above, but they have no evidence of a line near 780 Å. Our results are in agreement with their conclusions. For example, from Table XIV we see that there is a series line at 685.46 Å corresponding to the transition $^3P_2$-$(^2P^\circ)5d''$ $^3P^\circ$ and one at 685.44 Å a corresponding to $^3P_2$-$(^2P^\circ)5d''$ $^3D^\circ$.

The Rydberg series of levels, which belong to a discrete spectrum and which precede each threshold in the photon energy scale, provide an indirect method of ionization by absorption into a short-lived state followed by auto-ionization. Thus, peaks in the absorption cross section can be expected near (or at) each auto-ionization level. The analysis of these resonances will be treated elsewhere.

The good quantitative agreement between theory and experiment for the absorption-line series in atomic oxygen justifies our retention of only the terms of the ground-state configuration in the close-coupling expansion. It is noted that the same approximations were made for the atoms considered in the previous sections.

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Collisional Depolarization of the Rb 5*p* and Cs 6*p* Doublets*

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Using the Hanle-effect technique, the cross sections for collisional depolarization of the Rb $^2P_{1/2,\ 3/2}$ and Cs $^2P_{1/2}$ states by inert-gas collisions have been measured. (The natural lifetimes of these states were also obtained from the measurements.) The $^2P_{1/2}$-state depolarization cross sections are much smaller than any previously reported for $L\neq0$ states. This results from a strong selection rule governing nonresonant collisions. This selection rule is derived and its effect on $J=\frac{1}{2}$-state depolarization is demonstrated. A relationship between the reported cross sections and sensitized fluorescence cross sections is noted.

I. INTRODUCTION

THE collisional depolarization of atomic states by thermal atom-atom collisions has been calculated in several formalisms.[1-3] The theories have considered isolated fine-structure states, with hyperfine components, using the adiabatic approximation to evaluate cross sections for foreign gas (nonresonant) depolarization. Franz and Franz have discussed the effect of the fine-structure interval in the nonresonant collisional depolarization of the alkali *p* doublets, and the implications to optical pumping of a $\Delta m_J=0,\pm2$ selection rule that occurs in several theories for collisional transfer between the J state of these doublets.[4] It is demonstrated here that there exists in addition to that weak selection rule, a strong selection rule governing depolarization within a J level by inert gas collisions. This strong selection rule asserts that in the adiabatic approximation the entire electrostatic collisional interaction (including exchange) will not cause transitions between states JM and $J-M$ when J is half-integral. The effect of this selection rule on $J=\frac{1}{2}$-state depolarization is very apparent, since it requires that the depolarization must be produced by smaller nonadiabatic contributions to the interaction. We present experimental results for the nonresonant collisional depolarization cross sections σ of the Rb 5*p* and Cs 6*p* fine-structure levels, verifying that the $J=\frac{1}{2}$ state cross sections are unusually small. A discussion of the extension of the present adiabatic theories to the $J=\frac{1}{2}$ problem is included, and the meaning of the Δm_J selection rules in the adiabatic theories is reviewed. The effect of the reported cross sections on ground-state optical pumping can be established by use of formulas for relaxation of hyperfine components in Ref. 2 in conjunction with procedures in Ref. 4; no discussion is included.

The theories have treated the collision Hamiltonian V as a perturbation that alters the populations of the electronic states of the separated atoms, using the dipole-dipole interaction and straight-line paths for $V(\mathbf{b},\mathbf{v},t)$ of a collision with impact parameter $\mathbf{b}$ and relative velocity $\mathbf{v}$. For this V, the nonresonant collisional depolarization arises from the Van der Waals interaction, and if $[V^2(\mathbf{b},\mathbf{v},t),\ V^2(\mathbf{b},\mathbf{v},t')]$ is neglected,

$$S(\mathbf{b},\mathbf{v})=S_{dd}(\mathbf{b},\mathbf{v})\equiv\exp\left(-i\int_{-\infty}^{\infty}dt\ V_{dd}^2(\mathbf{b},\mathbf{v},t)/\hbar\Delta E_{\mathrm{av}}\right)$$

is obtained,[2] where we represent the effect of one collision by $\psi(\infty)=S(\mathbf{b},\mathbf{v})\ \psi(-\infty)$. This integral can be evaluated in the collision frame ($\mathbf{b}$ and $\mathbf{v}$ along coordi-

