K and T_1, an indication of different mechanisms for K and T_1. This is compatible with the analysis of Silbernagel *et al.*[27] showing that T_1 is governed by electron-spin correlation effects, while K is caused by conduction electrons (first-order effects) as well as the localized moments via the conduction electrons (second-order effects). Also, the fact that the T_1 values in the $LnAl_2$ compounds (as well as in β-UH_3) are in the msec range while in UP they are one order of magnitude shorter is suggestive that pair-correlation effects are not as effective in UP. In addition it should be pointed out that the validity range for the model proposed in Ref. 27 is for temperature higher than four or five times T_C or T_N, whereas the measurements reported here for UP are at $T\sim 2T_N$.

Measurements on UP are now being extended to the UP-US solid solutions, $UP_{1-x}S_x$. For $UP_{0.95}S_{0.05}$ and $UP_{0.90}S_{0.10}$ indications are that the generalized Korringa formula [Eq. (8)] describes the relation between K and T_1, with slightly increased values of the constant p. The complete results of the study on ^{31}P in the paramagnetic state of the $UP_{1-x}S_x$ system will be published elsewhere.

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Some Properties of a Zeeman Laser with Anisotropic Mirrors

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A laser subject to an axial magnetic field and having one end mirror which exhibits x-y-type loss anisotropy of an arbitrary degree is described from two points of view: (1) the resonance condition for a complete round-trip pass, and (2) the self-consistent-field equations with distributed loss. The behavior of the centrally tuned frequency-locked modes of linear polarization is described in detail as an illustration of the fact that these two theoretical methods are in general not equivalent. Except in the limiting case of weak anisotropies, where the differences become negligible, a correct description must be based on the resonance condition.

THE properties of lasers which are subject to external magnetic fields and whose end mirrors exhibit intrinsic anisotropies have been dealt with in varying degrees of thoroughness by a number of authors. There are two approaches which have been followed. The first is based on the condition of resonance for a complete round-trip pass; i.e., a wave which is reflected from one mirror, propagates, is reflected from the second mirror, and finally propagates back to its original position, is required to reproduce itself exactly. This treatment takes account directly of the properties of the mirrors and was discussed by de Lang.[1] Other authors[2–11] have treated asymmetries by generalizing the Lamb self-consistent field equations.[12] In the scalar theory of Lamb the losses due to the mirrors are treated as if they were distributed in a continuous fashion, becoming in effect a property of the medium. (The motivation for using such a procedure is that one thereby avoids the boundary value problem.) The scalar theory may be generalized by writing analogous equations for each of the electromagnetic vector field components. A generalization which is often made is to assign different Q values to x and y polarizations. Phase properties, including anisotropies, may be introduced by allowing Q to be complex.

Our purpose here is to point out that in general these two formalisms are not equivalent, although they may be nearly so if the anisotropies are weak. However, for strong anisotropies the descriptions are quite different. With the first method, in which a wave must satisfy the condition of self-reproducibility on a complete round-trip basis, there will nevertheless be a significant change in its state of polarization upon reflection from a strongly anisotropic mirror. The original state of polarization is then restored through a compensating

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change, also large, which must be effected either at the other mirror or through propagation. On the other hand, if the properties of the mirrors (losses, phase shifts, and anisotropies in either or both) have been described as properties of the medium, then the eigenstates of the self-consistent field equations satisfy the condition of self-reproducibility on a differential basis, that is, the eigenstate is preserved as it propagates from z to $z+\Delta z$, for any z, and an arbitrarily small Δz. So this description is a good approximation only when no major changes are experienced in reflection (and propagation), i.e., when the anisotropies are weak.

By way of illustration of the above discussion, and also because the results are of some intrinsic interest themselves, we shall discuss the case of a laser subject to a constant axial magnetic field and having one end mirror which displays x-y-type loss anisotropy, the other mirror being assumed isotropic. This problem has been discussed in the literature[2,3,6,9,10,13,14] for the case of small, residual-type anisotropy, because assumption of this configuration predicts the kind of behavior which has been observed[14–16] in small magnetic fields, namely, that for central tuning a single frequency-locked mode of linear polarization is observed, the direction of polarization rotating through an angle of about 45° as the field is increased from zero to some critical value H_c, beyond which the single-mode regime ceases to exist. We shall be concerned only with the single-mode operation. It will be shown that if the anisotropies are not small, the two formalisms outlined above lead to very different predictions. However, either one is adequate for the limiting case of weak anisotropy.

