Lorentzian curve (laid on its side since the frequency deviation is here the ordinate).

For simplicity, we now confine the discussion to the case $\eta_0=0$ (spin wave scattering only). As y increases, the Lorentzian shape is distorted so that more of the area is in the lower half plane. When y reaches 10, the distortion begins to take the form of a foldover (see Fig. 1) so that over a certain range of z the response is triple valued. Figure 2 shows the region in the z-y plane in which this occurs, so that bistability may be expected there.

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[6] This formula was derived independently by P. Gottlieb of Hughes Aircraft Company.

Terraces on Etchpits

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It is often observed that the side faces of etchpits do not have smooth surfaces, but are terraced. It is the purpose of this note to present evidence for a simple cause of such terraces that may act in many cases. No doubt there are several other possible causes, but they will not be considered here.

For some time it has been known that "aged"[1] dislocations in LiF crystals tend to produce etchpits with terraced sides, whereas "fresh" dislocations do not.[1] It has been suspected that segregated impurities that form nonuniform threads along the aged dislocations are responsible for the phenomenon, but there has been no proof. Direct evidence that discontinuous precipitates can form terraced etchpits is presented here.

Figure 1(a) shows an etchpit with very high terraces along with some other pits with smaller terraces. At the center of the pit with large terraces a string of small dots may be seen. These dots are caused by small particles which lie on the surface of the crystal. As the etching proceeded, each time a new terrace was observed, a new particle was also observed at the center of the pit. Figure 1(b) shows the etchpit of Fig. 1(a) after it had been etched additionally. The photograph was made just after a new terrace with its associated precipitate particle appeared. During the 4 min etching interval between the two pictures, the number of particles increased from 7 to 10; but the total number of terraces did not increase by the same number, because the terraces are not stable. They gradually become rounded off and disappear. The particles are 2–4000 A in diam, and the terraces have approximately this height when they first form.

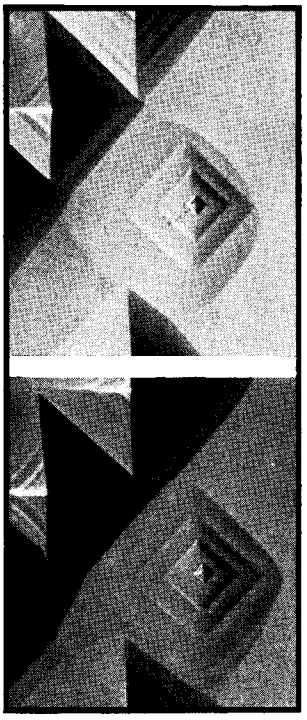

FIG. 1. Deeply terraced etchpit with an associated stringer of precipitate particles. The LiF crystal was etched in Etch A in 30 sec increments (Etch A = 1 pt HF, 1 pt glacial HA_c, and 2×10^{-6} pt FeF_3). (a) 7 min etch; (b) 11 min etch. 1000X.

The terraces in the case of Fig. 1 were caused by rapid development of the etchpit at a precipitate particle followed by slow etching between particles. In other cases, the rapid and slow stages might be reversed, because of different chemical factors, but the resulting terracing would look the same.

In Fig. 1, the precipitate particles are not located along a dislocation line. This may be deduced because the etchant (Etch A) is known to attack dislocations at a more rapid rate than was observed here. However, when dislocations *are* present together with precipitates, two sequences seem to operate which lead to terraced pits. In the first sequence, the rate of attack is most rapid at the parts of a dislocation that have precipitates, and is slower between precipitates where only the dislocation line is present. In the second sequence, the relative rates are reversed.

[1] "Aged" means dislocations that have been slowly cooled from the melting temperature or from a temperature above about 250°C. "Fresh" means dislocations introduced at room temperature and not allowed to stand more than several weeks.

Effect of Nuclear Polarization on the Behavior of Solid State Masers*

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(Received January 15, 1960)

Cowen, Schafer, and Spence[1] have shown that the saturation of the electron spin transitions of the Cr^{+++} ions in ruby produces significant polarization of the Al_{27} nuclei. Other investigators[2,3] have detected a similar effect in a variety of paramagnetic materials. Recently, we have observed what amounts to the inverse of this effect, i.e., that the application of rf power at or near frequency of the quadrupole resonance of the aluminum nuclei in ruby produced a marked change in the electron-spin resonant absorption under saturation conditions.[4] This communication deals with the effect of this interaction on the behavior of solid-state maser.

In the basic experiment, a two-turn coil was wound about a ruby sample containing nominally 0.1% chromium. This assembly was placed in a doubly resonant microwave cavity and located in the dc magnetic field in such a manner that the axis of the coil was perpendicular to the direction of the field, and the polar angle was approximately 60°. With the system cooled to 4.2°K, a K-band microwave pump was used to saturate the 1–3 transition, and stimulated emission at X-band frequency was obtained in the 3–2 transition. Subsequent application of rf power at 4.5 Mc produced an increase in the gain of the maser amplifier, and resulted in a change of mode of operation of the maser oscillator. The effect was most pronounced at the resonant frequency of the free aluminum nuclei; however, at higher levels of rf power substantial interaction was obtained over a band extending from 500 kc to 20 Mc. Typically, power levels required to produce a detectable effect were of the order of 10 mw on resonance, and about an order of magnitude higher off resonance.