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[1] F. W. Byron and H. M. Foley, Phys. Rev. **134**, A625 (1964). See Eq. (1) for the linewidths in doublet resonance or Hanle-effect experiments.

[2] Alain Omont, J. Phys. Radium **26**, 26 (1965).

[3] M. I. D'Yakonov and V. I. Perel', Zh. Eksperim. i Teor. Fiz. **48**, 345 (1965) [English transl.: Soviet Phys—JETP **21**, 227 (1965)].

[4] F. Franz and J. Franz, Phys. Rev. **148**, 82 (1966).

nate axes); then the integrals over collision directions, velocity distribution, and impact parameters lead to the observed damping rates $\gamma_i = n\bar{v}\sigma_i$ for the orientation ($i=1$) and alignment ($i=2$) of the excited-state density matrix referred to any laboratory frame (n is the foreign gas density, $\bar{v}$ is the average relative velocity).[2] Since the integral is proportional to b^{-5}, most of the depolarization is caused by "strong" collisions for which the integral exceeds 1 (b_0 is defined as the b at which the integral equals unity). In strong collisions, neither $[V_{dd}^2(\mathbf{b},\mathbf{v},t),\ V_{dd}^2(\mathbf{b},\mathbf{v},t')]$ nor the difference between the complete electrostatic interaction and V_{dd} are negligible, and $S_{dd}(\mathbf{b},\mathbf{v})$ is not a valid solution. The calculated cross sections are generally insensitive to this because the average effect of strong collisions is approximately complete mixing of all the m_J levels of each J state, whether S_{dd} or a more appropriate S is used. Thus, because the Van der Waals interaction (and presumably also the actual interaction) is a rapidly varying function of internuclear separation, the cross sections are about πb_0^2, typically $\sim 10^{-14}$ cm^2, regardless of minor difference between S_{dd} and S when $b \sim b_0$. But the ratio γ_1/γ_2 for a particular J state depends only on the ratios of the $\langle Jm_J|S|Jm_J'\rangle$ for the collision frame states $|Jm_J\rangle$, and if S_{dd} is used, the ratios obtained are not the same as those from the more appropriate model of complete m_J mixing by all strong collisions.[5] The largest difference arises because $\langle Jm_J|S_{dd}|Jm_J'\rangle = 0$ unless $m_J - m_J'$ is an even integer. This selection rule would also have important consequences in the case of a $J=\frac{1}{2}$ state because the collision-frame excited-state density matrix, and thus the laboratory-frame density matrix, will not be altered by collisions unless they cause transitions between the $m_J = \pm\frac{1}{2}$ levels (the phase shifts in both m_J levels will be the same). If these transitions were forbidden, $\sigma = 0$ would result. Since S_{dd} does not apply to strong collisions, one could hardly expect from the S_{dd} selection rule that σ of a $J=\frac{1}{2}$ state would differ greatly from 10^{-14} cm^2. But there is a much stronger argument against the J, $m_J = \frac{1}{2}, \frac{1}{2} \leftrightarrow \frac{1}{2}, -\frac{1}{2}$ transitions, and as a result σ of a $J=\frac{1}{2}$ state ($L>0$) can be much less than 10^{-14} cm^2. We will elaborate on this by considering the problem of the Rb $5p$ and Cs $6p$ doublets.