Let us first consider the problem from the point of view of the resonance condition for a complete round-trip pass. Since one mirror is being assumed isotropic we shall lump its reflection coefficients in with those for the anisotropic mirror. We assume that in the presence of an axial magnetic field the propagation is described by $\exp(-j\beta_\pm z)$ for right and left circular components. The condition for resonance is (with $\lambda=1$)

$$P^2R\mathbf{E}=\lambda\mathbf{E}, \tag{1}$$

where P^2 describes propagation back and forth through a cavity of length L, and R is the reflection matrix for the anisotropic mirror. In the circular basis

$$P^2=\begin{pmatrix} p_+ & 0 \\ 0 & p_- \end{pmatrix}, \tag{2}$$

where

$$p_\pm=e^{2\beta_\pm''L}e^{-2j\beta_\pm'L},$$
$$\beta_\pm'=\beta_0[1+\tfrac{1}{2}\chi_\pm'], \tag{3}$$
$$\beta_\pm''=\tfrac{1}{2}\beta_0\chi_\pm''-\alpha_0.$$

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Here $\beta_0=\omega/c$ for a wave of frequency ω; $\chi_\pm'$ and $\chi_\pm''$ are the real and imaginary parts of the susceptibilities for right and left circular polarizations; and α_0 is a phenomenological loss term (representing diffraction losses, etc.). The reflection matrix is diagonal in the Cartesian basis:

$$R'=\begin{pmatrix} r_x & 0 \\ 0 & r_y \end{pmatrix} \tag{4}$$

(here r_x and r_y are assumed to be real). When transformed to the circular basis it becomes

$$R=\begin{pmatrix} a & -b \\ -b & a \end{pmatrix}, \tag{5}$$

where

$$a\equiv\tfrac{1}{2}(r_x+r_y), \quad b\equiv\tfrac{1}{2}(r_x-r_y). \tag{6}$$

The transformation is $R=TR'T^{-1}$, where

$$T=2^{-1/2}\begin{pmatrix} -1 & -j \\ 1 & -j \end{pmatrix}. \tag{7}$$

Note that the eigenvectors of Eq. (1) represent the electric field vector just before it is reflected.

We solve the eigenvalue problem of Eq. (1) working in the circular basis. Setting the eigenvalues λ equal to

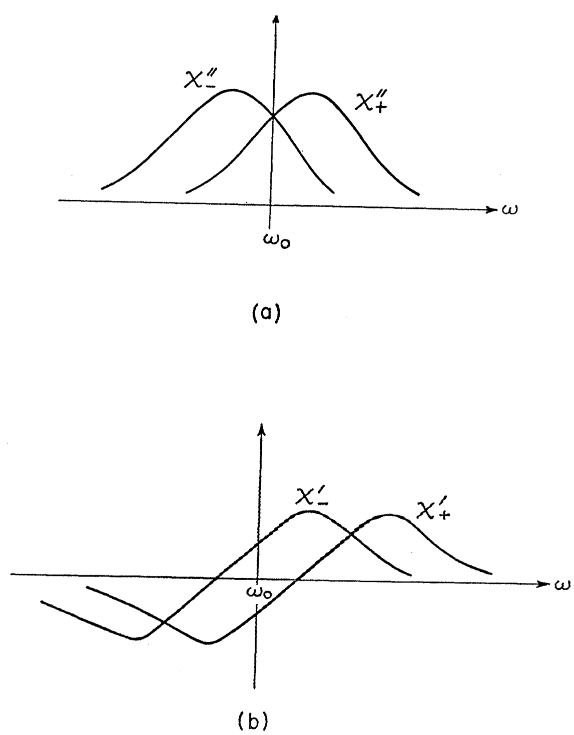

FIG. 1. The right and left circular components of the susceptibility $\chi=\chi'+j\chi''$ for the case of equal g factors and an axial magnetic field. The abscissa is the frequency of oscillation, and there is one set of these curves for each value of the mode intensity (Ref. 17). In the absence of magnetic field the function $\chi''(\omega)$ is symmetric about the natural atomic resonant frequency ω_0.

