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New possible manifestations of parity violation in forbidden magnetic transitions : chiral absorption of plane polarized light and circular dichroism in crossed transverse dc electric and magnetic fields

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Résumé. — Dans ce papier, nous présentons de nouvelles méthodes d'attaque du problème de l'observation d'une violation de la parité dans les transitions magnétiques interdites d'atomes lourds. Le principe de base est l'observation d'une interférence entre l'amplitude dipolaire électrique induite par l'interaction avec courants neutres violant la parité et l'amplitude dipolaire électrique induite par un champ électrique statique, dans les transitions à un seul photon $nS \rightarrow n' S$. Le caractère nouveau des expériences discutées ici est qu'elles nécessitent un champ magnétique statique de l'ordre de un kilogauss perpendiculaire au champ électrique. La grandeur mesurée est la section efficace d'absorption en lumière polarisée plane ou circulaire. L'asymétrie droite-gauche est discutée dans plusieurs cas simples où la nature physique de l'effet a une interprétation claire, puis on présente le résultat de calculs numériques correspondant à des conditions expérimentales réalistes. Les signaux de violation de la parité ont, dans tous les cas, une signature bien définie. Les expériences proposées dans ce papier constituent des alternatives intéressantes par rapport aux expériences actuellement en cours, qui sont réalisées en champ magnétique nul.

Abstract. — In this paper, we shall present new approaches to the problem of finding a parity violation in forbidden magnetic transitions of heavy atoms. The basic principle is the observation of an interference between the electric dipole amplitude induced by the parity violating neutral currents interaction and the stark induced electric dipole amplitude in $nS \rightarrow n' S$ radiative transitions. The new features of the experiments discussed here is that they have to be performed in a d.c. magnetic field, of the order of one kilogauss, crossed with the d.c. electric field. The experiments involve the measurement of the absorption cross-section of plane or circularly polarized light. The left-right asymmetries are discussed in some simple cases where the physical nature of the effect is clearly exhibited ; the results of numerical computations corresponding to realistic experimental conditions are presented. The parity violating signals have in all cases a well defined signature. The experiments proposed in this paper constitute interesting alternatives to experiments now under way, which are performed in zero magnetic d.c. field.

The study of parity violation in atoms has recently become a way to get information about the structure of weak neutral currents. The study of forbidden magnetic one-photon transitions in monovalent heavy atoms [1, 2], among several possible experimental approaches [3], looks rather promising. The experiments on Cesium [4] and Thallium [5], which are

under completion, involve the observation of an interference effect between the electric dipole amplitude induced by the parity violating neutral currents and the electric dipole amplitude $E_1^{\text{ind.}}$ induced by the static electric field E_0 . When the transition is excited with a beam of photons of momentum \mathbf{k} and circular polarization ξ , the $E_1^{\text{p.v.}} E_1^{\text{ind.}}$ interference gives rise to an electronic polarization \mathbf{P}_e in the final atomic state of the following form :

$$\mathbf{P}_e = 2p \frac{M_1 + \xi \text{Im} E_1^{\text{p.v.}}}{E_1^{\text{ind.}}} \hat{\mathbf{k}} \wedge \hat{E}_0, \quad (1)$$

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where M_1 is the magnetic dipole amplitude and \hat{k} , \hat{E}_0 unit vectors along \mathbf{k} and \mathbf{E}_0 ; the numerical factor p , a few tenths of unity, involves the nuclear spin and depends on the hyperfine transition under consideration. The $E_1^{p.v.}$ $E_1^{ind.}$ interference alone contributes to the ξ -dependent part of \mathbf{P}_e , and the pure-vectorial contribution $\xi \hat{k} \wedge \hat{E}_0$ to this angular momentum is a clear indication of a P -odd T -even part in the atomic hamiltonian. In the present paper, we want to show that, if a static magnetic field \mathbf{H} is applied along a direction perpendicular to the electric field \mathbf{E}_0 , new $E_1^{p.v.}$ $E_1^{ind.}$ interference effects appear. They are of the same order of magnitude as the one described by eq. (1) provided that the magnetic field is large enough (about one kilogauss); by symmetry considerations they can be shown to cancel out in zero magnetic field. They involve physical quantities which could be measured more conveniently or more reliably than \mathbf{P}_e and they deserve some interest not only from a pure theoretical point of view but also from an experimental one. In this paper, we shall not give a complete discussion of all the possible parity violating effects in crossed electric and magnetic fields, but limit ourselves to two types of experiments which constitute very interesting alternatives to experiments with $H = 0$.

The first experiment will consist in the measurement of a circular dichroism associated with the presence of the pseudoscalar $\xi \mathbf{k} \cdot \mathbf{E}_0 \wedge \mathbf{H}$ in the resonant absorption cross-section σ_a . The second effect, which could be described as *chiral* absorption of plane polarized light, involves the presence in σ_a of a pseudoscalar contribution which is quadratic in the magnetic field: $(\boldsymbol{\varepsilon} \cdot \mathbf{H})(\boldsymbol{\varepsilon} \cdot \mathbf{H} \wedge \mathbf{E}_0)$. From a theoretical point of view the two effects differ in the fact that the chiral absorption involving a plane polarized wave vanishes in the limit of L.S. coupling, while the circular dichroism stays finite in the same limit. This property, which applies only to nS - $n'S$ transitions where the effect of an external magnetic field \mathbf{H}_0 can be accounted for by a Zeeman Hamiltonian involving only the electronic spin \mathbf{S} , will be obtained from symmetry considerations in Appendix II.

Note that the two pseudoscalar quantities $\xi \mathbf{k} \cdot \mathbf{E}_0 \wedge \mathbf{H}$ and $(\boldsymbol{\varepsilon} \cdot \mathbf{H})(\boldsymbol{\varepsilon} \cdot \mathbf{H} \wedge \mathbf{E}_0)$ are even under the transformation $\mathbf{E}_0 \rightarrow \mathbf{E}_0$, $\mathbf{H} \rightarrow -\mathbf{H}$, $\mathbf{k} \rightarrow -\mathbf{k}$. This property reflects here the invariance of the parity violating Hamiltonian under time reflexion. In Appendix I, we show how possible rotational invariants, odd under space reflexion and even under time reflexion, can be constructed from the physical quantities which specify the experimental set-up. In the case of metastable 2S Hydrogen state parity violating effects in crossed electric and magnetic fields have also been considered [6]. The pseudo-scalars involved are superficially odd under time-reversal. This is due to the fact that the near degeneracy of the 2S state with the *unstable* 2P state prevents a straightforward application of time reflexion invariance.

We recall that one of the interests in looking for an interference between two electric dipole amplitudes ($E_1^{p.v.}$ and $E_1^{ind.}$), instead of an $E_1^{p.v.}$ M_1 interference (optical rotation experiment [3]) is the fact that the interference does not change sign under a normal mirror reflection of the beam allowing amplification by a multipass device.

We will first recall some of the properties of a forbidden magnetic dipole radiative transition between two $S_{1/2}$ states in an electric field (sect. 1) and then, in sections 2 and 3, we will discuss respectively circular dichroism and chiral absorption experiments, in crossed electric and magnetic fields. In each case, the experimental situation will be compared to the one currently under use in zero magnetic field with special concern about signal to noise ratio and discrimination against spurious effects. In Appendix I and II, symmetry considerations leading to the effects discussed in sections 2 and 3 are presented.

1. Forbidden $nS_{1/2} \rightarrow n'S_{1/2}$ transition in an external d.c. electric field : a review of the formalism. —

For an incident photon in a state of arbitrary polarization described by the complex vector $\boldsymbol{\varepsilon}$ of unit norm, the $nS_{1/2}$ - $n'S_{1/2}$ transition amplitude, for a one-valence-electron atom, can be obtained from the effective transition matrix T , acting only on the electronic spin states :

$$T_{m_S m'_S} = \langle n' S, m'_S | -\boldsymbol{\varepsilon} \cdot \mathfrak{D} - \mathbf{k} \wedge \boldsymbol{\varepsilon} \cdot \mathfrak{M} | n S, m_S \rangle,$$

or

$$T = a\mathbb{1} + \mathbf{b} \cdot \boldsymbol{\sigma}. \quad (2)$$

\mathfrak{D} represents the electric dipole operator and \mathfrak{M} the magnetic dipole operator. $a = \alpha \boldsymbol{\varepsilon} \cdot \mathbf{E}_0$ is the Stark induced electric dipolar contribution associated with the scalar polarizability α . The components of $\boldsymbol{\sigma}$ are the three Pauli matrices and

$$\mathbf{b} = i\beta \mathbf{E}_0 \wedge \boldsymbol{\varepsilon} - M_1 \mathbf{k} \wedge \boldsymbol{\varepsilon} + i \text{Im} E_1^{p.v.} \boldsymbol{\varepsilon}. \quad (3)$$

Three different contributions appear : one comes again from the Stark induced electric dipole, but now through the vectorial part β of the tensor polarizability (¹), and the others arise from the magnetic and the parity violating electric dipoles. Explicit expressions of α , β and the parity violating amplitude

$$E_1^{p.v.} = \langle nS_{1/2} | dz | n'S_{1/2} \rangle$$

in terms of one-particle wave functions can be found in ref. [2].

