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Investigation of the dynamic Stark effect of a $J = 0 \rightarrow J = 1$ optical transition

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Résumé. — Nous avons étudié l'effet Stark dynamique (dédoublement Autler-Townes et déplacement radiatif) d'une transition optique $J = 0 \rightarrow J = 1$. Le jet atomique de barium est éclairé par un faisceau laser intense (« pompe »), résonant ou légèrement non résonnant pour la raie de résonance de BaI ($\lambda = 5 \text{ 535 Å}$); la fluorescence produite par un deuxième faisceau laser, peu intense (« sonde »), est détectée en fonction de la fréquence sonde. Pour des polarisations linéaires des deux faisceaux on a affaire soit à un système à deux niveaux (polarisations parallèles), soit à un système à 3 niveaux (polarisations perpendiculaires). Dans ce dernier cas, pour de faibles valeurs du désaccord de fréquence du faisceau intense, nous avons observé le doublet Autler-Townes. Pour les grands désaccords de fréquence du faisceau pompe, nous observons une seule résonance déplacée, aussi bien pour des polarisations parallèles que pour des polarisations perpendiculaires. Dans le cas du système à trois niveaux, nous montrons le passage continu de l'effet Autler-Townes au déplacement radiatif. Nous présentons les calculs correspondant aux cas étudiés expérimentalement; nous nous intéressons particulièrement au cas des grands désaccords de fréquence pour le faisceau intense (déplacement radiatif). L'évolution calculée pour les courbes de résonance (positions et intensités des pics) quand on fait varier les directions de polarisation relatives est en bon accord avec les observations expérimentales.

Abstract. — The dynamic Stark effect (Autler-Townes splitting and radiative shift) of a $J = 0 \rightarrow J = 1$ optical transition have been investigated. The barium atomic beam is illuminated by a strong pump beam, resonant or slightly non-resonant with the BaI resonance line ($\lambda = 5 \text{ 535 Å}$); the fluorescence induced by a weak probe beam is detected versus the probe frequency. With linear polarizations for the two beams one deals either with a two-level (parallel polarizations) or with a three-level system (perpendicular polarizations). In the latter case and for small detunings of the pump beam, the Autler-Townes doublet has been observed. For large detunings of the pump beam a single, shifted resonance has been obtained with both polarization settings. In the three-level case the observation of the passage from the Autler-Townes situation to the light-shift situation, is reported. The corresponding calculations are presented; a particular attention is paid to the case of large detunings for the pump beam (light-shift situation). The evolution calculated for the resonance curve (position and relative intensities of the peaks) when the relative polarization directions are changed is consistent with the experimental data.

1. Introduction. — It has been known for a long time that atomic absorption lines could be split or shifted under the action of a strong resonant or slightly non-resonant electromagnetic field [1-3]. These effects, known as the a.c. Stark effect, have been intensively studied recently in the optical range, using laser beams as perturbing fields. The optical Autler-Townes splitting has been observed for atoms at rest (atomic beam experiment) [4] and for Doppler-broadened systems [5]. The spectral analysis of resonance fluorescence under strong monochromatic excitation has led to the observation of the expected triplet structure [6]. The optical light-shifts have also received some attention from experimentalists who have either observed the shift itself or used its properties for applications [7, 8].

Let us consider first a typical situation for the observation of the a.c. Stark effect: an atomic transition...
c ⇔ b is perturbed by a strong, resonant and monochromatic field; the absorption spectrum of a weak probe beam is measured either in the frequency range of the c ⇔ b transition (2-level system) or in the frequency range of a transition c ⇔ a sharing a common level. For the sake of simplicity, we suppose that, without any applied fields, the total atomic population is in the common level c. If the perturbing field is strong enough (i.e. if the Rabi nutation frequency is larger than the relaxation rates), the atomic levels are split; this splitting is responsible for the observation of the Autler-Townes doublet and can be physically understood using the dressed atom picture [9]. When the detuning of the strong field increases, the doublet is no longer symmetrical [10]: one component increases and goes towards the position of the unperturbed line and the other one decreases and goes away from this position. For large values of the detuning, the second component vanishes and one is left with a single resonance, slightly shifted from resonance. For an infinite detuning of the pump beam, the light-shift goes to zero and the unperturbed resonance is observed. This continuity and this relation between the Autler-Townes effect and the radiative shift has not, to our knowledge, been demonstrated experimentally in the optical range.

