COLLAPSE OF SHOCK-BROADENED MULTIPLETS

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The collapse (merging) of the spectral components due to the averaging of the field by nonadiabatic collisions is discussed in the special case of the Stark structure of hydrogen and helium lines. It is shown that the phase "memory" responsible for this phenomenon changes the scale of the broadening by a substantial factor in the presence of degeneracy. It is also shown that the effect of exchange narrowing of the spectrum is an extreme case of collapse.

INTRODUCTION

WHEN shock broadening of the components of a multiplet is comparable with their separation the transformation of the spectrum by pressure and temperature becomes a very complicated phenomenon. This must be remembered whenever one is dealing with radiative transitions between groups of degenerate or quasidegenerate states. In atomic spectroscopy this condition is best satisfied by optical transitions in hydrogen and helium, whose Stark structure is resolved and is broadened by collisions with ions and electrons, respectively.[1-3] Competition between these processes at pressures and temperatures feasible in practice leads to a qualitative change in the spectra, which depends on which is the dominating effect. In particular, when the contribution of electrons to the broadening becomes very important, electron collisions lead to the merging of previously resolved spectrum components at the "center of gravity of the multiplet," and the side satellites tend to be suppressed. An analogous collapse may occur in the case of a pair of lines, one of which is forbidden, but this leads to the opposite result, i.e., the two lines assume equal intensity after merging.

The interdependent behavior of the various components of a complicated spectrum is a general consequence of the nonadiabatic character of the collisions if the resulting phase memory is not averaged. By inducing transitions between different sublevels such collisions not only modulate the phase and amplitude of the radiation, but also give rise to a frequency exchange.

In the special case when frequency exchange is not accompanied by a change in the relative phase of radiatively connected states, this leads to merging and subsequent contraction of shock-connected components of the spectrum. The possibility of this type of collapse was demonstrated earlier in[4] in the case of a four-level system in which the upper and lower level pairs were found to broaden in the same way. In the case of three levels this broadening of optically connected states was not possible because the lower (nondegenerate) state was not disturbed at all.

However, it will be shown below that, under favorable conditions, when transitions between the upper terms proceed through an intermediate level, we again have the necessary conditions for collapse just as in the case of the four-level system described earlier,[4] although the consequences of this are somewhat different in character. Such conditions are realized, in particular, in the case of the L_α and L_β lines and, to a considerable extent, are responsible for the smearing out of the Stark structure of these lines under conditions that are usual in the discharge plasma.[1-3]

1. NONADIABATIC BROADENING OF OVERLAPPING LINES

The general theory of the broadening of multicomponent spectra containing lines of close or equal frequency was developed in[4-6]. Its only serious limitation is the assumption of the shock character of the interaction between the particles and the medium. This condition is well satisfied in gases whose density is not too high. In the case of optical spectra, the theoretical formalism is somewhat simplified because collisions cannot produce optical transitions between states $|a\rangle$ and $|\alpha\rangle$, but merely induce relaxation within each of these groups of the degenerate or quasidegenerate terms ($|a\rangle - |a'\rangle$, $|\alpha\rangle - |\alpha'\rangle$). In such cases the form of the spectrum[4,5]

$$I(\omega)=\frac{1}{\pi}\operatorname{Re}\sum W_a\,\mathbf{d}_{a\alpha}{}^*\,\mathbf{d}_{b\beta}\,G^{-1}_{a\alpha b\beta} \tag{1.1}$$

is expressed in terms of the fundamental matrix

$$G_{a\alpha b\beta}=i(\omega-\omega_{a\alpha})\delta_{ab}\delta_{\alpha\beta}+\langle\delta_{ab}\delta_{\alpha\beta}-S_{ab}S_{\alpha\beta}{}^*\rangle, \tag{1.2}$$

whose relaxation part is equal to the product of S matrices for the upper and lower level groups, respectively, averaged over an ensemble of binary collisions. The remaining symbols have the following meanings: $\mathbf{d}_{a\alpha}$ are the elements of the dipole moment matrix responsible for the optical transitions, $\omega_{a\alpha}$ is their frequency, and W_a is the statistical weight of the state $|a\rangle$, which is unimportant for further analysis because the scale of the broadening and the line splitting which is comparable with it is definitely less than kT and is equal for all the components.