It is shown that $\beta = 0$ in the limit of L-S coupling, i.e. if the spin orbit interaction is neglected. The ratio α/β has been found experimentally to be equal to -8.8 for the $6S_{1/2}$ - $7S_{1/2}$ transition of Cesium [4].

(¹) The selection rules associated with the α and β Stark induced transitions have been studied experimentally by M. A. Bouchiat and L. Pottier, *J. Phys. Lett.* **36** (1975) L-189.

M_1 is the magnetic dipolar amplitude which is also known experimentally for this transition [7], in terms of the induced electric dipolar amplitude

$$E_1^{ind.} = -\alpha E_0 \cdot \epsilon : \\ M_1/E_1^{ind.} = - (2.62 \pm 0.21)/E_0 \text{ (V/cm)}. \quad (4)$$

$E_1^{ind.}$ itself has been theoretically evaluated with an accuracy of about 15 % [2] :

$$E_1^{ind.} = 5.93 \times 10^{-8} |e| a_0 E_0 \cdot \epsilon \text{ (V/cm)}. \quad (5)$$

From the value computed for $E_1^{p.v.}$ within the framework of the Weinberg-Salam model [2], taking $\sin^2 \theta_w = 0.25$, one can expect :

$$\text{Im } E_1^{p.v.}/M_1 = 0.86 \times 10^{-4}. \quad (6)$$

2. Circular dichroism experiments in crossed electric and magnetic fields. — We shall first exhibit the physical nature of the parity violating effect under search and then consider as an illustration the low field limit where the calculations are particularly simple. Finally, we shall give the results of a numerical computation concerning the realistic case of the $6S_{1/2} \rightarrow 7S_{1/2}$ transition in Cesium where the experiment looks feasible.

2.1 PHYSICAL NATURE OF THE PARITY VIOLATING SIGNAL. — Let us take the photon momentum along the x axis and E_0 along the y axis, as shown on figure 1.

The polarization vector is written $\epsilon = \frac{1}{\sqrt{2}}(\hat{y} + i\xi\hat{z})$,

\hat{y} and \hat{z} being unit vectors directed along y and z axis. The cases $\xi = +1$ and -1 correspond to an helicity $+1$ and -1 of the incident photon. From eq. (2), we get for the transition matrix :

$$T = \frac{1}{\sqrt{2}} [\alpha E_0 - \xi\beta E_0 \sigma_x - (M_1 + \xi \text{Im } E_1^{p.v.}) \times \\ \times (\sigma_z - i\xi\sigma_y)]. \quad (7)$$

Let us momentarily ignore the nuclear spin. There are only four possible transitions from the ground $|n \pm 1/2\rangle$ to the excited $|n' \pm 1/2\rangle$ Zeeman substates. The corresponding probabilities and frequencies are presented in table I. $\omega_{n'n}$ denotes the transition frequency in zero magnetic field, $\omega_0 = 2\mu_B H$ is the Zeeman splitting.

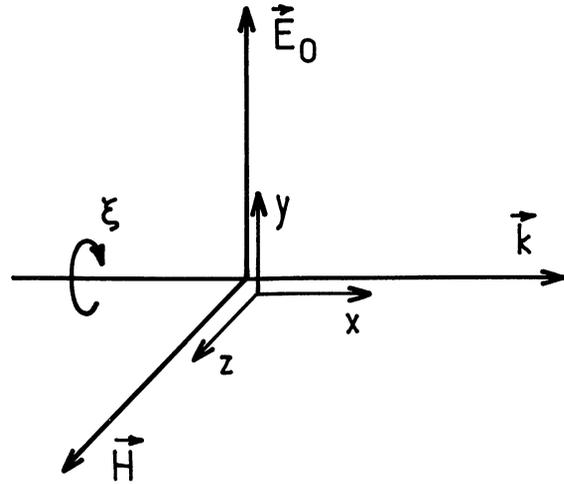


Fig. 1. — Geometrical configuration of the circular dichroism experiment in crossed transverse electric and magnetic fields.

An interference term between the Stark induced $E_1^{ind.}$ amplitude and the mixed $E_1^{p.v.}-M_1$ amplitude clearly appears in each transition probability. A circular dichroism (ξ dependence) is associated with the interference $E_1^{ind.} E_1^{p.v.}$. We have plotted in figure 2 the frequency dependence of interference terms associated with the four transitions, assuming an arbitrary line shape and taking $\alpha/\beta = -2$. It is seen that the Zeeman effect suppresses the cancellation of the interference terms for the $\Delta m_s = +1$ and -1 transitions induced by the spin dependent dipolar amplitude βE_0 but fails to do so for the $\Delta m_s = 0$ transitions

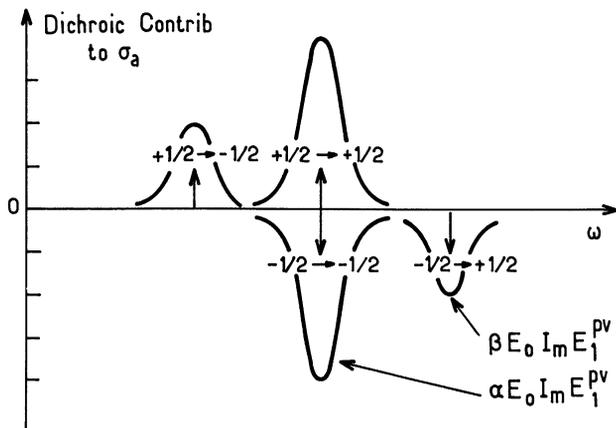


Fig. 2. — Frequency dependence for the interference effects in the absorption probabilities (no nuclear spin, α/β arbitrarily chosen equal to -2).

Table I.

| Transition | | Probability $\times 2$ | Frequency |
|----------------------|-------------------|---|-----------------------------|
| upper state | lower state | | |
| $\langle n', +1/2 $ | $ n, +1/2\rangle$ | $ \alpha E_0 - (M_1 + \xi \text{Im } E_1^{p.v.}) ^2$ | $\omega_{n'n}$ |
| $\langle n', +1/2 $ | $ n, -1/2\rangle$ | $ \beta E_0 + \xi(M_1 + \xi \text{Im } E_1^{p.v.}) ^2$ | $\omega_{n'n} + \omega_0/2$ |
| $\langle n', -1/2 $ | $ n, +1/2\rangle$ | $ \xi\beta E_0 + \xi(M_1 + \xi \text{Im } E_1^{p.v.}) ^2$ | $\omega_{n'n} - \omega_0/2$ |
| $\langle n', -1/2 $ | $ n, -1/2\rangle$ | $ \alpha E_0 + (M_1 + \xi \text{Im } E_1^{p.v.}) ^2$ | $\omega_{n'n}$ |

which are not affected by the magnetic field. We note that the resulting dichroism will have a dispersion line shape centered at the zero-field frequency.

In the above discussion, we have ignored the effect of the nuclear spin \mathbf{I} . As we shall see below, the hyperfine interaction modifies the previous results in an essential way. To first order in H the α -term still induces only transitions with $\Delta F = \Delta M_F = 0$: the operator $\alpha \mathbf{E}_0 \cdot \boldsymbol{\varepsilon}$ acting like a scalar cannot modify the total angular momentum $\mathbf{F} = \mathbf{I} + \mathbf{S}$. These transitions occur at the same frequency, the zero-field frequency, up to second order in H . But a first order perturbation of the initial and final spin wave functions gives rise to a difference between the electric and magnetic dipole transition probabilities associated with opposite values of M_F :

$$|(M_1 + E_1^{p.v.})_{M_F \rightarrow M_F}| \neq |(M_1 + E_1^{p.v.})_{-M_F \rightarrow -M_F}|.$$

As a result, the interference effects in $\Delta M_F = 0$ transitions no longer cancel out. In fact, since in the actual situation of the $6S_{1/2} \rightarrow 7S_{1/2}$ Cesium transition $|\alpha/\beta| = 8.8$, the $\alpha E_0 E_1^{p.v.}$ interference is dominant. Note that its frequency dependence will be of the absorptive type.