When considering the nonresonant collisional depolarization of atomic states characterized by Russell-Saunders coupling ($L \neq 0$), two extreme cases occur with the usual problematical region in between. When the fine-structure separation ΔE_J is much less than $\hbar/\tau_c$, $\mathbf{S}$ remains stationary while the orientation of $\mathbf{L}$ is changed by a collision ($\tau_c \cong \int V(\mathbf{b},\mathbf{v},t)dt/V_{\max}$ is the collision time). The resulting cross sections can be calculated using the results of Ref. 2, with $\mathbf{L}$ and $\mathbf{J}$ in place of $\mathbf{J}$ and $\mathbf{F}$; of course $\sigma \sim 10^{-14}$ cm^2 results for all the levels of such a multiplet. When $\Delta E_J \gg \hbar/\tau_c$, $\mathbf{S}$ and $\mathbf{L}$ remain coupled during a collision and $\mathbf{J}$ must be depolarized; the results in Ref. 2 may be used directly for this case. The ratios of the σ for depolarization within different levels of a multiplet will differ in this regime, but $\sigma \sim 10^{-14}$ cm^2 again results for $J > \frac{1}{2}$ cases. Thus collisional depolarization cross sections normally change by small amounts between these two regimes, whereas the cross sections for transfer between the levels separated by ΔE generally undergo large variations. Collisional transfer between the J levels of the lowest alkali p doublets is an example closely related to the measurements reported here. For Li and Na, $\Delta E_J < \hbar\tau_c$, so that inert gas collisions can change the orientation $\mathbf{L}$ without changing $\mathbf{S}$, thereby causing ΔJ transitions with cross sections $\sigma_{\Delta J}$ near 10^{-14} cm^2. But for Rb and Cs, $\Delta E_J \gg \hbar/\tau_c$, so that $\mathbf{J}$ is depolarized, and therefore the ΔJ transitions are not caused by the long-range interaction; $\sigma_{\Delta J} \ll 10^{-14}$ cm^2 results.[6,7] The depolarization of these ${}^2P_{1/2}$ states of the alkalis also tests the effects of ΔE_J versus $\hbar/\tau_c$ on collisional depolarization. The Li and Na ${}^2P_{1/2}$ states should have $\sigma \sim 10^{-14}$ cm^2 for inert-gas collisions. The Rb and Cs ${}^2P_{1/2}$ states must be depolarized primarily by J, $m_J = \frac{1}{2}, \frac{1}{2} \leftrightarrow \frac{1}{2}, -\frac{1}{2}$ transitions, but as noted already this will have an unusually small cross section. Thus the transition from $\mathbf{L}$ disorientation to $\mathbf{J}$ disorientation causes a drastic reduction in σ of these ${}^2P_{1/2}$ states, whereas the same transition would not greatly alter σ for a $J > \frac{1}{2}$ state.

II. EXPERIMENT

We have measured cross sections for the collisional depolarization of the Rb^{87} $5p$ ${}^2P_{1/2}$ and ${}^2P_{3/2}$ states and the Cs $6p$ ${}^2P_{1/2}$ state by inert-gas collisions. We used the Hanle-effect method because it offers several advantages over a direct measurement of the polarization of the scattered light, but otherwise obtains the same information. The principal advantages are the independence from corrections for instrumental scattering, a comparatively small dependence on optical depth, and a clear magnetic-field separation of the effects of ground-state pumping in the buffered cells. The Hanle-effect linewidths were obtained for each inert-gas pressure in the limit of zero optical depth. In contrast to zero-field polarization measurements, this was very easy because the nuclear spins of Rb^{87} ($I=\frac{3}{2}$) and Cs ($I=\frac{7}{2}$) make the Hanle-effect widths comparatively free from coherence narrowing. At the highest temperatures used, 22°C, less than 5% narrowing was obtained even from 3-in.-diam cells without buffer gas. The cross sections were obtained by fitting the Hanle-effect linewidths ΔH to $\Delta m|g_F|\mu_0\Delta H/\hbar = 1/\tau + n\bar{v}\sigma$,[1] where Δm is 2 or 1 for alignment or orientation, τ is the natural lifetime of the p state, and $|g_F|$ is the same for all the F levels of each p state investigated. In this relationship we have neglected the ($n=0$) line-shape corrections of about 1%

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[5] See Sec. II.D of Ref. 2 for a discussion of the strong-collision mixing and the two approximations for its evaluation. The Δm_J selection rule can be seen from Eq. (20).