We have paid a particular attention to this point in our experimental investigation of the dynamic Stark effect performed on a barium atomic beam. We have studied the BaI resonance transition

$$6s^2 \, ^1S_0 \rightarrow 6s \, 6p \, ^1P_1 .$$

With linear polarizations of the two beams one can perform either a three-level experiment (perpendicular polarizations, levels : b, c, b±) or a two-level experiment (parallel polarizations, levels : b, c) (Fig. 1). The choice of the probe polarization thus makes possible, for large detunings of the pump beam, the observation of a transition with only one shifted level or with the two levels shifted. Fluorescence detection has been used in the experiment rather than absorption detection [11] and this leads to some differences as shown in the theoretical investigation. In the absorption case, the density matrix element of interest (optical coherence) appears in the first order of the development in the probe electric field powers, while in the fluorescence case, the corresponding elements (populations and Hertzian coherences) appears in the second order. This leads to some additional, dispersion shaped, resonance terms discussed in the theoretical section of the paper; these terms arise when the angle $\alpha$ of the polarization directions of the two beams differs from 0 or $\pi/2$.

2. The a.c. Stark effect : theoretical investigation in a $J = 0 \leftrightarrow J = 1$ transition. — Let us consider an atomic beam illuminated at right angles by two laser beams with frequencies $\omega_1$ and $\omega_2$ very close to the frequency $\omega_{bc}$ of the atomic transition $c(J = 0) \leftrightarrow b(J = 1)$.

The two beams are propagating along Oy, and Oz is the direction of the atomic beam (Fig. 2). The weak probe beam (amplitude $E_1$) is linearly polarized along Oz, the strong pump beam (amplitude $E_2$) is also linearly polarized and $\alpha$ designates the angle between the polarization directions in the xOz plane. We consider the fluorescence light emitted in this xOz plane in a direction making an angle $\theta$ with Ox; this is just the fluorescence light $L_{f}(e_\theta)$ emitted with a polarization $e_\theta$ making an angle $\theta$ with the strong field polarization direction $E_2$ (quantization axis). For convenience, we had $\theta = \alpha$ for the experiment.

The fluorescence intensity $L_f(e_\theta)$ is proportional to $Tr_b \rho(e_\theta, d) (e_\theta, d)$, $d$ is the electric-dipole operator, $\rho$ is the density matrix of the atomic system and $b$ is the upper level of the transition. One gets:

$$L_f(e_\theta) \propto \rho_{b b_0} \cos^2 \theta + \frac{1}{2} \rho_{b \pm b_\pm} + \rho_{b_+ b_-} - 2 \text{Re} \rho_{b_+ b_-} \sin^2 \theta - \sqrt{2} (\text{Re} \rho_{b_+ b_0} - \rho_{b_+ b_0} \cos \theta \sin \theta) .$$

The elements of the density matrix depends upon $\alpha$ in the following way:

$$\rho_{b b_0} = A \cos^2 \alpha + B \sin^2 \alpha$$

$$\frac{1}{2} (\rho_{b_+ b_+} + \rho_{b_- b_-} - 2 \text{Re} \rho_{b_+ b_-}) = C \sin^2 \alpha$$

$$\sqrt{2} (\text{Re} \rho_{b_+ b_0} - \rho_{b_+ b_0} \cos \alpha) = 2 \text{Re} D \sin \alpha \cos \alpha .$$

Fig. 1. — Level scheme and laser beam polarizations for the observation of the a.c. Stark effect.

Fig. 2. — Geometrical arrangement of the atomic beam and of the laser beams. The quantization axis is along $E_2$; $\alpha$ is the angle between the polarization directions $E_1$ and $E_2$ of the beams; we had $\theta = \alpha$ for the experiment.
The analytical expressions of $A$, $B$, $C$ and $D$ are given in the Appendix. The calculations have been made according to the method developed by Ben Reuven and Klein [12, 13] i.e. using the linear response Green's functions $G_{ij}(n_1, n_2)$ to solve the non linear problem of concerns. The interaction with the weak field is treated to the lowest order (i.e. second order in the particular case) while the interaction with the saturating field is treated exactly to all orders by using diagrams characterized by the numbers $n_1$ and $n_2$ (differences of the number of photons for the fields $E_1$ and $E_2$ for the elementary processes). The Green's functions are:

$$G_{ij}(n_1, n_2) = [n_1 \omega_1 + n_2 \omega_2 - \omega_{ij} + i\Gamma_{ij}]^{-1};$$

$\omega_{1,2}$ and $\omega_{ij}$ are respectively the laser frequencies and the atomic frequencies; $i,j = b_0, b_+ , c$; $n_1, n_2 = 0, \pm 1, \pm 2$.