2.2 THE LOW MAGNETIC FIELD LIMIT WITH NUCLEAR SPIN EFFECTS. — We now take into account the hyperfine coupling of \mathbf{S} with a nuclear spin \mathbf{I} : A and A' are the hyperfine coupling constants of the $nS_{1/2}$ and $n'S_{1/2}$ states. For an arbitrary value of the magnetic field, we shall call $|nF_\varepsilon m_F\rangle$, $|n'F_\varepsilon m_F\rangle$ (with $\varepsilon = +1$ or -1) the initial and final atomic states which, in the limit $H \rightarrow 0$, reduce to eigenstates of $\mathbf{F}^2 = (\mathbf{I} + \mathbf{S})^2$ and $F_z = I_z + S_z$, with eigenvalues $F_\varepsilon = I + \frac{\varepsilon}{2}$ and m_F . If we make the simplification $\beta = 0$, only transitions with $\Delta m_F = 0$ are possible. We have for instance:

$$\begin{aligned} & |\langle n'F_\varepsilon m_F | T | nF_\varepsilon m_F \rangle|^2 = \\ &= \frac{1}{2} \alpha^2 E_0^2 \cos^2 \left(\frac{\theta_m - \theta'_m}{2} \right) - \frac{\varepsilon}{2} \hat{k} \cdot (\hat{E}_0 \wedge \hat{H}) \times \\ & \quad \times \alpha E_0 (M_1 + \xi \text{Im } E_1^{p.v.}) (\cos \theta_m + \cos \theta'_m), \quad (8) \end{aligned}$$

\hat{k} , \hat{E}_0 , \hat{H} being unit vectors along \mathbf{k} , \mathbf{E}_0 , \mathbf{H} . In eq. (8), terms of order $(M_1^2/\alpha^2 E_0^2) \ll 1$ have been neglected. The mixing angle θ_m is given by:

$$\text{tg } \theta_m = \frac{\sqrt{(I+1/2)^2 - m_F^2}}{m_F + \frac{\omega_0}{A}}, \quad \text{with } 0 \leq \theta_m \leq \pi. \quad (9)$$

To get θ'_m , A has to be replaced by A' .

In the limit of zero field

$$\theta_m = \theta'_m \quad \text{and} \quad \cos \theta_m = \frac{m}{I+1/2}.$$

The initial state being unpolarized, one has to average over m_F with equal weight (no Zeeman splitting) and the term involving $\hat{k} \cdot (\hat{E}_0 \wedge \hat{H})$ averages to zero, as expected.

To first order in \mathbf{H} , the energy differences

$$E(n'F_\varepsilon m_F) - E(nF_\varepsilon m_F)$$

remain independent of m_F . Consequently to compute the excitation cross-section, one has still to average over the values of m_F with equal weights. The quantity $\cos \theta_m + \cos \theta'_m$ is no longer proportional to m_F but is now a quadratic function of m_F :

$$\cos \theta_m = \frac{1}{I+1/2} \left[m_F + \frac{\omega_0}{A} \left(1 - \frac{m_F^2}{(I+1/2)^2} \right) + \mathcal{O}(\omega_0)^2 \right]. \quad (10)$$

The average of $\cos \theta_m$ is now clearly non zero to first order in ω_0 . In the limit $\beta = 0$, the absorption cross-section at the resonance peak associated with the transitions $nS_{1/2} \rightarrow n'S_{1/2}$ with $\Delta F = 0$ is given by:

$$\sigma_a(H) = \bar{\sigma}_a \left[u(H) - 2v(H) \frac{M_1 + \xi \text{Im } E_1^{p.v.}}{\alpha E_0} \times \hat{k} \cdot \hat{E}_0 \wedge \hat{H} \right], \quad (11)$$

where $\bar{\sigma}_a$ is the resonance cross-section in zero H -field limit, summed over the two transitions $\Delta F = 0$. To first order in H , one obtains for the two $\Delta F = 0$ transitions:

$$F = I + 1/2 \rightarrow F = I + 1/2$$

$$u = \frac{I+1}{2I+1} \quad v = \frac{1}{3} \frac{I(I+1)}{(I+1/2)^2} \mu_B H \left(\frac{1}{\Delta W} + \frac{1}{\Delta W'} \right) \quad (12)$$

$$F = I - 1/2 \rightarrow F = I - 1/2$$

$$u = \frac{I}{2I+1} \quad v = -\frac{1}{3} \frac{I(I+1)}{(I+1/2)^2} \mu_B H \left(\frac{1}{\Delta W} + \frac{1}{\Delta W'} \right) \quad (13)$$

where $\Delta W = A(I+1/2)$ and $\Delta W' = A'(I+1/2)$ are the hyperfine splittings of the initial and final states. The function $v(H)$ will reach a value of the order of unity for magnetic fields of the order of one kilogauss. A first order computation is clearly no longer adequate. Moreover, the effect of Zeeman splittings on the absorption line shape and the Boltzmann distribution of the absorbed frequencies have to be taken into account.

2.3 RESULTS OF A COMPLETE CALCULATION FOR THE $6S_{1/2}$ - $7S_{1/2}$ CESIUM TRANSITION. — A numerical computation (on UNIVAC 1110) of the absorption cross-section as a function of the photon frequency ω has been performed for different magnitudes of the

magnetic field (500 G, 1 000 G) where the Zeeman splitting is no longer small compared to the hyperfine splitting

$$(\Delta W/h = 9\,192\text{ MHz}; \Delta W'/h = 2\,175\text{ MHz}).$$

The numbers correspond to a realistic experimental situation : a saturated Cesium vapor in thermal equilibrium at $T = 472\text{ K}$; the Doppler half-height linewidth is 750 MHz ; the incident laser field is quasi-monochromatic (single mode c.w. tunable laser). If we neglect terms of the order of $(M_1/\alpha E_0)^2$ the absorption cross-section can be written as :

$$\begin{aligned} \sigma_a(H, \omega) = \bar{\sigma}_a \left[u_\alpha(\omega, H) + \frac{\beta^2}{\alpha^2} u_\beta(\omega, H) - \right. \\ \left. - 2 \left(v_\alpha(\omega, H) + \frac{\beta}{\alpha} v_\beta(\omega, H) \right) \frac{M_1 + \xi \text{Im } E_1^{p.v.}}{\alpha E_0} \times \right. \\ \left. \times \hat{k} \cdot \hat{E}_0 \wedge \hat{H} \right]. \quad (14) \end{aligned}$$

The numerical values of coefficients u_α , u_β , v_α and v_β are plotted *versus* the laser frequency ω for two field magnitudes $H = 500\text{ G}$ and $1\,000\text{ G}$, on figures 3 and 4. Note that $2 v_\alpha$ reaches values of the order of u_α for several frequencies. So the ξ -dependent part of σ_a , the parity violating circular dichroism :

$$\frac{\sigma_a^{p.v.}}{\bar{\sigma}_a} \simeq -2 \frac{v_\alpha + \frac{\beta}{\alpha} v_\beta}{u_\alpha} \cdot \frac{\text{Im } E_1^{p.v.}}{\alpha E_0} \xi \hat{k} \cdot \hat{E}_0 \wedge \hat{H} \quad (15)$$

reaches roughly the same value as the parity violating electronic polarization discussed in the introduction, i.e. the ξ -dependent part of \mathbf{P}_e in eq. (1) :

$$\mathbf{P}_e^{p.v.} = -2p \frac{\text{Im } E^{p.v.}}{\alpha E_0} \xi \hat{k} \wedge \hat{E}_0 \quad (16)$$

with $p = 5/12$ for the $4 \rightarrow 4$ transition and $p = 1/4$ for the $3 \rightarrow 3$ transition.