Comparing the performance of the maser amplifier with rf power on and off, as shown in Figs. 1(a) and 1(b), respectively,

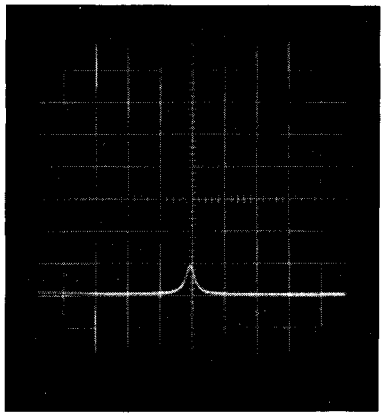
we may calculate the change in the magnetic Q due to the change in polarization of aluminum nuclei. The dependence of gain of the maser amplifier on the magnetic Q is given by

$$G=\left(\frac{1}{Q_c}-\frac{1}{Q_L}+\frac{1}{|Q_m|}\right)^2\Big/\left(\frac{1}{Q_c}+\frac{1}{Q_L}-\frac{1}{|Q_m|}\right)^2,$$

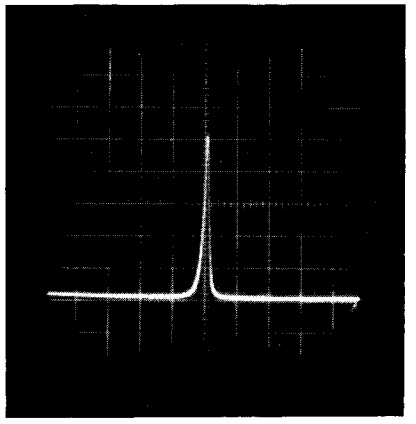
where Q_c is the coupling Q, Q_L is the loss Q, and Q_m is the negative magnetic Q. In our setup, Q_c and Q_L are typically 10^3. By using this value, we compute the decrease in Q_m caused by the application of rf to be approximately 20%. This change in magnetic Q was observed under conditions of partial saturation of the pumping transition.

The effect of decreased nuclear polarization on the behavior of the maser oscillator is shown in Fig. 2. The application of rf power tended to make the transient of the cw mode of the oscillator [Fig. 2(a)] less damped, and led eventually to the relaxation mode of operation [Fig. 2(b)]. The use of high rf power levels tended to diminish, or even reverse the effect.

It should be noted that the transient form of the interaction can be observed without the use of rf power. When maser action is initiated by bringing the dc magnetic field rapidly to the appropriate value, there results a transient during which the amplifier gain rises quickly to a high value, then decays slowly to a lower steady-state value. The time constant of the transient is of the order of seconds. The existence of the transient cannot be ascribed to spin-lattice relaxation, since T_1 for cooled ruby is of the order of 0.1 sec. Rather, it appears to result from the relaxation of polarization of aluminum nuclei.

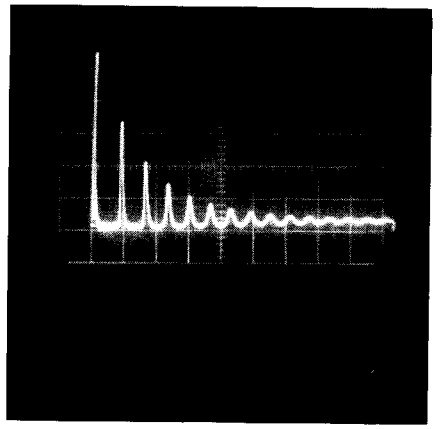


(a)

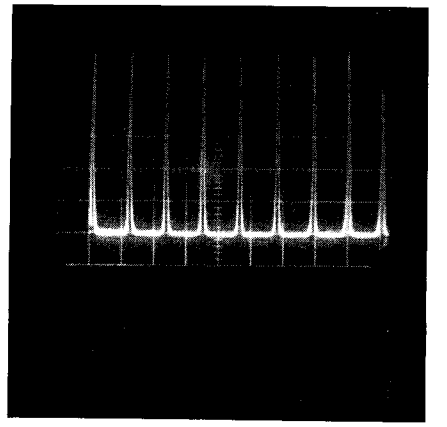


(b)

FIG. 1. Effect of change of nuclear polarization on maser amplifier gain. (a) rf off; power gain—15 db; (b) rf on; power gain—35 db.



(a)



(b)

FIG. 2. Effect of change of nuclear polarization on the behavior of maser oscillator. (a) rf off; cw mode of operation; (b) rf on; relaxation mode of operation.