unity, we obtain the oscillation condition for each of the two possible modes:

$$\frac{1}{\epsilon a\tau}=\frac{\sigma+1/\sigma}{2\epsilon}\pm\left[\frac{(\sigma-1/\sigma)^2}{4\epsilon^2}+1\right]^{1/2}, \tag{8}$$

where

$$\sigma^2\equiv p_+/p_- \quad \tau^2\equiv p_+p_- \tag{9}$$

and ϵ is the asymmetry parameter

$$\epsilon\equiv b/a=(r_x-r_y)/(r_x+r_y). \tag{10}$$

The eigenvectors are

$$\frac{E_+}{E_-}=\sigma\left\{\frac{-(\sigma-1/\sigma)}{2\epsilon}\mp\left[\frac{(\sigma-1/\sigma)^2}{4\epsilon^2}+1\right]^{1/2}\right\}. \tag{11}$$

The oscillation condition (8) is really two conditions, one for the phases (which determines the frequency of oscillation) and another for the magnitudes (which determines the intensity). If the levels involved have equal g factors, the right and left circular components of the susceptibility are as shown in Fig. 1, where the amount of splitting is proportional to the magnetic field.[17] We consider the phase part of Eq. (8), and we shall show that if $\omega=\omega_0$ (central tuning) is the solution when there is no magnetic field, then this solution persists for a range of fields $0\leqslant H\leqslant H_c$, where H_c is some critical field. We need only to assume that $\chi_\pm'(\omega_0)$ and $\chi_\pm''(\omega_0)$ obey the following symmetry relations:

$$\chi_+''(\omega_0)=\chi_-''(\omega_0),$$
$$\chi_+'(\omega_0)=-\chi_-'(\omega_0). \tag{12}$$

It is not necessary to make any further assertions about the dependence of χ on either the intensity E^2 or the frequency ω.

To verify that $\omega=\omega_0$ does indeed allow the phase condition to be satisfied, we first observe that with the relations (12) σ may be written

$$\sigma=e^{j\phi/2}, \tag{13}$$

where

$$\tfrac{1}{2}\phi\equiv\chi_-'(\omega_0)\beta_0\mathrm{L}. \tag{14}$$

Now we note that the phase of the left side of Eq. (8) (i.e., the phase of τ^{-1}) is simply $2\beta_0\mathrm{L}$, independent of magnetic field, whereas the right side is real provided that

$$\sin^2(\tfrac{1}{2}\phi)\leqslant\epsilon^2. \tag{15}$$

The phase condition then becomes

$$2\beta_0\mathrm{L}=2\pi q, \tag{16}$$

where q is an integer (the mode number). Since the cavity resonant frequency ω_c must satisfy $2\omega_cL/c=2\pi q$, we see that if the cavity has been tuned to the center of the unshifted Doppler line ($\omega_c=\omega_0$), the solution $\omega=\omega_0$ does then satisfy the phase equation. This solution is valid for a range of magnetic fields $0\leqslant H\leqslant H_c$, corresponding to $0\leqslant\phi\leqslant\phi_c$, where[18]

$$\sin^2(\tfrac{1}{2}\phi_c)=\epsilon^2, \tag{17}$$

since the phase condition remains unaffected for values of ϕ in this range. The solution $\omega=\omega_0$ is, of course, valid for both of the modes indicated ($\pm$) in Eq. (8). Although, in general, it would be expected that a magnetic field would split the mode into two frequency components, in this case the modes are degenerate in frequency. If $H>H_c$ two frequency operation becomes possible,[18] but we restrict our considerations here to the single frequency regime.