It can be shown that the v_α curve is odd with respect

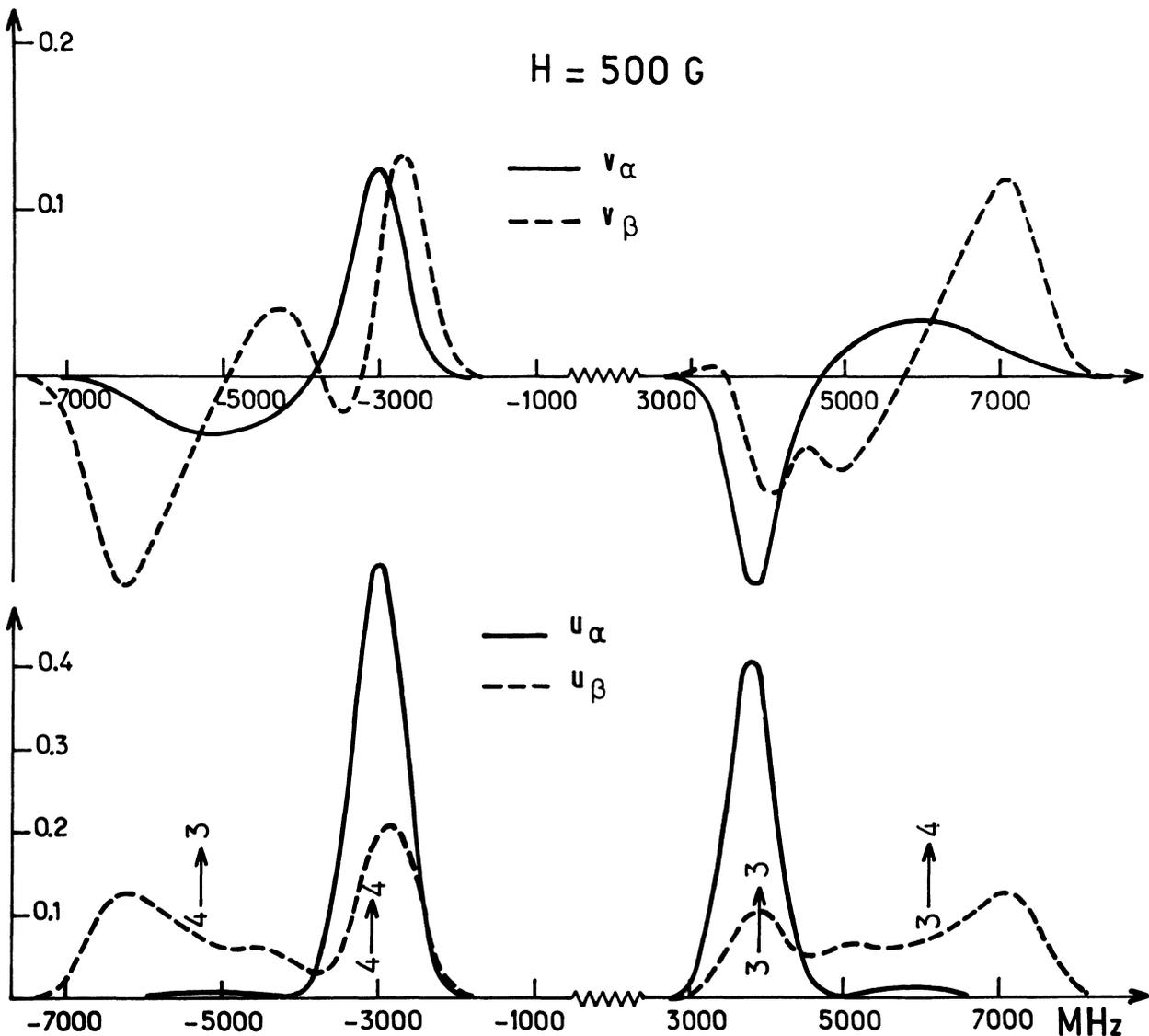
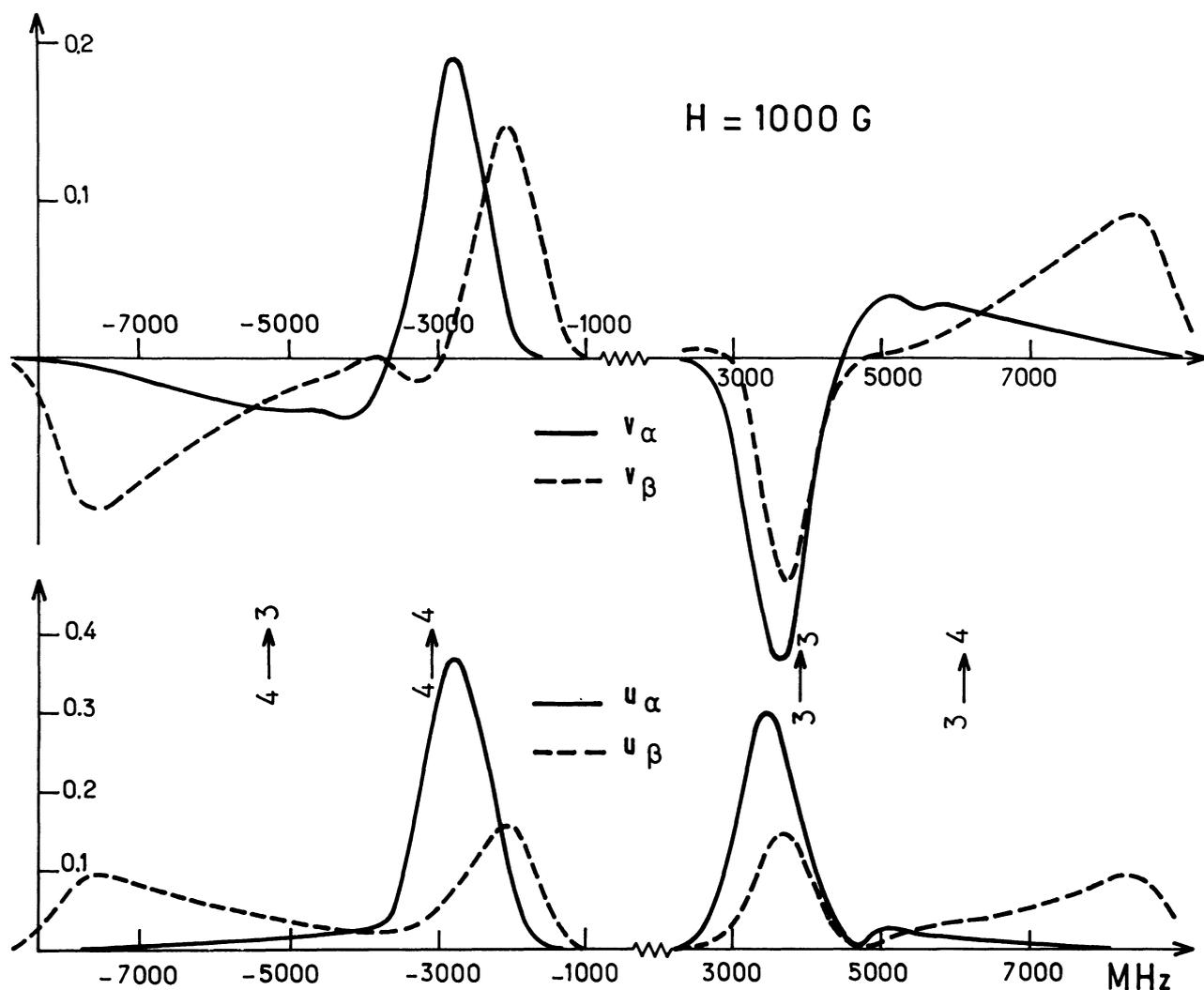


FIG. 3.



Figs. 3 and 4. — Absorption cross-section for the Cesium $6S_{1/2} \rightarrow 7S_{1/2}$ transition versus laser frequency, in crossed electric and magnetic fields with a circularly polarized beam.

Fig. 3 is relative to $H = 500$ G and fig. 4 to $H = 1000$ G

$$\sigma_a = \bar{\sigma}_a \left[u_\alpha + \frac{\beta^2}{\alpha^2} u_\beta - 2 \left(v_\alpha + \frac{\beta}{\alpha} v_\beta \right) \frac{M_1 + \xi \operatorname{Im} E_1^{p.v.}}{\alpha E_0} \hat{k} \cdot \hat{E}_0 \wedge \hat{H} \right].$$

Lower curves represent the main contributions : u_α and u_β . Upper curves give the dichroic contributions v_α and v_β .

The zero-field hyperfine components are indicated (3 \rightarrow 3, 3 \rightarrow 4, etc.). Zero frequency is the transition frequency without h.f.s.

to the frequency $(A - A')/4h$, the transition frequency averaged over the four hyperfine components.

The signal associated with the circular dichroism has several features which help its extraction from an eventual background :

i) Like $P_e^{p.v.}$, it is amplified by multipasses of the beam since $\xi \mathbf{k}$ does not change sign under reflexion of the beam at normal incidence. It should be remembered that this is not the case with the circular dichroism in zero electric and magnetic fields, associated with the $M_1 E_1^{p.v.}$ interference, first discussed in ref. [1]. Note also that the parity conserving term $E_1^{ind.} M_1 \hat{k} \cdot \hat{E}_0 \wedge \hat{H}$ is suppressed by multipasses of the beam.

ii) The change of sign which occurs when a reversal of \mathbf{E}_0 (or \mathbf{H}) is performed constitutes quite a good signature of the signal, more especially as the parity conserving term $M \hat{k} \cdot \hat{E}_0 \wedge \hat{H}$ also odd in \mathbf{E}_0 and \mathbf{H} , but not in ξ , is suppressed by multipasses of the beam. Thus a possible way of performing the experiment is to modulate E_0 and ξ at two different frequencies, ω_E and ω_i , and to monitor the p.v. signal at the frequencies $\omega_E + \omega_i$ and $\omega_E - \omega_i$ in the total fluorescence rate. With respect to the experiment currently under way which searches for $P_e^{p.v.}$ by monitoring the circularly polarized fluorescence rate, the removal of the circular analyzer should improve the efficiency of detection and consequently the signal to noise ratio by a significant factor.

iii) As it was apparent from the simplified analysis first presented, the circular dichroism changes with frequency. It will certainly help to eliminate spurious effects.