To be specific, let us consider the Stark effect in hydrogen and helium in the discharge plasma[1-3], which is a well-known phenomenon. In this case,[1,2] shock broadening of the line is largely due to perturbation by electrons of the upper levels which are readily polarized. Lower-level perturbation can be neglected in comparison with this (this assumption is rigorous for the Lyman lines because the ground state is spherically symmetric and not degenerate). A direct consequence of this fact is the linearization of the G matrix in S:

$$G_{\alpha\alpha\beta\beta} = G_{ab}{}^{\alpha} = i(\omega - \omega_{aa})\delta_{ab} - \langle T_{ab}\rangle, \tag{1.3}$$

$$\hat{T} = \hat{S} - \hat{1}. \tag{1.4}$$

Taking the interaction between the electron and the atom in the dipole approximation, using the classical description of the motion, and considering the effect of the collision as a resonance phenomenon,[2,7] we obtain the following expression for the T matrix in second-order perturbation theory:[2]

$$\hat{T} = -\frac{2}{3}\left(\frac{e^2 a}{\hbar \rho v}\right)^2 \mathbf{r}^2. \tag{1.5}$$

Averaging this expression with respect to the impact parameter ρ and velocity v does not alter the essence of the situation and affects only the absolute magnitude of the proportionality constant relating T and r^2 in Eq. (1.5):

$$\langle \hat{T}\rangle = -\frac{1}{3}\left(\frac{8\pi m}{kT}\right)^{1/2} N \left(\frac{\hbar}{m}\right)^2 \mathbf{r}^2 \int\limits_{y_{min}}^{\infty} \frac{e^{-y}}{y} dy = -g_n \mathbf{r}^2, \tag{1.6}$$

$$y_{min} = \frac{m v_{min}^2}{2kT} = \frac{4\pi}{3m} N \left(\frac{e\hbar n^2}{kT}\right)^2. \tag{1.6a}$$

In these expressions T is the temperature, N is the electron density, m is the electron mass, r^2 is in units of $a = \hbar^2/me^2$, and n is the principal quantum number. Therefore, the form of the G matrix is determined only by the matrix r^2.

The L_α spectrum. Neglecting the fine structure of the n = 2 level which, at densities in excess of 10^{14} cm^{-3} is negligible in comparison with the Stark structure, we obtain the well-known energy structure[8] in the ion field $\mathscr{E}$ (Fig. 1a):

$$E^{(1)} = 3ea\mathscr{E}(n_1 - n_2). \tag{1.7}$$

The two extreme components of the split level correspond to $n_1 = 1$, $n_2 = 0$ and $n_1 = 0$, $n_2 = 1$, whereas the central component with $n_1 = n_2$ remains doubly degenerate. If we evaluate the matrix elements of r^2 for the hydrogen functions in parabolic coordinates and use this result in Eqs. (1.3) and (1.6), we obtain

$$G_{ab} = \begin{pmatrix} i(\omega - \omega_{10}) + 18g & 0 & 0 & -9g \\ 0 & i(\omega - \omega_{20}) + 9g & 0 & 0 \\ 0 & 0 & i(\omega - \omega_{30}) + 9g & 0 \\ -9g & 0 & 0 & i(\omega - \omega_{40}) + 18g \end{pmatrix}. \tag{1.8}$$

We note that the presence of intermediate levels is a necessary condition for the appearance of the phase memory $\langle T_{ab}\rangle = -\Sigma g_2 r_{ac} r_{cb} = 9g_2$, which does not vanish after averaging with respect to the angles. This is illustrated by the transition scheme shown in Fig. 1a.