Although the two modes are degenerate in frequency, they are not degenerate in polarization. From Eqs. (11) and (13) we find that

$$E_+/E_-=\mp e^{j(\phi\pm\kappa)/2}, \tag{18}$$

where [cf. Eq. (15)]

$$\sin(\tfrac{1}{2}\kappa)\equiv\sin(\tfrac{1}{2}\phi)/\epsilon. \tag{19}$$

Equation (18) says that the modes are linearly polarized. In zero magnetic field ($\phi=0$)$E_+/E_-=\mp1$, so the upper (lower) sign refers to the mode which in the absence of magnetic field is polarized in the x direction (y direction). We shall refer to this as the X mode (Y mode). Generally, only one mode oscillates, the other mode being either below threshold or quenched by mode competition effects (see below). From Eq. (18) we see that in the presence of magnetic field the direction of linear polarization has been rotated from its zero-field value through an angle θ, where

$$\theta=\tfrac{1}{4}(\phi\pm\kappa). \tag{20}$$

The maximum rotation angle $\theta=\theta_c$ occurs for $\phi=\phi_c$, whence (assuming for the sake of definiteness that $\phi>0$ and $\epsilon>0$, i.e., $r_x>r_y$)

$$\theta_c=\tfrac{1}{4}(\phi_c\pm\pi). \tag{21}$$

[17] If the g factors are not equal, the right and left circular components of susceptibility each contain more than one frequency, so $\chi_\pm'$ and $\chi_\pm''$ become a superposition of the forms shown in Fig. 1. However, even in this case we shall assume that the symmetry relations (12) still hold. In Fig. 1 the susceptibility is to be considered as the net response of all atoms, i.e., an integral over the complete Doppler distribution. The susceptibility is a function both of intensity (through which it shows the effects of saturation) and frequency. Thus there are really many sets of curves of the type shown there, one for each value of the mode intensity. The particular value of intensity which actually results is such as will allow the magnitude part of the oscillation condition [Eq. (22)] to be satisfied.

[18] The exact dependence of ϕ upon H may be quite complicated, although it is suggested qualitatively in Fig. 1(b). It can be seen that in general there will be a second critical field $H_c'>H_c$ such that $\phi(H_c')=\phi(H_c)$, so that for $H>H_c'$ the condition (15) will again be satisfied, i.e., there will be a second region for which central tuning is accompanied by frequency locking and linear polarization. If the maximum value of $\chi_-'(\omega)$ is such that $\sin^2(\frac{1}{2}\phi)<\epsilon^2$ for all magnetic fields, then the two frequency operation would never be observed at $\omega=\omega_0$.

Thus, as H is varied from zero to H_c, the X mode is rotated by $\frac{1}{4}\pi+\frac{1}{4}\phi_c$, while the (nonoscillating) Y mode is rotated through the complementary angle in the opposite direction, so at $H=H_c$ the two are collinear.[19] The maximum rotation angle for the X mode is slightly more than 45° for weak anisotropy ($\epsilon\ll 1$), whereas it approaches 90° in the limit of strong anisotropy ($\epsilon\to 1$).

The magnitude part of the oscillation condition (8) may be written

$$\exp[\beta_0 L\chi_+''(\omega_0)-2\alpha_0 L] =[a\cos(\tfrac{1}{2}\phi)\pm b\cos(\tfrac{1}{2}\kappa)]^{-1}, \quad (22)$$

where the left side represents the net gain experienced through propagation and the right side represents the loss at the mirrors. It is easily verified that in zero magnetic field the mirror loss is r_x^{-1} for the X mode and r_y^{-1} for the Y mode. When $H=H_c$ the mirror loss is the same for both modes and is equal to the geometric mean of the zero-field losses, namely $(r_x r_y)^{-1/2}$. The variation of the mirror loss with magnetic field is clearly a geometric effect which is associated with the rotation of the direction of linear polarization (as may be verified directly through calculation of the reflection coefficient for a wave incident at an angle θ). We note that at the critical field H_c there is complete degeneracy between the two modes, in frequency, polarization, gain, and loss.