To conclude this section, we would like to discuss a circular dichroism linear in the electric field which is not associated with parity violation. In the geometrical configuration considered up to now, the magnetic field was taken to be perpendicular to the incident photon momentum \mathbf{k} . However, if \mathbf{H} acquires a small component along the momentum \mathbf{k} , a new term $\delta\sigma_a$ appears in the absorption cross-section :

$$\delta\sigma_a = \bar{\sigma}_a \xi(\hat{H} \cdot \hat{k}) (\hat{E}_0 \wedge \hat{H} \cdot \hat{k}) \frac{M_1}{\alpha E_0} g(H). \quad (17)$$

In other words, in a magnetic field having a component along the beam, parity conservation allows the M_1 - E_1^{ind} interference signal to take different values with incident circular polarizations of opposite signs. Even if a misalignment smaller than 10^{-3} radian cannot be achieved, this spurious circular dichroism can be reduced to an acceptable level using the two facts :

- i) $\delta\sigma_a$ is even under a reversal of \mathbf{H} while the true effect is odd.
- ii) $\delta\sigma_a$ is suppressed by multipasses of the light beam.

3. Chiral absorption of plane polarized light. —

We are concerned here with the search for parity violation through the measurement of the pseudoscalar $(\boldsymbol{\varepsilon} \cdot \mathbf{H})(\boldsymbol{\varepsilon} \cdot \mathbf{H} \wedge \mathbf{E}_0)$. For transitions between atomic states with $J = 1/2$ angular momentum, the above pseudoscalar is the only practical one if one wishes to use a plane polarized incident beam, which allows a multipass device with Brewster windows. As we shall see, another interesting feature of this new type of experiment is the fact that it is possible to find experimental conditions where the parity conserving E_1^{ind} M_1 interference vanishes while the parity violating term E_1^{ind} $E_1^{\text{p.v.}}$ stays near its maximum value.

If one forgets about the M_1 amplitude, the geometry of the experiment is entirely specified by two vectors (the linear polarization $\boldsymbol{\varepsilon}$ of the incident photon and the electric field \mathbf{E}_0) and one pseudovector (the magnetic field \mathbf{H}). Since we are interested in an interference term involving only the vectorial Stark induced amplitude $i\beta\boldsymbol{\sigma} \cdot \mathbf{E}_0 \wedge \boldsymbol{\varepsilon}$, we choose $\mathbf{E}_0 \cdot \boldsymbol{\varepsilon} = 0$. In order to keep the pseudoscalar at its maximum value, \mathbf{H} is taken to be normal to \mathbf{E}_0 and oriented at 45 degrees to the polarization vector $\boldsymbol{\varepsilon}$. The direction of the photon momentum \mathbf{k} , relevant only for the magnetic dipole amplitude can be chosen arbitrarily in the plane normal to $\boldsymbol{\varepsilon}$. In figure 4, we have drawn two mirror configurations with respect to a plane perpendicular to the magnetic field \mathbf{H} . The pseudoscalar $(\boldsymbol{\varepsilon} \cdot \mathbf{H})(\boldsymbol{\varepsilon} \cdot \mathbf{H} \wedge \mathbf{E}_0)$ takes — as it should — opposite values. We are going to show that the $\beta E_0 E_1^{\text{p.v.}}$ inter-

ference effect manifests itself by a contribution in the absorption cross-section with all transformation properties of the pseudoscalar $(\boldsymbol{\varepsilon} \cdot \hat{H})(\boldsymbol{\varepsilon} \cdot \mathbf{H} \wedge \mathbf{E}_0)$: it takes opposite signs in the two mirror configurations.

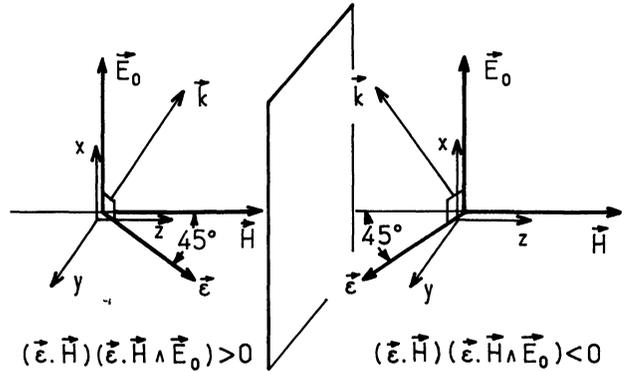


Fig. 5. — Chiral absorption experiment : image configurations with respect to reflexion by a mirror perpendicular to \mathbf{H} .

3.1 THE CHIRAL ABSORPTION IN A SIMPLE CASE : zero nuclear spin. — As an illustration, we shall first consider an $nS_{1/2} \rightarrow n'S_{1/2}$ transition in the simple case of zero nuclear spin. The effective transition matrix reduces to $T = \mathbf{b} \cdot \boldsymbol{\sigma}$ with \mathbf{b} given by eq. (3). Taking the quantization axis z along \mathbf{H} and defining b_{\pm} as :

$$b_{\pm} = b_x \pm ib_y.$$

The four transition probabilities between the atomic states $|nm_s\rangle$ and $|nm'_s\rangle$ ($m_s, m'_s = \pm 1/2$) are :

$$\begin{aligned} | \langle n' \pm 1/2 | T | n \pm 1/2 \rangle |^2 &= |b_z|^2 = (\mathbf{b} \cdot \hat{H})^2 \\ | \langle n \mp 1/2 | T | n \pm 1/2 \rangle |^2 &= |b_{\pm}|^2. \end{aligned} \quad (18)$$

Using the explicit expression of \mathbf{b} and keeping the first order term in M_1 and $E_1^{\text{p.v.}}$, one gets immediately :

$$\begin{aligned} | \langle n' \pm 1/2 | T | n \pm 1/2 \rangle |^2 &= \frac{1}{2} \beta^2 E_0^2 + \\ &+ 2 \beta E_0 \text{Im} E_1^{\text{p.v.}}(\boldsymbol{\varepsilon} \cdot \hat{H}) (\hat{E}_0 \wedge \boldsymbol{\varepsilon} \cdot \hat{H}), \\ | \langle n' \mp 1/2 | T | n \pm 1/2 \rangle |^2 &= \frac{1}{2} \beta^2 E_0^2 - \\ &- 2 \beta E_0 \text{Im} E_1^{\text{p.v.}}(\boldsymbol{\varepsilon} \cdot \hat{H}) (\hat{E}_0 \wedge \boldsymbol{\varepsilon} \cdot \hat{H}) \\ &\pm 2 \beta M_1(\boldsymbol{\varepsilon} \cdot \hat{H}) (\hat{E}_0 \wedge \boldsymbol{\varepsilon} \cdot \hat{k}). \end{aligned} \quad (19)$$

In the zero magnetic field limit where the four transitions occur at the same frequency, both the $\beta E_0 E_1^{\text{p.v.}}$ and $\beta E_0 M_1$ interference effects disappear. When \mathbf{H} is non zero, the $\Delta m_s = 0$ transition frequency remains unchanged while the $\Delta m_s = \pm 1$ ones are shifted by $\pm \omega_0/2$. Let us call $I(\omega)$ the line shape function characterized by a linewidth $\Delta\omega$. (In the limit $\Delta\omega \rightarrow 0$, $I(\omega) \rightarrow \delta(\omega)$.) The chiral part of the cross-section is proportional to the quantity :

$$\begin{aligned} 2 I(\omega - \omega_{nn'}) - I\left(\omega - \omega_{nn'} + \frac{\omega_0}{2}\right) - \\ - I\left(\omega - \omega_{nn'} - \frac{\omega_0}{2}\right) \end{aligned}$$

which in low magnetic field limit reduces to :

$$-\frac{\omega_0^2}{4} \frac{d^2 I}{d\omega^2} (\omega - \omega_{m'}) + \mathcal{O}(\omega_0^4). \quad (20)$$

In the simple case considered here of zero nuclear spin, the frequency dependence of the chiral part of the absorption cross-section for low magnetic field is given by the second derivative of the line shape function.

Concerning the $E_1^{\text{ind.}} M_1$ (or $\beta E_0 M_1$) interference term, we note that it involves the true scalar

$$(\boldsymbol{\varepsilon} \cdot \hat{H}) (\boldsymbol{\varepsilon} \cdot \hat{k} \wedge \hat{E}_0)$$

which is identical for the two mirror configurations of figure 5. In the low field limit, the $\beta E_0 M_1$ interference has a dispersive line shape, so it vanishes when the chiral term is maximum (see Fig. 6). We shall see that the above properties remain approximately the same in the more complicated case of non-zero nuclear spin with large magnetic field. Moreover, if we choose a configuration where the photon momentum \mathbf{k} is taken along the electric field, the $E_1^{\text{ind.}} M_1$ term vanishes while the $E_1^{\text{ind.}} E_1^{\text{p.v.}}$ stays unchanged.