It is clear from the structure of the above matrix that the 2—0 and 3—0 lines are broadened independently of one another. Therefore, when Eq. (1.8) is substituted in Eq. (1.1), their shape is specified by the usual shock profile

$$I_a(\omega) = \frac{|\mathbf{d}_a|^2}{\pi} \frac{\Gamma_a}{(\omega - \omega_{a0})^2 + \Gamma_a^2} \quad (\Gamma_a = 9g,\ a = 2, 3), \tag{1.9}$$

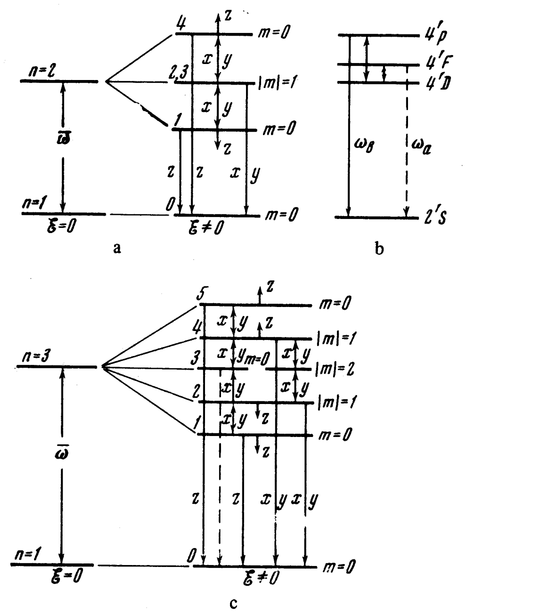

FIG. 1. Structure of the multiplets: a—L_α, b—HeI 3965Å, c—L_β; x, y, z are the components of the dipole moment inducing the transitions and polarization.

whose position and intensity are unaltered.

The situation is quite different for the other pair of lines, namely, 1—0 and 4—0 whose behavior is mutually connected through the nondiagonal matrix element representing the phase memory of the system and relating the two components when the $\hat{G}$ matrix is inverted. This problem is a special case of the more general model system, namely, the coupled doublet

$$G_{ab} = \begin{pmatrix} i(\omega - \bar{\omega} + {}^1\!/_2\,\Delta) + \Gamma_a & -\beta \\ -\beta^* & i(\omega - \bar{\omega} - {}^1\!/_2\,\Delta) + \Gamma_b \end{pmatrix} \tag{1.10}$$

and is obtained from it by substituting in Eq. (1.8) the following expressions:

$$\Gamma_a = \Gamma_b = \Gamma = 2\beta = 18g_2,$$
$$\Delta = \omega_{40} - \omega_{10} = 6ea\mathscr{E}/\hbar, \qquad \bar{\omega} = (\omega_{10} + \omega_{40})/2. \tag{1.11}$$

We shall show below that other pairs of lines in the spectra of hydrogen and helium can be fitted into the framework of this model.

The spectrum of λ = 3965 Å HeI. It is clear from Fig. 1b that, because of collisions with electrons, the allowed 4^1P-2^1S and forbidden 4^1F-2^1S transitions are coupled through the intermediate term 4^1D. Because of the substantial fine splitting, the Stark structure will be neglected in this case. Each of the three m components of the 4^1P level is coupled to the corresponding 4^1F component by analogy with the foregoing, so that the $\hat{G}$ matrix splits into three identical blocks of the form given by Eq. (1.10). Although the exact magnitude of the parameters will not be established in this case, the simplification of the problem to a simple model enables us to perform a qualitative analysis of the transformation of the spectrum, i.e., the appearance of the forbidden component.

The L_β spectrum. The Stark structure of the third level of the hydrogen atom in a field parallel to the z axis, and all the dipole transitions, namely, optical (to the ground state) and relaxation (between the components) are shown in Fig. 1c. The matrix of the operator given by Eq. (1.5) is diagonal in the quantum number m (as is, in general, the matrix of any scalar). It follows that the $\hat{G}$ matrix splits into blocks corresponding to two identical doublets (2—0, 4—0) and a triplet (1—0, 3—0, 5—0), the central component of which in this quantization is forbidden. As regards transitions from terms with $|m| = 2$, these are optically forbidden and are not relaxationally coupled to others. They do not, therefore, appear in the spectrum even when the collisions are taken into account. Equally so, the doublets and triplets are not coupled in any way, so that we can consider them separately. Direct evaluation of the $\hat{G}$ matrix for the n = 3 level over the states with $n = n_1 + n_2 + |m| + 1$ enables us to establish that blocks referring to the $|2\rangle - |4\rangle$ doublets have the same structure as that indicated by Eq. (1.10), and even with the same relationship between the parameters:

$$\Gamma_a = \Gamma_b = \Gamma = 2\beta = 81 g_3. \tag{1.12}$$

The situation is quite different in the case of the triplet. The corresponding $\hat{G}$ matrix in second-order perturbation theory is of the form

$$G_{ab} = \begin{pmatrix} i(x-\Delta)+\Gamma & -\beta & 0 \\ -\beta & ix+\Gamma-\delta & \beta \\ 0 & \beta & i(x+\Delta)+\Gamma \end{pmatrix}, \tag{1.13}$$

$$x = \omega - \bar{\omega} = \omega - {}^1\!/_2(\omega_{50} + \omega_{10}), \tag{1.14a}$$

$$\beta = \delta = {}^1\!/_3\Gamma = {}^{81}\!/_2 g_3. \tag{1.14b}$$

The appearance of the additional parameter δ is connected with the fact that the width of the forbidden line is different from that for the side components.

As regards splitting, its magnitude is given by the linear Stark effect[8], the theory of which yields

$$\Delta = 9ea\mathscr{E} \tag{1.15}$$

both for the triplet and the two doublets.

2. THE SPECTRUM OF THE DOUBLET

Let us now consider the three-level model, i.e., two optical transitions from different initial states to a common final state. It is important to note that this model has been discussed earlier in[4,5] in the hypothetical case where the phase memory appears in the first-order perturbation theory. Averaging of the multipole interactions in odd orders leads to the fact that it appears only in the second order and this modifies qualitatively all the conclusions because the nondiagonal elements of the $\hat{G}$ matrix turn out to have the same signs rather than different signs, as before.[4,5]

Inverting the matrix given by Eq. (1.10), and substituting the result in Eq. (1.1), we can readily show that the spectrum of the doublet consists of two Lorentz lines

$$I(\omega) = \frac{1}{\pi}\,\mathrm{Re}\left(\frac{A_+}{\omega-\omega_+} + \frac{A_-}{\omega-\omega_-}\right), \tag{2.1}$$

the position and width of which are given by the roots of the determinant of the matrix given by Eq. (1.10)

$$\omega_\pm = \bar{\omega} + i(\Gamma_a + \Gamma_b)/2 \pm \Omega, \tag{2.2}$$

whereas the intensity is given by

$$A_\pm = -i\frac{|\mathbf{d}_a|^2 + |\mathbf{d}_b|^2}{2} \mp \frac{\Gamma_a - \Gamma_b + i\Delta}{2\Omega}\,\frac{|\mathbf{d}_b|^2 - |\mathbf{d}_a|^2}{2} \mp \frac{\mathrm{Re}(\mathbf{d}_a^*\mathbf{d}_b\beta)}{\Omega} \tag{2.3}$$

where

$$\Omega = {}^1\!/_2\{[\Delta - i(\Gamma_a - \Gamma_b)]^2 - 4|\beta|^2\}^{1/2}. \tag{2.4}$$

It is readily seen that, in the absence of the phase memory, the broadening of both components of the spectrum would be independent, since for $\beta = 0$ the doublet is automatically reduced to that given by Eq. (1.9), i.e., two completely identical components which are obtained from each other by a permutation of the subscripts 1 and 2.

The presence in Eq. (2.4) of the phase-memory element ensures that this radical makes different contributions to the width and shift of the components, depending on the relative magnitude of the competing terms, and the appearance of the correction term proportional to β/ω in Eq. (2.3) violates the permutation symmetry of the line for $|\beta/\Omega| \sim 1$.

The most striking effect of the phase memory on the transformation of the spectrum is found in the special case

$$\Gamma_a = \Gamma_b, \tag{2.5}$$

to which the above doublet spectra of hydrogen reduce. This is clear from Eqs. (1.11) and (1.12). In this case, there is a strong alternative:

$$\Omega = \sqrt{(\Delta/2)^2 - |\beta|^2} \tag{2.6}$$

can be either real ($\Delta > 2|\beta|$) or purely imaginary ($\Delta < 2|\beta|$), so that in the first case Eq. (2.6) gives the shift of the component, whereas in the second it gives the correction to their width.