If the above results are taken in the limit of $\epsilon\ll 1$ we find that (1) the critical field H_c corresponds to a value ϕ_c, where

$$\lim_{\epsilon\ll 1}(\tfrac{1}{2}\phi_c)^2=\epsilon^2 \quad (23)$$

and (2) the rotation angle for the direction of polarization is given by

$$\lim_{\epsilon\ll 1}\theta=\pm\tfrac{1}{4}\kappa, \quad (24a)$$

which can also be written

$$\lim_{\epsilon\ll 1}\sin 2\theta=\pm\phi/2\epsilon. \quad (24b)$$

At $H=H_c$ we have $\theta=\theta_c$, where

$$\lim_{\epsilon\ll 1}\theta_c=\pm\tfrac{1}{4}\pi. \quad (25)$$

These results for weak anisotropy are in agreement with observation[14–16] and with previous calculations.[2,3,6,13,14]

A physical interpretation has been given by D'Yakonov and Fridrikhov.[13] They observe that in the presence of a magnetic field a linearly polarized wave experiences Faraday rotation of an amount $\frac{1}{2}\phi$. At the same time if this wave is incident at an angle θ upon mirrors with unequal reflection coefficients, the reflected wave experiences a rotation θ_r (the authors assume $\epsilon\ll 1$):

$$\lim_{\epsilon\ll 1}\theta_r=\epsilon\sin 2\theta. \quad (26)$$

If we equate the Faraday rotation to the reflection rotation we then obtain Eq. (24b), and also the corollary results (23) and (25), as shown in Ref. 13 [see especially Eq. (5.7)]. This same argument is, of course, valid for asymmetries which are not small. In the general case, θ_r is given by

$$\tan\theta_r=\tan\theta\frac{2\epsilon}{1+\tan^2\theta+\epsilon(1-\tan^2\theta)}. \quad (27)$$

The requirement that the angle θ_r be equal to the Faraday rotation $\frac{1}{2}\phi$ leads to the same expression for θ that was given above in Eq. (20).

If this problem is approached from the point of view of the self-consistent field equations with distributed loss, then different Q values are assigned for the x and y directions ($Q_x\neq Q_y$). These equations have been handled in simplest fashion by D'Yakonov[2] for an arbitrary degree of anisotropy. The asymmetry parameter ϵ is expressed in terms of Q_x and Q_y through

$$\epsilon=\tfrac{1}{2}\beta_0 L(1/Q_y-1/Q_x). \quad (28)$$

It is found[2,20] that for central tuning there exist frequency locked modes which are linearly polarized if the magnetic field lies in the range $0\leqslant|H|\leqslant|H_c|$, corresponding to $0\leqslant|\phi|\leqslant|\phi_c|$, where (in the present notation)

$$(\tfrac{1}{2}\phi_c)^2=\epsilon^2. \quad (29)$$

The direction of polarization is rotated from its zero-field value through an angle θ given by

$$\sin 2\theta=\pm\phi/2\epsilon, \quad (30)$$

the maximum value θ_c occurring when $H=H_c$:

$$\theta_c=\pm\tfrac{1}{4}\pi. \quad (31)$$

These results are to be compared to those [Eqs. (17), (20), and (21)] which were obtained from the resonance condition (1). It is seen that, in general, the two sets of results do not agree except in the limit of small anisotropy [$\epsilon\ll 1$; cf. Eqs. (23)–(25)]. An especially significant difference concerns the amount of rotation which has been achieved when $H=H_c$. In the self-consistent field calculation this is always $\frac{1}{4}\pi$, but in the calculation based on the resonance condition this maximum rotation could be anywhere between $\frac{1}{4}\pi$ and $\frac{1}{2}\pi$, depending on the value of ϵ [see Eq. (21)].

Thus, we have illustrated through a specific example the remarks made earlier concerning the nonequivalence

[19] In the second region of linear polarization (see Ref. 18) the rotation angle θ would start at $\theta=\theta_c$ for $H=H_c'$ and then decrease toward zero as H was further increased. If $\sin^2(\frac{1}{2}\phi)<\epsilon^2$ for all magnetic fields the maximum amount of rotation would be determined by the maximum value of $\chi_-'(\omega)$.