[6] Bruce Pitre, A. G. A. Rae, and L. Krause, Can. J. Phys. 44, 731 (1966).

[7] C. Czajkowski, D. A. Mcgillis and L. Krause, Can. J. Phys. 44, 91 (1966).

TABLE I. Collisional depolarization cross sections.

State	Foreign gas	σ (10^{-16} cm^2) Orientation	Alignment
Rb $5p\ ^2P_{3/2}$	Neon	57	100
	Argon	130	210
Rb $5p\ ^2P_{1/2}$	Helium	9.0	
	Neon	6.0	
	Argon	9.7	
	Krypton	10.6	
Cs $6p\ ^2P_{1/2}$	Helium	2.1	
	Neon	0.8	
	Argon	1.7	

due to the finite size of the level width compared to the hfs separations. An example of the data is given in Fig. 1, the measured cross sections are given in Table I, and the natural lifetimes obtained from the $n=0$ linewidths are reported in the last paragraph. The reported cross sections represent an average over those for depolarization within and between different hyperfine components, with the relative pumping rates of different components somewhat uncertain due to uncertainties in the (electrodeless bulb) lamp profile. The relative sizes of these different cross sections within one J state typically differ by 20 to 50%, hence this is not important to the present discussion.[2] The (average) σ_2 was obtained by scattering linear polarized light and observing the $\Delta m_F=2$ crossing signals; the (average) σ_1 by scattering circular polarized light and measuring only the $\Delta m_F=1$ crossing signals (we did not measure σ_2 of the $^2P_{1/2}$ states because they cannot be aligned by simple optical pumping, nor the Cs $^2P_{3/2}$-state signals because they are complicated by different $|g_F|$ values). It can be seen that as ΔE_J increases compared to $\hbar/\tau_c$, σ of the $^2P_{1/2}$ states becomes one to two orders of magnitude less than 10^{-14} cm^2. Yet σ remains orders of magnitude greater than the known $\sigma_{\Delta J}$ and its magnitude must be explained by depolarization within the $J=\frac{1}{2}$ levels.[6,7]

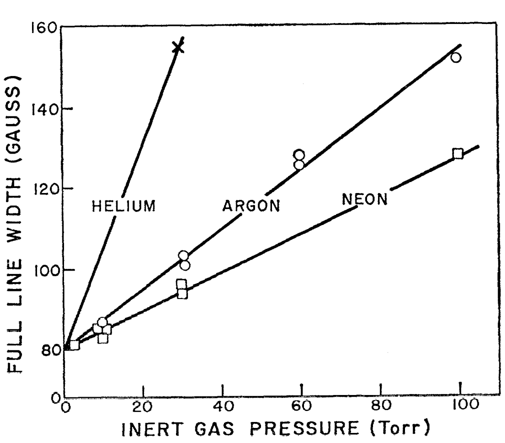

FIG. 1. Cs $6p\ ^2P_{1/2}$ state Hanle-effect linewidths (zero Cs density limits).