When both components of the spectrum are allowed, which occurs in hydrogen,

$$\mathbf{d}_a = \mathbf{d}_b = \mathbf{d} = 2^{-1/2}e\langle 1s_0|z|2p_0\rangle \quad (L_\alpha), \tag{2.7a}$$

$$-\mathbf{d}_a = \mathbf{d}_b = \mathbf{d} = 2^{-1/2}e\langle 1s_0|r_{x,y}|3p_{\pm 1}\rangle \quad (L_\beta), \tag{2.7b}$$

It then follows from the general formulas and from Eqs. (2.5) and (2.7) that

$$I(\omega) = \frac{|\mathbf{d}|^2}{\pi}\left[\frac{\Gamma \pm \beta\Omega^{-1}(\omega - \bar{\omega} + \Omega)}{(\omega-\bar{\omega}+\Omega)^2 + \Gamma^2} + \frac{\Gamma \mp \beta\Omega^{-1}(\omega-\bar{\omega}-\Omega)}{(\omega-\bar{\omega}-\Omega)^2+\Gamma^2}\right]$$
$$\Delta > 2\beta. \tag{2.8a}$$

$$I(\omega) = \frac{|\mathbf{d}|^2}{\pi}\left[\left(1 \pm \frac{\beta}{\Omega'}\right)\frac{\Gamma - \Omega'}{(\omega-\bar{\omega})^2 + (\Gamma-\Omega')^2} + \left(1 \mp \frac{\beta}{\Omega'}\right)\frac{\Gamma+\Omega'}{(\omega-\bar{\omega})^2+(\Gamma+\Omega')^2}\right]$$
$$\Delta < 2\beta, \quad \Omega' = |i\Omega|. \tag{2.8b}$$

The upper signs in Eq. (2.8) refer to the L_α and the lower to the L_β. It is clear from these formulas that as long as $2\beta < \Delta$ the spectrum consists of two split components which are asymmetric to an extent which depends on the ratio β/Ω. On the other hand, when $2\beta \geq \Delta$ the collapse sets in, and the two lines merge at the center of gravity of the spectrum, i.e., their shifts vanish. The doublet components differ from each other only in their width and intensity. It is noticeable that with decreasing $\Delta/2\beta$ one of the components disappears altogether (Fig. 2) and the spectrum becomes exactly the same as in zero fields. As a result, the spectrum of L_α contains only one narrow component whose half width is given by

$$\Delta\omega_{1/2} = \Gamma - \Omega' = \Gamma - \sqrt{\beta^2 - (\Delta/2)^2} \approx (\Gamma - \beta) + {}^1\!/_2\Delta^2/4\beta \quad (2.9a)$$

when $\Delta \ll 2\beta$. We note that fact that when there is a random compensation of the first terms of this formula there remains only the term which is of the second order of small quantities in $\Delta/2\beta$, and this decreases with increasing β, ensuring the narrowing of the spectrum by pressure:

$$\Delta\omega_{1/2} = \Delta^2/8\Gamma \qquad (\Gamma = \beta). \qquad (2.9b)$$

This type of phenomenon is known in molecular spectroscopy and is, in fact, observed after the averaging of the rotational structure of the nuclei[9] and electron paramagnetic resonances[10] by pressure. Its origin was discussed in detail within the framework of the four-level model[4], and it was shown that the necessary condition for its realization was $\Gamma = \beta$. It is not very probable that this condition will be accidentally satisfied and one would not, therefore, expect such a radical change in the nature of the broadening in atomic spectroscopy. However, the generality of the phenomenon is quite clear. This is so simply because of a different relation between Γ and β, established in Eqs. (1.11) and (1.12), so that the effect which is of zero order in $\Delta/2\beta$ predominates:

$$\Delta\omega_{1/2} = \frac{\Gamma}{2} + \frac{\Delta^2}{4\Gamma} \approx \frac{\Gamma}{2} \qquad (\Gamma = 2\ ). \qquad (2.9c)$$

Although line-narrowing does not occur in this case, the effect of the phase memory on the line width remains very important: it reduces it by a factor of two. Further increase in the ratio Γ/β, i.e., a reduction in the relative importance of the phase memory, results in the fact that the disproportion in the line width under the collapse conditions is smoothed out, as illustrated in Fig. 3.

