of opposite sign and of the same order (by (8), Φ_{ij}^r is greater than Γ_{ij}''). Then this term has a maximum for $\theta = 0$, a zero for $\theta \simeq \pi/4$ and a minimum for $\theta = \pi/2$, the same as for the second-order term of (1).

The characteristic term of the second order is proportional to M_x^2, that is [9]

$$M_x^2 = 4H_1^2(\chi' \cos \Omega t + \chi'' \sin \Omega t)^2 =$$
$$= 2H_1^2(\chi'^2 + \chi''^2) + 2H_1^2(\chi'^2 - \chi''^2) \cos 2\Omega t + 4H_1^2 \chi' \chi'' \sin 2\Omega t \,.$$

Then a static signal of the type

$$\frac{H_1^2}{2}\chi_0^2 \frac{\Omega_0^2 T_2^2[1 + (\Delta\Omega T_2)^2]}{[1 + (T_2\Delta\Omega)^2 + \gamma^2 H_1^2 T_1 T_2]^2} \tag{11}$$

and an oscillating signal at a frequency 2Ω of the type

$$\frac{H_1^2}{2}\chi_0^2 \frac{\Omega_0^2 T_2^2[(\Delta\Omega T_2)^2 - 1]}{[1 + (\Delta\Omega T_2)^2 + \gamma^2 H_1^2 T_1 T_2]^2} \cos 2\Omega t + H_1^2 \chi_0^2 \frac{\Omega_0^2 T_2^2 \Delta\Omega T_2}{[1 + (\Delta\Omega T_2)^2 + \gamma^2 H_1^2 T_1 T_2]^2} \sin 2\Omega t \tag{12}$$

result.

For low r.f. field H_1 the static signal depends on Ω as an absorption curve whereas, at stronger H_1 fields, it has the shape of a square of a dispersion curve. Moreover for strong H_1 intensities, the $\cos 2\Omega t$ signal depends on Ω as the square of a dispersion curve, but for lower r.f. fields it turns out to be a mixture of that curve with an absorption curve.

3. – Experimental method.

A schematic picture of the experimental apparatus is shown in Fig. 2. A pyrex glass cell of diameter 10 cm, placed in a oven at temperature 120 °C, contained saturated sodium vapour and a few Torr of Ar as buffer gas. D_1 pumping light was produced by a 50 MHz radiofrequency excited Na lamp filtered by means of a quartz Lyot filter and circularly polarized by a quarter-wave plate. The pumping light trasmitted through the absorption cell was monitored by a photocell and displayed on a dual-beam oscilloscope. The

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[9] A. ABRAGAM: *The Principles of Nuclear Magnetism* (Oxford, 1961).

absorption cell was located at the centre of a pair of 50 cm diameter Helmholtz coils, used to produce a static magnetic field parallel to the light beam. This field was swept linearly over the resonances by the sawtooth voltage of the oscilloscope horizontal deflection plates. A second pair of coils, at right angles to the first, supplied the r.f. magnetic field and was driven by a 3.90 MHz quartz oscillator.

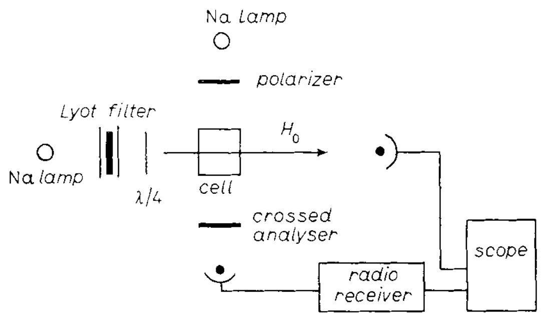

Fig. 2. – Experimental apparatus in the arrangement with crossed polarizers.

The monitoring transversal beam, of about 2 cm² cross-section, was produced by a radiofrequency excited Na lamp. A Lyot filter could select the D_1 or D_2 component, but a good signal-to-noise ratio was obtained also without the filter (2,10). After crossing the cell the beam passed through an analyser and was detected by a 2P2I (RCA) photomultiplier. The output was amplified in a frequency bandwidth of 13 kHz by a radio receiver tuned at the double of the quartz oscillator frequency. At the input of the radio a second-harmonic signal from the r.f. generator was also present so that phase detection on the radiofrequency could be achieved.

The radio output was finally displayed on the second beam of the oscilloscope. The cross-beam light was polarized and analysed by means of various linear and circular polarizers.

4. – Experimental results.

Investigations have been firstly carried out on the dependence of the second-order effects on the polarization of the monitoring light beam and on the detection technique. Thereafter the line shapes of the magnetic resonances at the modulation frequencies have been examined for a wide range of r.f. power.

1) A modulation at twice the Larmor frequency was observed in monitoring light absorption for any kind of polarization of the transversal beam. This is due to the fact that both the two circular polarizations contribute equally to the modulation of light beam as theoretically foreseen in the previous Section.

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(10) A. L. Bloom: *Journ. Phys. Rad.*, **19**, 881 (1958).

In order to observe the transversal Faraday effect, linearly polarized light has been used, the transmitted light being observed through a linear analyser and detected by the photomultiplier. In particular, it must be stressed that a signal at twice the Larmor frequency can be obtained either with a crossed or with a parallel analyser. As has been illustrated theoretically above, the two signals have an opposite phase relationship, the intensity being weaker with parallel than with crossed polaroids.