These are the occurring values for the harmonic numbers [12, 13] when the non-resonant Green's functions are neglected (this is equivalent to the well-known rotating wave approximation).

Let us discuss now some particular cases for the relative linear polarizations of the two beams; $\alpha$ is the angle between the polarization vectors and we assume that the fluorescence light is detected along Ox ($\theta = \alpha$) (Fig. 2). We only consider the part of the fluorescence light $L_F$ depending upon the probe field intensity; the corresponding signal can, in principle, be obtained experimentally using an appropriate modulation technique.

2.1 $\alpha = \pi/2$; $L_F(\epsilon_1) = C$. — The resonance conditions for the system are given by the zeroes of the denominator of $C$, i.e. of the quantity:

$$\det Y_b = \{1 - \beta_2^2 G_{bb}(1, -1) - G_{bc}(1, 0)\}$$

$$G_{bb}(1, -1) = [\delta_1 - \delta_2 + i\Gamma_b]^{-1};$$

$$G_{bc}(1, 0) = [\delta_1 + i\Gamma_b/2]^{-1};$$

$\delta_1 = \omega_1 - \omega_{bc}$ and $\delta_2 = \omega_2 - \omega_{bc}$ are the frequency detunings; $2\beta_2 = d_{bc} E_2/\hbar$ is the Rabi nutation frequency for the strong beam; $\Gamma_b$ is the radiative relaxation rate of level $b$. The equation $\det Y_b = 0$ leads, in the case $\Gamma_b \ll \beta_2$, to the resonance condition:

$$\delta_1 = \frac{1}{2} \{ \delta_2 \pm \sqrt{\delta_2^2 + 4 \beta_2^2} \}. \quad (3)$$

For $\delta_2 = 0$, one has the Autler-Townes doublet

$\delta_1 = \pm \beta_2$ and for large values of $\delta_2 (\delta_2 \gg \beta_2)$ one gets the shifted one-photon resonance:

$\delta_1 \simeq -\beta_2^2/\delta_2 = \Delta$ (radiative shift) and the shifted two-photon resonance : $\delta_1 = \delta_2 - \Delta$. Under the conditions $\delta_2 \gg \beta_2 \gg \Gamma_b$, the simplified expression for the signal of interest is just:

$$C \simeq \beta_1^2 \rho_c^0 \{[(\delta_1 - \Delta)^2 + \Gamma_b^2/4] \}. \quad (4)$$

The resonance curves obtained with the exact expression of $C$ are drawn on figures 3 and 4. For zero detuning of the strong field ($\delta_2 = 0$), one gets the well-known symmetrical Autler-Townes doublet as in the three-level system case [1, 4, 10]. The strong field splits both the $b_0$ and $c$ levels, but the weak field only probes the $c$ level (Fig. 1); the corresponding spectra shows two symmetrical peaks separated by the Rabi frequency $2\beta_2$ (Fig. 3). When $\delta_2$ increases, one of the peaks increases and approaches the unperturbed resonance position while the other one decreases drastically and goes away from this position (Fig. 3). For large values of the detuning, one gets a single
resonance which is found to be very close to the expected position $\delta_1 = - \beta_2^2 / \Omega_2 = d$ (Fig. 4).

2.2 $\alpha = 0$ ; $L_4(e_{i}) = A.$ — The resonant denominator is now:

$$\det Y = \left\{ 1 - \beta_2^2 (G_{cd}(1, -1) + G_{ce}(1, -1)) - g_{cd}(1, -1)[G_{bc}(1, 0) + G_{cd}(1, -2)] \right\}.$$ 

The resonance condition $\det Y = 0$ can be written in the strong field limit ($\beta_2 \gg \Gamma_b$):

$$- \delta_1^2 + 3 \delta_2^2 + \delta_3 (4 \beta_2^2 - 2 \delta_3^2) - 4 \delta_2 \beta_2^2 = 0.$$ 

For $\delta_2 = 0$ one has three resonances : $\delta_1 = 0$ and $\delta_3 = \pm 2 \beta_2$; for large values of $\beta_2$ ($\beta_2 \gg \beta_3$), one gets $\delta_3 \approx -2 \beta_2^2 / (\delta_2) = 2 \Delta$. In the latter case the approximation consists in neglecting the term $G_{cb}(1, 2)$ (three-photon process).