A diametrically opposite effect in the sense of the behavior of the intensities of the narrow and broad components is obtained for L_β. It is clear from Eq. (2.8b) that when $\Delta \ll 2\beta$ the spectrum of L_β retains only the broad component whose half-width is given by

$$\Delta\omega_{1/2} = \Gamma + \Omega' \approx (\Gamma + \beta) - {}^1\!/_2\Delta^2/4\beta \approx {}^3\!/_2\Gamma \quad (\Gamma = 2\beta). \quad (2.10)$$

The narrowing of the line in this case is impossible in principle, even when the necessary condition $\Gamma = \beta$, mentioned above, is satisfied. The difference in the behavior of the collapsing doublets of the L_α and L_β lines is connected with the sign of γ_{ab} in Eq. (2.3)

$$|\mathbf{d}_a|\,|\mathbf{d}_b|\,\gamma_{ab} = \Omega^{-1}\operatorname{Re}(\mathbf{d}_a^*\mathbf{d}_b\beta) \sim \operatorname{Re}\sum_c (\mathbf{R}_b\mathbf{R}_a^*)(\mathbf{r}_{ac}\mathbf{r}_{cb}), \qquad (2.11)$$

which is conserved under any unitary transformation of the adopted representation (with the z axis along the field).

The quantity γ_{ab} can be interpreted as the mutual coherence function:[11] the value $\gamma_{ab} = 0$ corresponds to two independent oscillators ($\beta = 0$), whereas the state with full coherence $|\gamma_{ab}| \to 1$ corresponds to a correlated behavior of the two oscillators, and the spectrum of the collapsing doublet reduces to the form which it has in the absence of the field ($\Delta \ll 2\beta$).

When β is positive definite it can be interpreted as the frequency of the transitions between the fine-structure sublevels $|a\rangle$ and $|b\rangle$, and $\mathbf{d}_a$ and $\mathbf{d}_b$ are the oscillating polarizations at frequencies ω_a and ω_b. Collisions produce exchange between these frequencies without affecting the phase of the oscillations, but only when $\mathbf{d}_a$ and $\mathbf{d}_b$ have the same sign ($\gamma_{ab} > 0$). This leads to the narrowing of the line. On the other hand, when $\mathbf{d}_a$ and $\mathbf{d}_b$ have different signs, each collision results in a change of both frequency and phase ($\gamma_{ab} < 0$) and this in its turn results in a monotonic broadening of the line.[9,12] It is interesting to follow the behavior of the intensities of the doublet components, one of which is initially forbidden, as in the case of the helium atom:

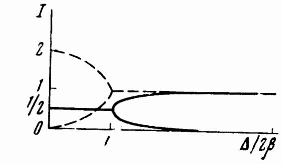

FIG. 2. Change in the intensities of the doublet components: broken curve—hydrogen (both transitions allowed), solid line—helium (one transition forbidden).

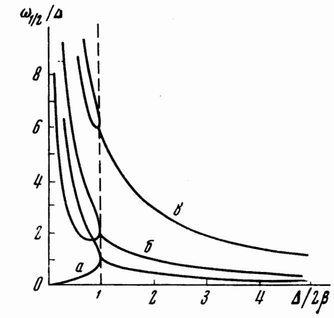

FIG. 3. Doublet line widths: a—$\Gamma = \beta$ (effect of narrowing), b—$\Gamma = 2\beta$, c—$\Gamma = 6\beta$. The collapse region lies between the ordinate axis and the broken line.