[20] It may be verified through direct calculation that the results using the self-consistent field equations are also obtainable by equating the differential Faraday rotation to the differential rotation induced when a wave propagates with unequal decay constants [$\alpha_x\neq\alpha_y$, $\alpha=\beta_0/2Q$, decay $\sim\exp(-\alpha z)$], considering only the propagation from z to $z+dz$.

of the two methods. The solution to the self-consistent field equations was an eigenstate characterized by a particular direction θ [Eq. (30)]. However, if the mirror is highly anisotropic the wave suffers a significant amount of rotation upon reflection [cf. Eq. (27)], so a single direction of polarization is, in this case, not an adequate kind of description. If the anisotropy is weak, the reflection rotation is small, and it then becomes a good approximation to say that the wave is always in one particular direction.

Lasers having mirrors which show a weak loss anisotropy have been studied experimentally,[14–16] but no studies have been reported for the case of a substantial loss anisotropy. If observations were to be made for this latter case and compared to the results presented here, it should be remembered that the wave which is described in Eq. (1) is the wave incident upon the anisotropic mirror. To find the transmitted wave one must take account of the transmission matrix for the anisotropic mirror (assuming, as is most likely, that the output is taken there).

In summary, then, we have considered in detail a centrally tuned laser subject to an axial magnetic field and having one end mirror which shows an x-y-type of loss anisotropy. It was found that the behavior of the frequency-locked modes of linear polarization is to a large extent independent of the complex details of the nonlinear interactions, with the result that the mathematics required to describe these modes is relatively simple.[21] Thus, a consideration of these modes is especially appropriate to the purpose of illustrating through a concrete example the differences between two theoretical bases, namely, (a) application of the resonance condition for a round-trip pass, and (b) the use of the self-consistent field equations with distributed loss. In general, the former method (a) should be used, although the differences between the two become negligible in the limit of small asymmetries. In the example considered here, the theory predicts, in the case of mirrors with a strong loss anisotropy, some new features which have yet to be observed.

[21] This mathematical simplicity, it should be noted, derives from the fact that the symmetry relations (12) allow for a rather full treatment of the eigenvalue problem; in particular, the solutions for the frequency of oscillation and the eigenstate of polarization do not depend on the details of the nonlinearities, which are contained in the susceptibility χ. Thus for the frequency-locked modes of linear polarization the simplicity of their physical behavior is reflected in the mathematics. The situation would, of course, not be so simple for other configurations (e.g., two frequency operation, detuning of the cavity from line center, the inclusion of phase anisotropies).

PHYSICAL REVIEW VOLUME 178, NUMBER 2 10 FEBRUARY 1969

Spin-Lattice Relaxation of *F* Centers in Alkali Halides: Theory and Optical Measurements to 50 kG*

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We have measured the ground-state spin-lattice relaxation of F centers in KBr and KI over the range of magnetic field from 0 to 50 kG, and at the temperature of 1.6°K. Essentially continuous measurement over this large field range was made possible by a detection technique in which the magnetic circular dichroism of the optical absorption band was monitored. In the range 10 to 50 kG, the relaxation rate is predominantly that of single F centers. Here, the field dependence of the relaxation rate followed very closely a curve of the form $(AH^3+BH^5)\coth(g\beta H/2KT)$. The first term represents a mechanism involving phonon modulation of the hyperfine contact interaction with neighboring nuclei, and is the larger term for $H_0 \lesssim 25$ kG. The second term is due to relaxation via the Kronig–Van Vleck process. A theoretical evaluation of that part of the rate due to the hyperfine mechanism has produced close agreement with experimental results. In the evaluation, the interaction with the second shell of nuclei (halogens) was found to be of greatest importance for KBr and KI.

I. INTRODUCTION

At very low temperatures, where the Raman process is frozen out, spin-lattice relaxation of F centers in alkali halides takes place by the direct process.[1] In its ground state, the F-center electron may be approximately described by an s-type wave function which spreads out over many lattice sites. At all but the highest fields, a relaxation process by way of modulation of the hyperfine interaction may predominate over the common mechanism of modulation of the crystal field, which for an s-type electron should be very small. From measurements of field dependence

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† Visiting Research Fellow of the Consejo Nacional de Investigaciones Cientificas y Tecnicas, Argentina.
[1] For a good collection of papers, see *Spin-Lattice Relaxation in Ionic Solids*, edited by A. A. Manenkov and R. Orbach (Harper and Row, New York, 1966).