Under the conditions $\beta_2 \gg \beta_3 \gg \Gamma_b$ the resonance shape is actually given by the simplified expression:

$$A \approx \beta_2^2 \rho_c^2 [(\delta_1 - 2 \Delta)^2 + \Gamma_b^2 / 4].$$

Using the exact expression for $A$, one gets resonance curves which are similar for $\delta_2 = 0$ to that previously shown by Mollow [14] and by Wu et al. [11], taking into account that in our particular case we consider fluorescence rather than absorption for detection. The corresponding curves are not given in the present paper since it has not been possible to obtain the corresponding experimental curves; they can be found in part [15]. As the experiment was made for a non-resonant strong beam, we have calculated the corresponding curve (Fig. 4) which shows a single shifted resonance; the shift is nearly equal to $2 \Delta$. This corresponds to the fact that the probe beam is sensitive to the combined shifts of levels $b_0$ and $c$ (Fig. 1) resulting in a net modification of the resonance position of $2 \Delta$.

2.3 General case ($\alpha \neq 0, \pi/2$). — Except for the two particular cases discussed above, the general expression for the fluorescence signal is not a simple linear combination of $A$ and $C$ but also depends on $B$ and $D$.

For instance:

$$L_4(e_{i, 4}) = (A + C + B - 4 \text{Re } D)/4$$

$$L_4(e_{i, 0}) = (9 A + C + 3 B - 12 \text{Re } D)/16.$$ 

The corresponding calculation for the absorption of the weak field under the same geometry gives a linear combination of $\chi_{ii} (\alpha = 0)$ and $\chi_{ii} (\alpha = \pi/2)$ for the susceptibility $\chi_a$. In the fluorescence calculation, we had to include second order terms relative to the weak field and this gives rise to the $B$ and $D$ contributions.

The simplified expression for $D$, when $\delta_2 \gg \beta_2$, is the following:

$$D \approx \beta_2^2 \rho_c^2 \frac{1}{\Delta + i \Gamma_b} \left( \frac{1}{(\delta_1 - \Delta) + i \Gamma_b/2} + \frac{1}{(\delta_1 - 2 \Delta) - i \Gamma_b/2} \right).$$

The real part of $D$ appears as a difference between two dispersion shaped functions centred respectively at $\delta_1 = \Delta$ and at $\delta_1 = 2 \Delta$. This results in dissymmetrical peaks and negative parts in the spectra for $\alpha \neq 0, \pi/2$ (Fig. 4).

The term $B$ is centred at $\delta_1 = \Delta$ but is smaller than $A$ and $C$ when $\delta_2$ is large. This small term could be observed for intermediate values of $\delta_2$ under the geometrical conditions $\theta = 0$ and $\alpha = \pi/2$. The expression for $B$ is a difference of two terms which produces a nearly total compensation, each of them being non zero.

The calculated spectra corresponding to $\alpha = \pi/4$ and $\alpha = \pi/6$ are shown on figure 4; they exhibit two peaks approximately centred at $\Delta$ and $2 \Delta$.

2.4 Circularly polarized saturating beam. — As the corresponding spectrum has been obtained experimentally, we have also considered the case of a circularly polarized strong beam. The weak field is linearly polarized along Ox and the fluorescence intensity $L_4(e_{i})$ emitted along Oy and polarized along Ox is given by:

$$L_4(e_{i}) = \left( A + C + B - 2 \text{Re } D \right).$$

The corresponding spectrum, under light-shift conditions ($\gamma_2 \gg \gamma_1$) is similar to that obtained for $L_4(e_{i, 4})$ (Fig. 4), with, as expected, the same relative amplitude of the two peaks.

3. Experimental investigation of the a.c. Stark effect in the Bal resonance line ($6s^2 1S_0 \rightarrow 6s 6p 1P_1$).

3.1 Experimental arrangement (Fig. 5). — The atomic beam apparatus and the dye lasers are the same as in the previous experiment on the optical Hanle effect (light-shift induced zero-field level crossing) [8, 16].