$$\mathbf{d}_b = \mathbf{d}, \quad \mathbf{d}_a = 0. \qquad (2.12)$$

Using the simplification given by Eqs. (2.5) and (2.12), we find from the general formulas that

$$I(\omega) = \frac{|\mathbf{d}|^2}{\pi}\left[\left(1 - \frac{|\beta|^2}{\Delta^2}\right)\frac{\Gamma}{(\omega-\omega_b)^2+\Gamma^2} + \frac{|\beta|^2}{\Delta^2}\frac{\Gamma}{(\omega-\omega_a)^2+\Gamma^2}\right],$$
$$2|\beta| \ll \Delta, \qquad (2.13a)$$

$$I(\omega) = \frac{|\mathbf{d}|^2}{\pi}\left[\frac{1}{2}\frac{\Gamma - \Omega' + (2\Omega')^{-1}\Delta(\omega-\bar\omega)}{(\omega-\bar\omega)^2 + (\Gamma-\Omega')^2}\right.$$
$$\left. + \frac{1}{2}\frac{\Gamma + \Omega' - (2\Omega')^{-1}\Delta(\omega-\bar\omega)}{(\omega-\bar\omega)^2 + (\Gamma+\Omega')^2}\right]$$
$$2|\beta| > \Delta, \quad \Omega' = |i\Omega|. \qquad (2.13b)$$

Although the widths and shifts of the lines behave in the same way as before, the spectroscopic picture in both limiting situations is substantially different. In the allowed structure [Eq. (2.13a)] there are two symmetric lines of equal intensity and the forbidden component is allowed,1) depending on the ratio $|\beta|^2/\Delta^2$.

As the phase memory increases, the intensities gradually become equal and eventually become identical when collapse sets in (Fig. 2). However, the shape of the spectrum remains relatively complicated because of the different widths and the asymmetric distortion of the merging components.

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1) This conclusion is unaffected when Eq. (2.5) is not satisfied, which is the case of a helium atom. All we need do is introduce the replacements $\Gamma \to \operatorname{Re}\Gamma_a$, $\operatorname{Re}\Gamma_b$ and, correspondingly, $\omega_a \to \omega_a + \operatorname{Im}\Gamma_a$, $\omega_b \to \omega_b + \operatorname{Im}\Gamma_b$ in Eq. (2.13a).

3. THE SPECTRUM OF THE TRIPLET

Inverting the matrix given by Eq. (1.13), and using it together with the formulas

$$\mathbf{d} = \mathbf{d}_{10} = \mathbf{d}_{50} = 2^{-1/2} e \langle 1s_0 | z | 3p_0 \rangle, \quad \mathbf{d}_{30} = 0 \tag{3.1}$$

in Eq. (1.1), we obtain

$$I(\omega) = \frac{|\mathbf{d}|^2}{\pi} \operatorname{Re} \frac{2z^2 - 2\beta z - 4\beta^2}{(z^2 + \Delta^2)(z - \delta) - 2\beta^2 z}, \tag{3.2}$$

where $z = ix + \Gamma = i(\omega - \overline{\omega}) + \Gamma$. It is clear from this result that when the phase memory is neglected ($\beta = 0$) the spectrum is always a doublet consisting of the two extreme components of equal intensity, since the central component in Eq. (3.1) is forbidden.

The situation is completely modified when phase memory is taken into account: the central line is partially allowed and it remains after the collapse. However, it is difficult to investigate this transformation of the spectrum in the general case and we shall simplify the problem by using the random coincidence of the parameters ($\delta = \beta$) established in Eq. (1.14b).

Since the mutual disposition of the lines in the spectrum is determined by the roots of the denominator in Eq. (3.2), the analysis of the structure of the spectrum reduces to the analysis of the discriminant of the cubic equation. Simple algebraic transformation yields

$$D(\eta) = -\frac{\Delta^6}{27}\left[\left(1 - \frac{7}{3}\eta^2\right)^3 + 3\eta^2\left(1 + \frac{10}{9}\eta^2\right)^2\right], \quad \eta = \frac{\beta}{\Delta}. \tag{3.3}$$

So long as the phase memory is small ($\eta < \eta_0$) the discriminant is negative and the equation has two complex and one real root:

$$z_\pm = \pm i\Delta(1 - \eta^2), \quad z_0 = \beta(1 + 2\eta^2) \quad (\eta \ll \eta_0). \tag{3.4}$$

The first two represent a doublet (two lines of equal intensity, separated by $\pm\Delta$) and the last represents the central component. When η increases the discriminant decreases in absolute magnitude, passes through the point $D(\eta_0) = 0$, and changes sign. When $\eta > \eta_0$ all the roots are real:

$$z_+ = \beta\left(2 - \frac{1}{6\eta^2}\right), \quad z_- = \beta\left(-1 + \frac{2}{3\eta^2}\right), \quad z_0 = -\frac{3}{2\eta^2} \quad (\eta \gg \eta_0). \tag{3.5}$$

Thus, here again we encounter collapse: the total merging of the spectrum components at the "center of gravity." The point at which collapse occurs is defined by the condition

$$D(\eta_0) = 0, \quad \eta_0 = 1.55, \tag{3.6}$$

which is similar to that which we have in the case of the doublet ($\eta_0 = 1/2$).