III. THEORY

As noted above, the $^2P_{1/2}$ states must be depolarized by collisional transfer between the $m_J=\pm\frac{1}{2}$ levels, so we will now consider this transfer by a collision of arbitrary **b** and **v**. We will consider the complete electrostatic interaction, neglecting magnetic interactions with the alkali spin since these are known to produce extremely small cross sections for ground-state depolarization. If we expand the state of the outer alkali electron and the inert gas outer-shell electrons in antisymmetrized products of atomic orbitals, this will be almost equivalent to using single products of atomic orbitals and adding an exchange interaction $-\sum_k r_{1k}^{-1}p_{1k}$ to the direct collisional interaction $\sum_k V_{1k}$.[8] Here P_{1k} exchanges the alkali electron, 1, with the kth electron of the inert gas outer shell, and $V_{1k}\equiv r_{1k}^{-1}-r_{1N}^{-1}-r_{kN'}^{-1}+r_{NN'}$, where N and N' refer, respectively, to the positions of the inert gas and alkali nuclei. Due to the nonorthogonality of the atomic states on different centers, some additional smaller exchange contributions occur. We will discuss them below. The essential feature of this interaction is its spin independence, and this should be appropriate even during the parts of a collision when the actual state is a considerably mixed set of atomic states. In the interaction representation, the probability amplitudes $a_m(t)$ of the states $|mg\rangle$ are given by

$$\begin{aligned} i\hbar\dot{a}_m(t)=&\langle|V(t)gm|m'g\rangle a_{m'}(t)\\ &+i\hbar^{-1}\sum_{p,q}e^{-i\omega_{pq,ng}t}\langle mg|V(t)|pq\rangle\\ &\times\int_{-\infty}^{t}dt'e^{i\omega_{pq,ng}t'}\langle pq|V(t')|m'g\rangle a_{m'}(t')+\cdots. \end{aligned} \quad (1)$$

Here $|mg\rangle$ is a single product,

$$|(\tfrac{1}{2}m)_1\prod_{k=1}^{N}(g^k)_{k+1}\rangle,$$

of the $n^2P_{1/2\ m}(\mathbf{r}_1)$ state of the alkali electron and the ground-configuration single-electron states, $g^k(\mathbf{r}_{k+1})$, of the N outer-shell electrons of the inert gas; p is another alkali state, q is one state of another configuration of the inert gas, $\hbar\omega_{pq,ng}=(E_p+E_q)-(E_n+E_g)$, and V is the sum of the direct and exchange interactions. We will limit the discussion to the $m'=-\frac{1}{2}$, $m=\frac{1}{2}$ terms in (1). The first-order interaction in (1) becomes[8]

$$\begin{aligned} \langle\tfrac{1}{2}g|V|-\tfrac{1}{2}g\rangle\equiv&\langle(\tfrac{1}{2}\,\tfrac{1}{2})_1\prod_{k'}(g^{k'})_{k'+1}|\sum_k(V_{1k}-r_{1k}^{-1}P_{1k})|\\ &\times|(\tfrac{1}{2}\,-\tfrac{1}{2})_1\prod_{k'}(g^{k'})_{k'+}\rangle=\sum_k\langle(\tfrac{1}{2}\,\tfrac{1}{2})_1(g^k)_{k+1}\\ &\times|V_{1k}-r_{1k}^{-1}P_{1k}|(\tfrac{1}{2}\,-\tfrac{1}{2})_1(g^k)_{k+1}\rangle\equiv\sum_{k=1}^{N}\langle(\tfrac{1}{2}\,\tfrac{1}{2})_1\\ &\times(j^km^k)_2|V_{12}-r_{12}^{-1}P_{12}|(\tfrac{1}{2}\,-\tfrac{1}{2})_1(j^km^k)_2\rangle, \end{aligned} \quad (2)$$

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[8] E. U. Condon and G. H. Shortley, *The Theory of Atomic Spectra* (Cambridge University Press, New York, 1959); see Sec. 7 except as noted.