The barium atomic beam is crossed at right angle by two counter propagating light beams produced by c.w. single mode dye lasers operating at $\lambda = 535$ Å. The strong beam is linearly polarized and produces the resonant or the off-resonant a.c. Stark effect. The weak beam is linearly polarized along the beam

Fig. 5. — Experimental set-up for the observation of the a.c. Stark effect.
A.C. Stark Effect \(0 \leftrightarrow 1\) Optical Transition

Direction; its frequency is continuously tuned across the BaI line. The intensity of the probe beam is kept at a constant level by means of a servo-loop acting upon the Ar\(^+\) laser current. The absorption spectrum of the weak beam is observed, using the fluorescence light emitted at right angle from the atomic beam and from the laser beams for detection. The probe beam is mechanically chopped and a lock-in amplifier is used to discriminate the part of the fluorescence signal due to this beam. Rectangular diaphragms are used to select in the beam image the fluorescence light emitted in the overlap region of the two beams. A part of the probe beam is sent to an auxiliary interaction region of the atomic beam to get an unperturbed reference spectra; as the isotopic structure of the BaI resonance line is well known [17] this also provides a frequency calibration.

Let us point out now two limiting characteristics of the experiment. First, we had to work with natural barium which includes undesirable isotopes. However, under our polarization conditions, the \(^{138}\text{Ba}\) resonance line is much more intense than the 8 other components which are all located on the high frequency side [17]. Negative values of the detuning \(\Delta_2\) have been used for the experiment and the presence of the other isotopes does not prevent the observation of the expected phenomena.

Another point is that we use Gaussian beams. Even with a probe beam smaller than the pump beam, one has to deal with a non-uniform power density at the interaction region. The effective Rabi frequency values are spread over a certain range and this results mainly in a broadening of the shifted resonances (Fig. 6).

3.2 Experimental results. — 3.2.1 Autler-Townes doublet and radiative shift. — For small values of \(\Delta_2\) (resonant strong beam) the experiment could only be performed for crossed polarizations of the two fields (\(\alpha = 90^\circ\), three-level system). For parallel polarizations of the beams (\(\alpha = 0^\circ\), two-level system), the fluorescence light produced by the strong beam is very large (\(J = 0 \rightarrow J = 1\) transition, geometry of the experiment on figure 5) and the useful signal i.e. the fluorescence light due to the weak beam is not detectable with an acceptable signal-to-noise ratio. For perpendicular polarizations (\(\alpha = 90^\circ\)) no fluorescence is produced by the strong beam in the direction parallel to its polarization vector and the signal-to-noise ratio of interest is good. For large values of \(\Delta_2\) (off-resonant strong beam), no fluorescence from the strong beam occurs and the experiment has been possible even with \(\alpha = 0^\circ\).

Figure 6 shows, in the perpendicular polarization case, the evolution of the spectrum for increasing values of \(\Delta_2\). For \(\Delta_2 = 0\), one gets the well-known symmetrical Autler-Townes doublet with a separation \(2\beta_2/2\pi \approx 600\) MHz corresponding to a power density of 35 W/cm\(^2\) for the strong beam at the interaction region.

For increasing values of \(\Delta_2\), one gets a large peak coming closer to the unperturbed atomic line position (this position is obtained from the reference spectrum) and a small peak going away from this position. For large values of \(\Delta_2\), only one peak could be detected; this peak is just the light-shifted resonance. The observed evolution is in complete agreement with the theoretical prediction (Fig. 3). On figure 6, one also sees an additional and small peak at the unperturbed atomic line position. This unexpected signal is probably due to spurious fluorescence light coming from a part of the atomic beam not illuminated by the strong laser beam. It has been considerably reduced to get the spectra of figure 6 by adding very small slits just before the photomultiplier (Fig. 5).

Figure 6 clearly illustrates the relation between the Autler-Townes effect and the radiative shift effect; to our knowledge this is the first observation in the optical range of the continuous passage from one effect to the other.

This could not be obtained in the previous experimental investigations of the Autler-Townes effect in atomic beams [4] since the probe transition in the previous experiment occurs between levels normally not populated in the absence of the strong beam. The evolution of the doublet for non-zero detunings strongly differs in this case [4, 10] from that corresponding to our situation (three-level system with a common, populated, lower level for the two transitions).

We have measured the optical light-shift for different values of the detuning \(\Delta_2\) in order to check experimentally the \(1/\Delta_2\) law. The results are given on figure 7. The observed shifts appear to be proportional to \(1/\Delta_2\) except for the lowest values of the detuning \((\Delta_2/2\pi < 750\) MHz\). The slope of the straight line of figure 7 gives a coefficient of proportionality of \((270\) MHz\(^2\)) between the observed shift and \(1/\Delta_2\).
The value of 270 MHz can be compared to half the separation of the Autler-Townes doublet in the same experiment which was measured to be 290 MHz.