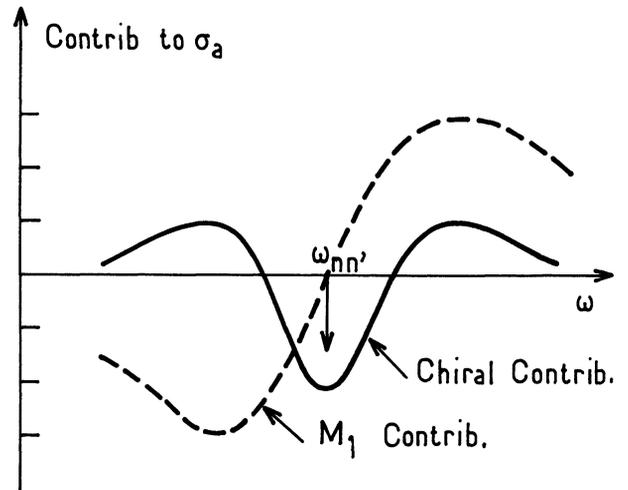


Fig. 6. — Frequency dependences for the chiral ($\beta E_0 E_1^{\text{p.v.}}$) contribution and the parity conserving contribution ($\beta E_0 M_1$) to the absorption cross-section (no nuclear spin).

3.2 RESULTS OF A COMPLETE CALCULATION FOR THE $6S_{1/2} \rightarrow 7S_{1/2}$ CESIUM TRANSITION. — The chiral effect being quadratic in the magnetic field H , rather high fields (of the order of one kilogauss) are necessary if one wishes to reach values of the parity violating

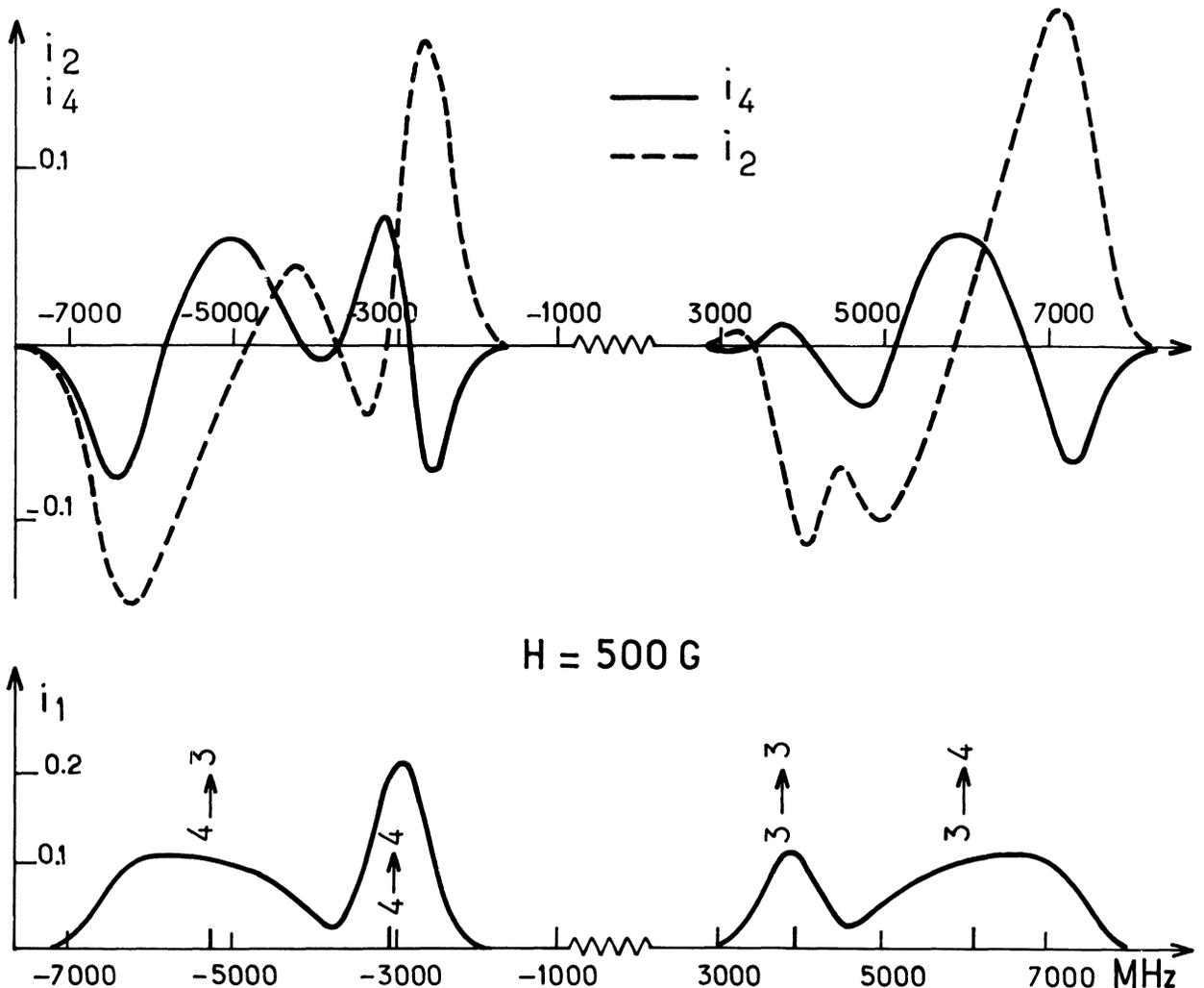


FIG. 7.

contribution to the cross-section of the same order as those obtained in a zero-field experiment. A perturbative treatment is no longer adequate. A numerical computation of the absorption cross-section for the forbidden 6S_{1/2} → 7S_{1/2} transition of Cesium has been performed (on UNIVAC 1110) as a function of the photon frequency, with different values of the magnetic fields 500 G and 1 000 G, for which the Zeeman splitting (0.35 MHz per Gauss) is larger than the Doppler half-height linewidth (750 MHz at 420 K).

Neglecting terms of the order of (M₁/βE₀)², the absorption cross-section can be written as :

$$\sigma_a(H, \omega) = \bar{\sigma}_a \left[i_1(\omega, H) - i_2(\omega, H) \frac{\sqrt{2} M_1}{\beta E_0} \times \right. \\ \left. \times (\boldsymbol{\varepsilon} \cdot \hat{H}) (\boldsymbol{\varepsilon} \cdot \hat{k} \wedge \hat{E}_0) + i_4(\omega, H) \frac{2 \operatorname{Im} E_1^{P.v.}}{\beta E_0} \times \right. \\ \left. \times (\boldsymbol{\varepsilon} \cdot \mathbf{H}) (\boldsymbol{\varepsilon} \cdot \hat{H} \wedge \hat{E}_0) \right] \quad (21)$$

$\bar{\sigma}_a$ being the resonance cross-section in zero magnetic field summed over the four hyperfine components.

The above expression is valid only for the special configuration described previously

$$(\boldsymbol{\varepsilon} \cdot \mathbf{E}_0 = \mathbf{E}_0 \cdot \mathbf{H} = 0 ; |\boldsymbol{\varepsilon} \cdot \hat{H}| = 1/\sqrt{2}).$$