where the inert gas single-electron states have been characterized by jm quantum numbers in the final expression. Using the definition of single-electron eigenfunctions in Eqs. (14^31) and (4^52) of Ref. 8, $Y_{lm}^* = (-1)^m Y_{l,-m}$ from Sec. 4^3 of Ref. 8, and $(lm_lSm_S|jm) = (-1)^{l+S-j}(l-m_lS-m_S|j-m)$, it can be verified that if the operator $\mathcal{V}$ satisfies

$$\int\int \phi_1^*(1)\phi_2^*(2)\{\mathcal{V}\phi_3(1)\phi_4(2)\} = \int\int \{\mathcal{V}\phi_1^*(1)\phi_2^*(2)\}\phi_3(1)\phi_4(2),$$

then

$$\langle (l^1j^1m^1)_1(l^2j^2m^2)_2|\mathcal{V}|(l^3j^3m^3)_1(l^4j^4m^4)_2\rangle = \{\prod_{i=1}^{4}(-1)^{l^i-j^i-m^i}\}\langle (l^3\,j^3\,-m^3)_1(l^4\,j^4\,-m^4)_2|\mathcal{V}| \times (l^1\,j^1\,-m^1)_1(l^2\,j^2\,-m^2)_2\rangle. \quad (3)$$

Since P_{12} can exchange $\mathbf{r}_1$ and $\mathbf{r}_2$ of either pair of states in the matrix elements of (2), $V_{12}-r_{12}^{-1}P_{12}$ satisfies the condition on $\mathcal{V}$ and we define it as $\mathcal{V}$ below (note that a general spin-dependent interaction would not satisfy this condition on $\mathcal{V}$). Since all the configurations of single-electron states appear an even number of times in all orders of (1), the various $(-1)^{l^1+l^2+l^3+l^4}$ that occur when (3) is used in (1) will always reduce to inconsequential $+1$ factors. For brevity we will delete the l quantum numbers in the following discussion. Since the inert gas outer shell is filled, the sum over k in (2) will include a $j^km^k = j\,-m$ state for every jm state. If we apply (3) to such a pair from the k sum in (2), we obtain

$$\sum_{\pm}\langle (\tfrac{1}{2}\,\tfrac{1}{2})_1(j\,\pm m)_2|\mathcal{V}|(\tfrac{1}{2}\,-\tfrac{1}{2})_1(j\,\pm m)_2 = \sum_{\pm}(-1)^{-2j-1\mp 2m} \times\langle (\tfrac{1}{2}\,\tfrac{1}{2})_1(j\,\mp m)_2|\mathcal{V}|(\tfrac{1}{2}\,-\tfrac{1}{2})_1(j\,\mp m)_2\rangle. \quad (4)$$

Since j and m are half-integral, $(-1)^{-2j-1\mp 2m} = -1$ and one side of (4) is the negative of the other. Thus both sides must be zero, and $\langle \tfrac{1}{2}g|V|-\tfrac{1}{2}g\rangle$ in (1) is zero for all collisions.

The second-order part of (1) reduces to[8]

$$\frac{1}{i\hbar}\int_{-\infty}^{t} dt' \sum_{g,q,p} e^{i(t'-t)\omega_{pq,ng}}\langle (\tfrac{1}{2}\,\tfrac{1}{2})_1(j^gm^g)_2 \times|\mathcal{V}(t)|(j^pm^p)_1(j^qm^q)_2\rangle\langle (j^pm^p)_1(j^qm^q)_2|\mathcal{V}(t')| \times(\tfrac{1}{2}\,-\tfrac{1}{2})_1(j^gm^g)_2\rangle a_{-1/2}(t'). \quad (5)$$

The excitation energy $\hbar\omega_{pq,ng}$ now depends on j^pm^p, j^qm^q, and j^gm^g (within each configuration), but the signs of m^p, m^g, and m^g can be reversed without significantly changing $\omega_{pq,ng}$. Consequently, the matrix elements in (5) that differ only by the signs of m^p, m^g, and m^q are summed with a common coefficient. Using (3), we find for such a pair of terms in (5):

$$\sum_{\pm}\langle (\tfrac{1}{2}\,\tfrac{1}{2})_1(j^g\,\pm m^g)_2|\mathcal{V}(t)|(j^p\,\pm m^p)_1(j^q\,\pm m^q)_2\rangle \times\langle (j^p\,\pm m^p)_1(j^q\,\pm m^q)_2|\mathcal{V}(t')|(\tfrac{1}{2}\,-\tfrac{1}{2})_1(j^g\,\pm m^g)_2\rangle$$