3.2.2 Polarization dependence of the optical light-shift. — The experimental resonance curves to be considered now have been obtained in the light-shift regime (large $\delta_2$ values). The evolution of the spectrum when the angle $\alpha$ between the polarization directions is changed, is illustrated on figure 7. The whole set of experimental curves have been obtained for the same value of the power and of the detuning for the strong beam i.e. for a specified value of $\delta_2 = -\beta_2^2/\delta_2$.

$L_F(e_\perp)$ shows, as in figure 6 for large $\delta_2$ values, a single peak shifted of $\Delta$ from the reference position; this corresponds to the observation of the shift of the c level (Fig. 1).

For parallel polarizations of the two beams ($\alpha = 0^\circ$) a single peak is also obtained and the corresponding shift is approximately twice as large as in the $\alpha = 90^\circ$ case. As discussed previously this corresponds to the observation of the combined shifts of the levels c and b$_0$ (Fig. 1). Apart from a rather poor resolution in the experimental spectra, the agreement with the theoretical predictions is good.

For intermediate values of $\alpha$, one gets more or less an unresolved superposition of the $\alpha = 0^\circ$ and of the $\alpha = 90^\circ$ spectra. The linewidth (mainly due to the power inhomogeneity of the strong beam) is too large to resolve the $\alpha = 45^\circ$ and $\alpha = 30^\circ$ spectra (Fig. 8). Thus, the predicted asymmetry of the peaks (due to the term $D$ of the theory) cannot be checked. Nevertheless, the relative heights of the two peaks and consequently the asymmetrical aspect of the experimental curves are fully compatible with the results of figure 4.

We have also studied the case of a circular polarization for the strong beam. The linearly polarized probe beam can be considered as a superposition of $\sigma_+$ and $\sigma_-$ circularly polarized beams. The saturating beam acts upon the c and b$_-$ levels; one circularly polarized component of the weak beam probes the shift of the c level only while the other component probes the combined shifts of the c and b$_-$ levels. The situation is very similar to that considered previously for a linear polarization of the strong beam and $\alpha = 45^\circ$. In the light-shift regime

$$(\delta_2/2\pi = -750 \text{ MHz})$$

Fig. 7. — Variation of the measured shift with the inverse of the detuning of the strong beam $1/\delta_2$.

![Graph showing variation of measured shift](image1)

Fig. 8. — Experimental observation of the dependence of the optical light-shift with the polarization direction of the strong beam.

![Experimental observation of dependence](image2)

Fig. 9. — Comparison between the experimental curves $I_F(e_{45})$ (linear polarization of the strong beam, $\alpha = 45^\circ$) and $I_F(e_\perp)$ (circular polarization of the strong beam).
we have indeed obtained very similar spectra in the case of circular polarization ($L_T(e_\perp)$) and in the case of linear polarization ($L_T(e_\parallel)$) as shown on figure 9.

4. Conclusion. — We have studied both experimentally and theoretically some aspects of the a.c. Stark effect, namely the passage from the Autler-Townes doublet to the light-shifted resonance on the one hand and the polarization dependence of the radiative shift on the other hand.

To make a complete test of the theoretical predictions in the case of fluorescence detection i.e. to observe the term $B$ or the asymmetrical peaks due to the dispersion-shaped $D$ term, it would be necessary to improve the experimental conditions. First, we need a pure $^{138}\text{Ba}$ atomic beam with a rather good collimation ratio. Second, the power density for the strong beam at the interaction region must be as constant as possible; this would require much more output power and much less focusing for the strong beam. Then, some attention must be paid to the total elimination of parasitic light. Under these conditions, it would be interesting to measure the variation of the peak positions with the strong laser power for different values of the detuning $\delta_2$. At resonance ($\delta_2 = 0$) the proportionality of the splitting to the square root of the light power has been observed several times [5].

The light-shift domain (large $\delta_2$) and the intermediate domain could be investigated in a similar way.

Acknowledgments. — The authors would like to thank Professor P. Berman for helpful discussions concerning the theoretical part of the work.