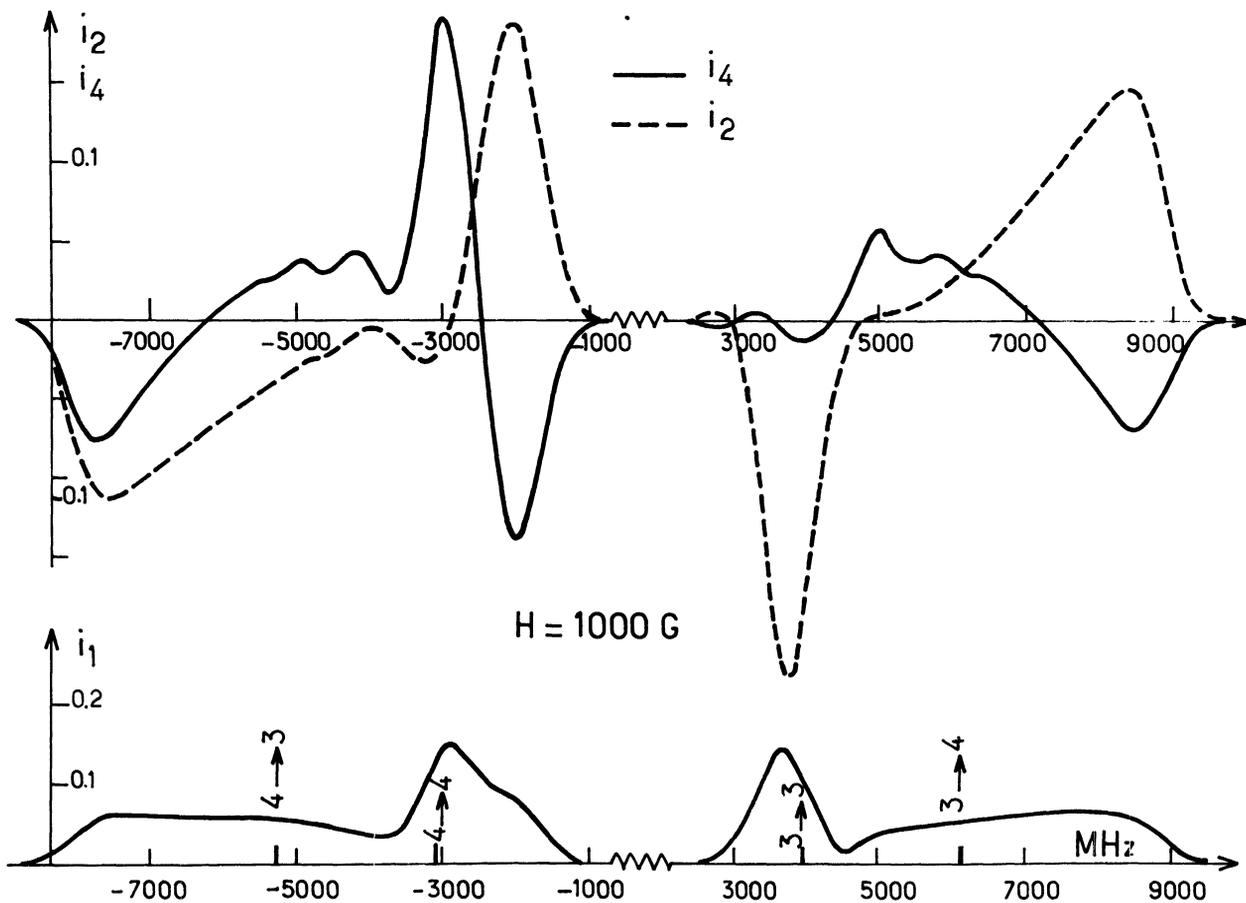
The numerical values of i_1 , i_2 and i_4 are plotted *versus* the laser frequency ω for the two values $H = 500$ G and $H = 1\,000$ G on figures 7 and 8. We note that there are frequencies ω for which both i_1 and i_4 reach values of the order of 0.1. The right

$$(\boldsymbol{\varepsilon} \cdot \hat{H}) (\boldsymbol{\varepsilon} \cdot \hat{H} \wedge \hat{E}_0) > 0,$$

left $(\boldsymbol{\varepsilon} \cdot \hat{H}) (\boldsymbol{\varepsilon} \cdot \hat{H} \wedge \hat{E}_0) < 0$, asymmetry characteristic of parity violation is :

$$\mathcal{A} = \frac{\sigma_a(R) - \sigma_a(L)}{\sigma_a(R) + \sigma_a(L)} = \frac{i_4}{i_1} \frac{\operatorname{Im} E_1^{P.v.}}{\beta E_0}. \quad (22)$$

\mathcal{A} is of the order of $\operatorname{Im} E_1^{P.v.}/\beta E_0$ as in experiments performed on $|\Delta F| = 1$ transitions in zero magnetic field.



Figs. 7 and 8. — Absorption cross-section for the Cesium 6S_{1/2}-7S_{1/2} transition *versus* laser frequency, in crossed electric and magnetic fields with a plane polarized beam. Fig. 7 is relative to $H = 500$ G and fig. 8 to $H = 1\,000$ G

$$\sigma_a = \bar{\sigma}_a \left[i_1 - i_2 \frac{\sqrt{2} M_1}{\beta E_0} (\boldsymbol{\varepsilon} \cdot \hat{H}) (\boldsymbol{\varepsilon} \cdot \hat{k} \wedge \hat{E}_0) - i_4 \frac{2 \operatorname{Im} E_1^{P.v.}}{\beta E_0} (\boldsymbol{\varepsilon} \cdot \hat{H}) (\boldsymbol{\varepsilon} \cdot \hat{H} \wedge \hat{E}_0) \right].$$

Lower curve represents the main contribution i_1 and lower curves give the interference terms : in broken line, the parity conserving term, i_2 , and in continuous line, the parity violating chiral term, i_4 .

The frequency dependence of $i_2(\omega)$ and $i_4(\omega)$ is rather complex but qualitative features of the nuclear spin zero case are still present. In particular the function $|i_4(\omega)|$ stays near its maximum at frequencies where the $E_1^{\text{ind.}}-M_1$ contribution vanishes ($i_2(\omega)=0$). For example, the frequencies $-3\,200$ MHz for $H = 500$ G and $-2\,900$ MHz for $H = 1\,000$ G look particularly interesting. More details about the calculation and discussion of the results can be found in reference [8].

The signal associated with the chiral absorption has many characteristic features which will allow discrimination against eventual spurious effects :

i) It is amplified by multipasses of the beam like the zero-field effect (search for $\mathbf{P}_e^{\text{p.v.}}$) and like the circular dichroism in crossed electric and magnetic dc fields. It should be noted that the parity conserving term $E_1^{\text{ind.}} M_1$ involving the scalar $(\boldsymbol{\varepsilon} \cdot \hat{H}) (\boldsymbol{\varepsilon} \cdot \hat{k} \wedge \hat{E}_0)$ is on the contrary suppressed by multipasses of the beam.

ii) The change of sign which occurs when either a reversal of \mathbf{E}_0 or a reversal of $\boldsymbol{\varepsilon} \cdot \mathbf{H}$ is performed constitutes quite a good signature of the signal. The parity conserving term $\beta E_0 M_1 (\boldsymbol{\varepsilon} \cdot \hat{H}) (\boldsymbol{\varepsilon} \cdot \hat{k} \wedge \hat{E}_0)$, also linear in \mathbf{E}_0 and in $\boldsymbol{\varepsilon} \cdot \hat{H}$, can be suppressed by one or several of the following means :

- multipasses of the laser beam,
- H reversals,
- judicious choice of the laser frequency,
- judicious choice of geometric configuration ($\mathbf{k} \not\parallel \mathbf{E}_0$).

iii) The peculiar variation *versus* the laser frequency will be also very helpful to identify the effect.

With respect to the experiment currently underway which searches for $\mathbf{P}_e^{\text{p.v.}}$ in zero-field, the chiral absorption experiment may present several advantages :

i) Removal of the circular analyzer on the fluorescence beam allows one to increase the photon flux collected by the detector and consequently the signal to noise ratio.

ii) Use of linearly polarized incident light simplifies the technological problem of the high efficiency multipass cell. It is no longer necessary, as in the case of circular polarization, to put the mirrors inside the alkali vapor cell since Brewster windows are now allowed.

Conclusion. — We have analyzed possible manifestations of parity violation in forbidden radiative M_1 transitions, in two new chiral configurations involving crossed transverse dc electric and magnetic fields. As for the zero magnetic field experiment presently under progress [4, 5], they involve an interference effect taking place between the Stark induced and the parity violating electric dipole amplitudes of transitions. The asymmetry under search is of the

same order of magnitude but the main point is that the experimental apparatus may be simplified (for instance multipasses of the beam in a Brewster cell) and the efficiency of collection may be improved (no circular analyzer). At the same time, the signature of parity violating effects remains well defined owing to their special behaviour under normal incident light reflexion, under reversals of the applied fields and incident light polarization and under variation of incident light frequency. The numerous and varied tests which can be achieved make these experiments relatively safe from systematic errors and certainly quite complementary to the present one.

Appendix I. — In this appendix, we would like to construct, from the physical quantities which specify the experimental set-up (external electric field \mathbf{E}_0 , photon momentum \mathbf{k} , etc.), the possible rotational invariants odd under space reflexion and even under time reflexion (hereafter called $P_- T_+$ invariants). The appearance of a $P_- T_+$ invariant in the transition probability is a signature for a parity violating time reflexion invariant interaction in the atomic Hamiltonian. We assume, here, that the initial and final states can be considered as stationary states with respect to the total Hamiltonian of the system. We restrict our analysis to experiments involving an unpolarized initial atomic state and no polarization measurement will be performed on the final state.