$$\sum_{\pm}(-1)^{-1-2j^g-2j^q-2j^p\mp 2m^g\mp 2m^q\mp 2m^p}\langle (j^p\,\mp m^p)_1 \times(j^q\,\mp m^q)_2|\mathcal{V}(t)|(\tfrac{1}{2}\,-\tfrac{1}{2})_1(j^g\,\mp m^g)_2\rangle \times\langle (\tfrac{1}{2}\,\tfrac{1}{2})_1(j^g\,\mp m^g)_2|\mathcal{V}(t')|(j^p\,\mp m^p)_1(j^q\,\mp m^q)_2\rangle. \quad (6)$$

The exponent of the -1 will be an odd integer, and because each intermediate state always appears an even number of times this will be a general result for all orders in (1). The final expression in (6) would thus be the negative of the first one if $\mathcal{V}(t)$ were equal to $\mathcal{V}(t')$, and since the entire sum of the type in (5) reduces to such pairs of states, it will be zero if $\mathcal{V}(t)=\mathcal{V}(t')$.

The smaller exchange terms that arise due to the nonorthogonality of the atomic states on different centers add to (2) some matrix elements of the form $\langle (\tfrac{1}{2}\,\tfrac{1}{2})_1(g^k)_2(g^{k'})_3|r_{13}^{-1}|(\tfrac{1}{2}\,-\tfrac{1}{2})_2(g^k)_1(g^{k'})_3\rangle$. It is clear that the extension of (3) and the succeeding steps to these terms also leads to a null result. Similarly the extra second-order terms sum to zero when $\mathcal{V}(t)=\mathcal{V}(t')$. The importance of this cancellation to these slow collisions is best demonstrated by integrating the second-order term in (1) by parts to obtain

$$-\sum_{pq}\langle \tfrac{1}{2}g|V(t)|pq\rangle\langle pq|V(t)|-\tfrac{1}{2}g\rangle a_{-1/2}(t)/\hbar\omega_{pq,ng} + \sum_{pq} e^{-i\omega_{pq,ng}t}\langle \tfrac{1}{2}g|V(t)|pq\rangle\int_{-\infty}^{t} dt' \times e^{i\omega_{pq,ng}t'}\left\langle pq\left|\frac{d}{dt'}V(t')a_{-1/2}(t')\right|-\tfrac{1}{2}g\right\rangle \Big/ \hbar\omega_{pq,ng}. \quad (7)$$

In the adiabatic approximation only the first sum, which includes the Van der Waals interaction, would be used, but the V are evaluated at the same time in that sum so it is zero in the present case. The applicability of this result to higher order terms in (1) is clear; the conclusion reached is that repeated application of the adiabatic approximation leads to a null result for $\tfrac{1}{2}\,\tfrac{1}{2}$ to $\tfrac{1}{2}\,-\tfrac{1}{2}$ mixing. This conclusion clearly applies to any JM to $J\,-M$ mixing when J is half-integral.[9]

We will now investigate the residual (nonadiabatic) interaction in the second sum of (7). By inspection, or another parts integration, it may be verified that the pq term in the second sum of (7) is about $(\omega_{pq,ng}\tau_c)^{-1}$ as large as the pq term in the first sum ($dV/dt \sim V/\tau_c$). Because of the spherical symmetry of the inert-gas ground state, matrix elements of the direct electrostatic interaction between $\langle pg|$ and $|p'g\rangle$ are zero. But if q is

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[9] Since $m_J \rightarrow m_J' \rightarrow -m_J$ transitions will occur in the adiabatic approximation to $J > \tfrac{1}{2}$ states, the average effect of strong collisions should still be almost complete m_J mixing when $J > \tfrac{1}{2}$.