Appendix

The analytical expression for the quantities $A$, $B$, $C$ and $D$ (equation (2)) are given in this appendix. The model used for the atomic system and for the fields as well as the notations are the same as in the previous paper by the same authors [16] concerned with the optical Hanle effect.

\[ A = - \beta_1^2 \rho_0^2 G_{bb}(0, 0) X_1!(\det X_2)^2 \]
\[ - 2 \beta_1^3 \rho_2^2 \rho_0^2 \text{Re} \{ G_{bb}(0, 0) [G_{bb}(1, -1) + G_{cc}(1, -1) - g_{cb}(1, -1)] Z^2!(\det X_2)^2 \text{det } Y \} \]
\[ - \beta_2^2 \rho_0^2 G_{bb}(0, 0) X_2!(\det X_2) \]
\[ B = - \beta_1^2 \rho_2^2 \rho_0^2 [G_{bb}(0, 0, 0) - g_{cb}(0, 0)] X_1 X_2 [1 - \beta_2^2 G_{bb}(0, 0) X_2!(\det X_2)^2] \]
\[ - 2 \beta_1^3 \rho_2^2 \rho_0^2 \text{Re} [G_{bb}(0, 0) G_{bb}(1, -1) - 1] \{ Z - \beta_2^2 G_{bb}(0, 0) G_{bc}(1, 0) X_2 \} \{ G_{cb}(0, -1) \}
\[ + \beta_2^2 [G_{cc}(0, 0, 0) - g_{cb}(0, 0)] X_2 G_{bc}(1, 0)!(\det X_2)^2 (1 - \beta_2^2 G_{bb}(1, -1) G_{bc}(1, 0)) \]
\[ C = - \beta_1^2 \rho_2^2 G_{bb}(0, 0) X_1[1 - \beta_2^2 G_{bb}(0, 0) X_2!(\det X_2]\]
\[ - 2 \beta_1^3 \rho_2^2 \rho_0^2 \text{Re} [G_{bb}(0, 0) G_{bb}(1, -1) - 1) G_{bc}(1, 0) \]
\[ \times [Z - \beta_2^2 G_{bb}(0, 0) G_{bc}(1, 0) X_2]/(\det X_2[1 - \beta_2^2 G_{bb}(1, -1) G_{bc}(1, 0))] \]
\[ D = \beta_1^2 \rho_2^2 G_{bb}(0, 0) \{ X_1 - \beta_2^2 G_{bb}(0, 0) X_2 G_{bb}(1, 0) \}!(\det X_2[1 - \beta_2^2 G_{bb}(0, 0) G_{bc}(1, 0)] + \]
\[ + \beta_1^2 \rho_2^2 \rho_0^2 G_{bb}(0, 0) G_{bb}(1, -1) G_{bc}(1, 0) \]
\[ \times [Z - \beta_2^2 G_{bb}(0, 0) G_{bc}(1, 0)]!(\det X_2[1 - \beta_2^2 G_{bb}(0, 0) G_{bc}(1, 0)] + \beta_2^2 \rho_2^2 \rho_0^2 G_{bb}(0, 0) Z* [G_{bc}^*(1, -1) G_{cb}(-1, 0)
\[ + [G_{bc}^*(1, -1) - g_{cb}^*(1, -1)] Z*]!(\det X_2[1 - \beta_2^2 G_{bb}(0, 0) G_{bc}(1, 0)]) \text{det } Y* \]

The terms $g_{cb}(0, 0)$ and $g_{cb}(1, -1)$ are due to spontaneous emission and are respectively equal to:

\[ i\Gamma_s G_{cc}(0, 0) G_{bb}(0, 0) \text{ and } i\Gamma_c G_{cc}(1, -1) G_{bc}(1, -1) \cdot \]

$X_1$, $X_2$, $Z$ and $Y$ are sums of Green functions:

\[ X_1 = G_{bc}(1, 0) + G_{cb}(-1, 0) \]
\[ X_2 = G_{bc}(0, 1) + G_{cb}(0, -1) \]
\[ Z = G_{bc}(1, 0) + G_{cb}(0, -1) \]
\[ Y = G_{bc}(1, 0) + G_{cb}(1, -2) \]

and

\[ \det X_2 = 1 - \beta_2^2 [G_{bb}(0, 0) + G_{cc}(0, 0) - g_{cb}(0, 0)] X_2 \]
\[ \det Y = 1 - \beta_2^2 [G_{bb}(1, -1) + G_{cc}(1, -1) - g_{cb}(1, -1)] Y \]
\[ = \det Y*. \]
References