Let us first consider the case when only a static electric field \mathbf{E}_0 is present ($\mathbf{H} = 0$) with a linearly polarized incident light beam of polarization $\boldsymbol{\varepsilon}$. If we neglect the effects associated with the M_1 amplitude, the momentum of the photon should not appear explicitly. The experiment is then completely specified by the two vectors \mathbf{E}_0 and $\boldsymbol{\varepsilon}$. It is clearly impossible to build a $P_- T_+$ invariant from these two vectors. With a circularly polarized beam of momentum \mathbf{k}_i and circular polarization ζ_i , the only $P_- T_+$ invariant is ζ_i itself; but we are interested in an interference between $E_1^{\text{p.v.}}$ and $E_1^{\text{ind.}}$ so that the $P_- T_+$ invariant should be linear in the electric field. In conclusion : there is no interference term between the Stark induced electric amplitude $E_1^{\text{ind.}}$ and the parity violating amplitude $E_1^{\text{p.v.}}$ in the absorption cross-section on unpolarized atoms if no external magnetic field is present ($\mathbf{H} = 0$).

Let us turn now to the situation of interest, where \mathbf{H} is non-zero, and look, first, at the case of a plane polarized light beam. We have three vectors at our disposal, $\boldsymbol{\varepsilon}$, \mathbf{H} , \mathbf{E}_0 , from which we can build the four basic rotational invariants :

$$\boldsymbol{\varepsilon} \cdot \mathbf{H}, \quad \boldsymbol{\varepsilon} \cdot \mathbf{E}_0, \quad \mathbf{H} \cdot \mathbf{E}_0, \quad \boldsymbol{\varepsilon} \cdot \mathbf{H} \wedge \mathbf{E}_0.$$

Among all the monomials obtained from these basic invariants, we have to pick those satisfying the following requirements : they have to be :

- i) linear in \mathbf{E}_0 ($E_1^{\text{p.v.}} E_1^{\text{ind.}}$ interference term),

ii) quadratic in $\boldsymbol{\varepsilon}$ (each transition amplitude is clearly linear in $\boldsymbol{\varepsilon}$),

iii) odd under : $\mathbf{E}_0 \rightarrow -\mathbf{E}_0$, $\boldsymbol{\varepsilon} \rightarrow -\boldsymbol{\varepsilon}$, $\mathbf{H} \rightarrow \mathbf{H}$,

iv) even under : $\mathbf{E}_0 \rightarrow \mathbf{E}_0$, $\boldsymbol{\varepsilon} \rightarrow \boldsymbol{\varepsilon}$, $\mathbf{H} \rightarrow -\mathbf{H}$.

The conditions iii) and iv) are those for a $P_- T_+$ invariant. A simple enumeration leads to

$$(\boldsymbol{\varepsilon} \cdot \mathbf{H}) (\boldsymbol{\varepsilon} \cdot \mathbf{H} \wedge \mathbf{E}_0)$$

as the only possible invariant.

The case of a circularly polarized incident light beam is treated in a similar way. The experiment is described by the three vectors : \mathbf{k} , \mathbf{H} , \mathbf{E}_0 and the pseudoscalar $\xi = \pm 1$. The basic rotational invariants are :

$$\mathbf{k} \cdot \mathbf{H}, \quad \mathbf{k} \cdot \mathbf{E}_0, \quad \mathbf{H} \cdot \mathbf{E}_0, \quad \mathbf{k} \cdot \mathbf{H} \wedge \mathbf{E}_0, \quad \xi.$$

We build from them the monomials satisfying the three conditions :

i) linear in \mathbf{E}_0 ,

ii) odd under $\mathbf{k} \rightarrow -\mathbf{k}$, $\mathbf{H} \rightarrow \mathbf{H}$, $\mathbf{E}_0 \rightarrow -\mathbf{E}_0$, $\xi \rightarrow -\xi$,

iii) even under $\mathbf{k} \rightarrow -\mathbf{k}$, $\mathbf{H} \rightarrow -\mathbf{H}$, $\mathbf{E}_0 \rightarrow \mathbf{E}_0$, $\xi \rightarrow \xi$.

The only $P_- T_+$ invariant linear in \mathbf{E}_0 is readily obtained by inspection : $\xi \mathbf{k} \cdot (\mathbf{E}_0 \wedge \mathbf{H})$.

The two invariants

$$(\boldsymbol{\varepsilon} \cdot \mathbf{H}) (\boldsymbol{\varepsilon} \cdot \mathbf{H} \wedge \mathbf{E}_0) \quad \text{and} \quad \xi (\mathbf{k} \cdot \mathbf{E}_0 \wedge \mathbf{H})$$

can be multiplied by an arbitrary function of \mathbf{H}^2 .

Appendix II. — Using symmetry considerations, we want to show, here, that the two parity violating interference effects differ in the fact, that the chiral absorption involving a plane polarized wave vanishes in the limit of L-S coupling, while the circular dichroism stays finite in the same limit.

This result applies only to an $nS \rightarrow n'S$ transition, where the effect of an external magnetic field \mathbf{H} can be accounted for by a Zeeman Hamiltonian involving only the electronic spin \mathbf{S} :

$$H_Z = g\mu_B H_0 S_z. \quad (\text{A.1})$$

Note that we do not need to suppose that the S states are necessarily single electron states.

The effect of the magnetic field \mathbf{H}_0 on the virtual S-P transitions involved in the calculation of $E_1^{P.v.}$ are negligible.

In the limit of L-S coupling with only electric dipole interaction with the radiation field, the parity conserving part of the total Hamiltonian $H^{P.c.}$ depends on the spin variables of the electron only through the Zeeman Hamiltonian H_Z and the hyperfine Hamiltonian :

$$H_{h.f.} = \sum_i A_i \mathbf{s}_i \cdot \mathbf{I} \quad (\text{A.2})$$

\mathbf{I} being the nuclear spin.

Thus $H^{P.c.}$ is invariant under the unitary transformations $U_z(\varphi)$

$$U_z(\varphi) = e^{-i(S_z + I_z)\varphi},$$

associated with arbitrary rotations around the z -axis in spin S space. Since the two pseudoscalars to be observed (neither $\xi_i \mathbf{k}_i \cdot \mathbf{E}_0 \wedge \mathbf{H}$, nor $(\boldsymbol{\varepsilon} \cdot \mathbf{H}) (\boldsymbol{\varepsilon} \cdot \mathbf{H} \wedge \mathbf{E}_0)$) do not involve quantities associated with the spin degrees of freedom, only the part of the parity violating Hamiltonian $H^{P.v.}$, also invariant under the transformation $U_z(\varphi)$, will contribute.

The parity violating Hamiltonian induced by the neutral currents, relevant for heavy atoms [2], has the following form :

$$H^{P.v.} = \sum_{i=1}^Z \kappa (\mathbf{s}_i \cdot \mathbf{p}_i \delta^3(\mathbf{r}_i) + \delta^3(\mathbf{r}_i) \mathbf{s}_i \cdot \mathbf{p}_i) \quad (\text{A.3})$$

\mathbf{p}_i , \mathbf{s}_i , \mathbf{r}_i are respectively the momentum, spin, position operator of the i -electron and κ is a nuclear constant given by :

$$\kappa = \frac{G_F}{2\sqrt{2}m_e} Q_W K_r,$$

where Q_W is the weak charge of the nucleus, and K_r a relativistic correction [2].

Let us write the part of $H^{P.v.}$ which is invariant under $U_z(\varphi)$

$$H_z^{P.v.} = \kappa \sum_i (s_{iz} p_{iz} \delta^3(\mathbf{r}_i) + \delta^3(\mathbf{r}_i) s_{iz} p_{iz}). \quad (\text{A.4})$$

Let us, now, look at the symmetry property of the two Hamiltonian $H^{P.c.}$ and $H_z^{P.v.}$ under another transformation Θ , acting only on the atomic Hamiltonian $(^2)$, and such that :

$$|\psi'\rangle \rightarrow |\psi\rangle = U_y(\pi) |\psi\rangle \quad (\text{A.5})$$

and

$$\mathbf{H} \rightarrow -\mathbf{H},$$

where U_y is the unitary transformation associated with a rotation of π around the y -axis in the spin space :

$$U_y(\pi) = e^{-i(S_y + I_y)\pi}. \quad (\text{A.6})$$

Still in the limit of L-S coupling, the parity conserving part of the total Hamiltonian $H^{P.c.}$ is invariant under the transformation Θ , while the *diagonal* part of the parity violating Hamiltonian, $H_z^{P.v.}$, is odd.

(²) With a suitable representation of the radiation field states, it can be proved that the product of Θ by the time reflexion operation T is nothing but the complex conjugation operator K .

We remark that the pseudoscalar $(\boldsymbol{\epsilon} \cdot \mathbf{H})(\boldsymbol{\epsilon} \cdot \mathbf{H} \wedge \mathbf{E}_0)$ is invariant under Θ . Since it results from an interference effect between two amplitudes of transition induced one by $H^{p.c.}$, even under Θ , and the other by $H_z^{p.v.}$, odd under Θ , it necessarily vanishes in the limit of L-S coupling. On the other hand, the pseudoscalar $\xi k \cdot E_0 \wedge H$, which is odd under Θ , stays finite in the same limit.

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