If we now consider the above roots of the denominator in Eq. (3.2) we can readily establish from this formula the spectra which correspond to the two limiting situations:

$$I(\omega) = \frac{|\mathbf{d}|^2}{\pi}\left[(1 - 2\eta^2)\frac{\Gamma}{(\omega - \overline{\omega} - \Delta)^2 + \Gamma^2} + (1 - 2\eta^2)\frac{\Gamma}{(\omega - \overline{\omega} + \Delta)^2 + \Gamma^2} + 4\eta^2 \frac{2\Gamma/3}{(\omega - \overline{\omega})^2 + (2\Gamma/3)^2}\right] \tag{3.7a}$$

when $\beta/\Delta \ll \eta_0$, and

$$I(\omega) = \frac{|\mathbf{d}|^2}{\pi}\left[\left(2 + \frac{3}{2\eta^2}\right)\frac{\Gamma}{(\omega - \overline{\omega})^2 + \Gamma^2} - \frac{1}{6\eta^2} \cdot \frac{\Gamma/3}{(\omega - \overline{\omega})^2 + (\Gamma/3)^2} - \frac{4}{3\eta^2}\frac{4\Gamma/3}{(\omega - \overline{\omega})^2 + (4\Gamma/3)^2}\right] \tag{3.7b}$$

when $\beta/\Delta \gg \eta_0$. The line widths are determined here with allowance for Eq. (1.14b) ($\Gamma = 3\beta$). In the first case, as the ratio $(\beta/\Delta)^2$ is allowed to vary the intensity is transferred to the forbidden component of the spectrum and all three components are quite symmetric and have the Lorentz shape. In the opposite case, collapse is eventually (when $\eta \gg \eta_0$) found to result in a single line of twice the intensity, but as we return to the point η_0 from the side of large η the spectrum is found to be transformed. Because of the second term in Eq. (3.7b) (narrow line with negative weight) there is a valley at the center of the spectrum which becomes deeper as η decreases, and only after passage through η_0 when the extreme lines separate by $\pm\Delta$ do we find at this point an additional maximum corresponding to the forbidden component.

CONCLUSIONS

The phase memory of an atom represented by the nondiagonal elements of G_{ab} was shown above to produce exchange between the corresponding spectral lines by analogy with the situation when the nondiagonal elements of a Hamiltonian mix different energy states. It is, therefore, natural that its role increases when close frequencies are involved in the exchange. Since the phase memory scale is determined by the line broadening ($\beta \leq \Gamma$), it is clear that this must be taken into account whenever one considers poorly resolved structures or radiation from a degenerate term.

When the phase memory is taken into account the result is that collisions do not simply smear out the structure and broaden its components but, in fact, transform it to the form which one would expect in the absence of splitting. In other words, collisions completely average the field acting on the atom. This was to be expected: exchange between the multiplet components takes place because nonadiabatic collisions take the atom from one Stark level to another, change the sign of the dipole moment, and thus average the energy and remove the polarization of the states which appears in the field. In this sense, the collapse of the Stark structure is quite similar to the averaging of local magnetic fields because of exchange[9] or rotational structure and reorientations of the angular momenta due to collisions.[13,14] The appearance of the effect is not therefore connected with the number of levels involved in the process but with the fact $G_{ab} = G^*_{ba}$ ($a \neq b$) (whereas in[4,5] $G_{ab} = -G^*_{ba}$).

However, the narrowing of the lines is possible only when $\langle T_{aa}\rangle = \langle T_{bb}\rangle = \langle T_{ab}\rangle$, $\gamma_{ab} > 0$, Eqs. (2.9b) and (2.11).

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