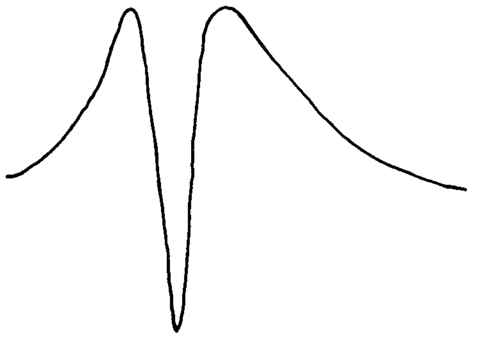

Fig. 3. – Transversal Faraday effect at twice the Larmor frequency as observed on the $(2, -2) \leftrightarrow (2, -1)$ transition for very low r.f. field intensity.

As in any case the same observables are monitored, the effects on the beam can be observed in any desired particular arrangement. In the investigations reported below, the arrangement with crossed polarizer and analyser was chosen, since it gives the best signal-to-noise ratio.

2) Figure 3 shows the $(2, -2) \leftrightarrow (2, -1)$ transition line at a frequency 2Ω as observed with σ^- as pumping light and at very low r.f. power. Keeping into account that the radio receiver acts as a phase detector, this curve follows exactly the theoretical form as calculated for the $\cos 2\Omega t$ term at small values of H_1.

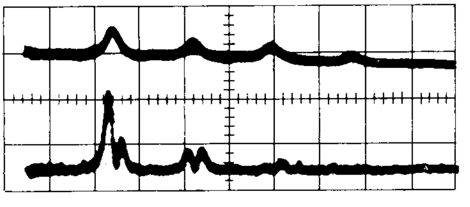

Fig. 4. – $\Delta m = 1$ transitions observed in the intensity of the pumping light (upper trace) and in the double frequency Faraday effect (lower trace).

Figure 4 shows the whole pattern of lines for stronger r.f. field as displayed on the dual-beam oscilloscope. The upper trace reports the lines as obtained in observing the absorption on the pumping light along the H_0 direction, the lower trace shows the $\cos 2\Omega t$ r.f. signals at the radio output. The four lines due to the transitions between $F = 2$ sublevels are superposed on those due to the $F = 1$ ground state. The double-frequency modulation lines show the form of a square of a dispersion signal.

At higher r.f. field intensities broadening of single photon lines occurs and, as shown in Fig. 5, more intense and narrow two-photon transitions appear. The behaviour of the signals in the r.f. pattern is rather complex and can be described only in qualitative terms. In the two-photon transitions with $\Delta m = 2$, Zeeman coherences of the type $\sigma'_{m,m-2}$ are produced, but owing to the fact that the $B_i(t)$ do not depend on them, these terms are not observable. Nevertheless Zeeman coherences of small amplitude of the type $\sigma^f_{m,m-1}$ and $\sigma^f_{m-1,m-2}$ result [11]. These terms contribute to $\langle T^{(1)}_0 \rangle_i$ and allow the observation in the monitoring transversal beam of weak signals due to the two-photon transitions. On the other hand $\Delta m = 1$ transitions overlap so remarkably that only a single lobe of the expected r.f. curve is recognizable on the less perturbed side of the line pattern.

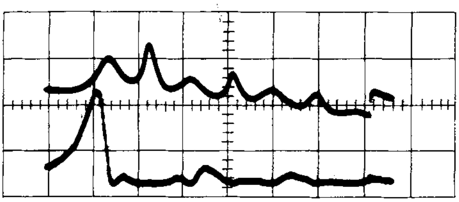

Fig. 5. – $\Delta m = 1$ and $\Delta m = 2$ resonances. The r.f. Faraday effect in the lower trace.

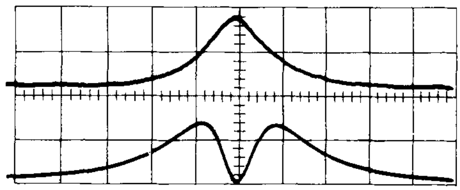

Fig. 6. – Magnetic resonance between Zeeman sublevels for very strong r.f. field as observed in pumping light (upper trace) and in the Faraday effect (lower trace).

For an overly strong r.f. field the different Zeeman transition lines are broadened and overlap so widely that the whole system behaves like a single-spin assembly: the r.f. signal takes again its simple form of the square of a dispersion curve (Fig. 6).

Static Faraday rotation has been observed by monitoring directly on the oscilloscope the d.c. response of the photomultiplier. Figure 7 shows in the lower trace this d.c. response: owing to the intensity of the r.f. field, the static signal has the behaviour of the square of a dispersion curve, and single- and two-photon lines show the same characteristics of Fig. 5.

Fig. 7. – The output of the photomultiplier directly displayed on the oscilloscope (lower trace).

5. – Conclusion.

The higher-order analysis of the polarization matrix of a monitoring light beam propagating through an oriented alkali vapour produces the harmonics

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([11]) K. Rosiński: *Acta Phys. Pol.*, **31**, 107, 173 (1967).

of the Larmor frequencies of the spins. Nevertheless the light beam always detects the observables of the first order, although they are examined in a different way.

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The authors wish to thank Dr. P. MINGUZZI for many helpful discussions on the argument.

RIASSUNTO

Utilizzando un fascio di luce trasversale al campo magnetico statico, sono stati esaminati gli effetti del secondo ordine dovuti alla risonanza magnetica in atomi di sodio orientati per mezzo del pompaggio ottico. Per tener conto di questi effetti la teoria di Cohen-Tannoudji e Laloë è stata estesa. I risultati sperimentali sono confrontati con quelli teorici ottenendo pieno accordo.

Модуляция второго порядка для поперечного светового пучка посредством магнитного резонанса в ориентированных парах Na.

Резюме (*). — С помощью мониторной техники поперечного пучка были исследованы эффекты второго порядка, обусловленные магнитным резонансом атомов натрия, ориентированных посредством оптической накачки. Было проведено расширение теории Кохена-Танноуджи и Лалоё, чтпбы включить эти эффекты, и ее применения сравниваются с экспериментальными результатами.

(*) *Переведено редакцией.*