an excited inert-gas configuration $(\omega_{pq,ng}\tau_c)^{-1}$ is about 10^{-4}, whereas $(\tau_c\omega_{n^2P_{3/2}g,n^2P_{1/2}g})^{-1}\equiv(\Delta E_J\tau_c/\hbar)^{-1}$ is between 10^{-1} and 10^{-2}. It appears therefore that in the second sum of (7), the matrix elements of the exchange interaction to the intermediate $|n^2P_{3/2\ m}g\rangle$ state should make the greater contribution to the depolarization. In addition, the direct interaction can mix in $|n^2P_{3/2\ m}g\rangle$ as the second intermediate states in fourth-order perturbation theory. Using the adiabatic approximation and closure for the first and third intermediate states (but not for the second), one can obtain a modified (7) to represent this fourth-order interaction with $V^2(t)/\Delta E_{\mathrm{av}}$ in place of $V(t)(\Delta E_{\mathrm{av}}\sim 15$ eV.) [The first sum in the modified (7) will still be zero.] From order-of-magnitude arguments, it can be established with reasonable certainty that the fourth-order mixing of the $|{}^2P_{3/2\ m}g\rangle$ intermediate state will cause more depolarization than the direct second-order interaction, which cannot reach the $|{}^2P_{3/2\ m}g\rangle$ intermediate state. Although we are not at this time able to estimate the relative importance of the fourth-order direct and second-order exchange mixing of the $|n^2P_{3/2\ m}g\rangle$ intermediate states, it seems clear that they should dominate the depolarization. These $|n^2P_{1/2\ m}g\rangle$ to $|n^2P_{3/2\ m'}g\rangle$ matrix elements are the same ones that determine collisional transfer between the $j=\frac{1}{2}$ and $\frac{3}{2}$ states of the alkali doublet. In the transfer problem, the size of the Fourier components of these matrix elements (at frequencies $\sim\Delta E_J/\hbar$) determines the cross section, whereas in (7) the magnitude of the matrix element is more important. *Thus these* $J=\frac{1}{2}$ *depolarization measurements impose independent requirements on any theoretical solutions of the collisional transfer problem.* Since the ${}^2P_{1/2}$ depolarization is strongly dependent upon the magnitude of this matrix element during the close parts of the collision, it is essential that the straight-line path $V(t)$ should be replaced by a $V(t)$ with magnitude and frequency components limited by the repulsive potential. As a consequence, no order-of-magnitude estimates of the cross section will be offered, but it is clear from (7) that it should decrease rapidly as ΔE_J increases beyond $\hbar/\tau_c$.

IV. MEASURED LIFETIMES

In the course of the investigations, the natural lifetimes of the Rb $5p\ {}^2P_{1/2}$ and ${}^2P_{3/2}$ states and the Cs $6p\ {}^2P_{1/2}$ state were obtained from the $n=0$ Hanle-effect linewidths. The limit of zero optical depth was obtained, but no attempt was made to check lamp profile or other uncertainties, so the results have 5 to 10% uncertainty. Nonetheless, our results, $\tau(\mathrm{Rb}\ 5\ {}^2P_{1/2})=3.0\times10^{-8}$ sec and $\tau(\mathrm{Rb}\ 5\ {}^2P_{3/2})=2.7\times10^{-8}$ sec, are in excellent agreement with the more thorough Hanle-effect measurements of Feichtner[10]: $\tau(\mathrm{Rb}\ 5\ {}^2P_{1/2})=3.0\pm0.3\times10^{-8}$ sec and $\tau(\mathrm{Rb}\ 5\ {}^2P_{3/2})=2.8\pm0.2\times10^{-8}$ sec. Also, these lifetimes and our Cs result, $\tau(\mathrm{Cs}\ 6\ {}^2P_{1/2})=3.4\times10^{-8}$ sec, are in excellent agreement with Link.[11]

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[10] John D. Feichtner, thesis, University of Colorado, 1965 (unpublished).

[11] John K. Link, J. Opt. Soc. Am. 56, 1195 (1966).