Subsequent experiments have indicated that the application of rf power does not affect thermal relaxation processes to any perceptible degree. Rather, it can be thought of as added pumping, resulting in an increase in the magnetic Q. The degree of enhancement of the punping process is of the order of change in the degree of nuclear polarization, estimated to be about 1% under our experimental conditions.[4] Under conditions of marginal saturation of the pumping transition, this may lead to substantial decrease in magnetic Q, as noted in the foregoing.

These considerations suggest that the recovery of a solid-state maser device from saturation, or from fluctuations in pumping power, is characterized in general by two time constants. There is rapid recovery, of the order of 0.1 sec, determined by spin-lattice relaxation; and slow recovery, of the order of seconds, determined by nuclear-electronic relaxation. These effects have been observed experimentally. The latter effect, of course, becomes significant only when the on-period of the saturating signal, or variations in the degree of saturation of the pumping transition, are on a time scale comparable to that of nuclear relaxation.

In conclusion, we would like to conjecture that the observed increase in gain of the maser amplifier resulting from the application of rf power may be used to advantage in detecting weak nuclear resonances as well as in studies of ENDOR-type effects. In essence, the effect on the electron system produced by resonant pumping of the nuclear system, is amplified by the practically noiseless maser amplifier. When a high degree of regeneration is used, a small change in the magnetic Q results in a very con-

siderable change in gain, or output, as attested by Figs. 1(a) and 1(b), and the corresponding calculation. A disadvantage of this method is, when using ruby, that maser action cannot be obtained in the vicinity of $\theta=0°$, whereas nuclear resonance effects are conveniently observed at that orientation. Presumably, this can be remedied by using a paramagnetic material, such as iron-doped sapphire, which permits maser action at the above orientation.

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† Now with the Scientific Laboratories, Ford Motor Company, Dearborn, Michigan.
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Microwave Measurement of Semiconductor Carrier Lifetimes

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IN a recent communication, Ramsa, Jacobs, and Brand[1] have described a microwave technique for the measurement of carrier lifetime in semiconductors. In this technique, a pulse of minority carriers is injected into a slab of semiconductor placed transversely across a waveguide. The carrier lifetime is deduced from observation of the resulting transient variation of microwave attenuation by the sample. The authors find, from an analysis of the attenuation time decay curve shape, that the relative carrier injection level must be kept small in order that the increment of attenuation shall decay exponentially with the same time constant as the surplus carrier density. It may also be shown that the conditions under which this attenuation decay has the same time constant as the surplus carrier decay may be derived from an analysis wave propagation problem. We treat here the propagation of dominant mode waves in a waveguide containing an imperfect dielectric medium in which the conductivity is a function of time.

On assuming charge neutrality for the medium, as must be the case macroscopically, for example, with optical injection of carriers, the electric field must satisfy a wave equation of the form

$$\nabla^2 E_y=\mu\sigma(\partial E_y/\partial t)+\mu\epsilon(\partial^2 E_y/\partial t^2). \tag{1}$$

Here, σ is the conductivity and μ and ϵ the magnetic permeability and dielectric constant of the medium, respectively. The dominant mode solution of Eq. (1) is

$$E_y=\sin(\pi x/b)e^{-\gamma z}e^{j\omega t}, \tag{2}$$

in which b is the waveguide width, ω is the angular frequency of the waves, and γ is a complex propagation constant defined as

$$\gamma=\alpha+j\beta. \tag{3}$$

From Eqs. (1) and (2) the propagation constant is

$$\gamma=\pm[(\pi/b)^2-\mu\epsilon\omega^2+j\mu\sigma\omega]^{\frac{1}{2}}. \tag{4}$$

If the waveguide is operating far from its cutoff frequency, $(\pi/b)^2$ in Eq. (4) may be neglected in comparison with $\mu\epsilon\omega^2$. If it is further assumed that $(\sigma/\omega\epsilon)<1$, and expanding to second order in $(\sigma/\omega\epsilon)$, Eq. (4) becomes approximately

$$\gamma\approx(\sigma/2)(\mu/\epsilon)^{\frac{1}{2}}+j\omega(\mu\epsilon)^{\frac{1}{2}}[1+\sigma^2/(8\omega^2\epsilon^2)]. \tag{5}$$

If the second term in the square brackets of Eq. (5) is not negligible with respect to unity, the wave propagation constant β will depend on σ. The resulting change in the impedance of the sample may cause a change in reflected power, and hence an incorrect deduction of transmitted power. It may also be shown that a similar result can occur if $(\sigma/\omega\epsilon)>1$. The condition assumed to hold in the reference of footnote 1, namely that only the absorbed power varies with the excess conductivity, requires that $(\sigma/\omega\epsilon)\ll 1$. By using $\epsilon=16\epsilon_0$, for germanium, this condition would limit the usable conductivity, for X-band measurements, to about 10 (ohm-m)$^{-1}$, or for K-band operation, to about 30 (ohm-m)$^{-1}$. If the restriction of small $(\sigma/\omega\epsilon)$ is retained, the behavior of the microwave power attenuation by the sample may be deduced as follows. The power propagated along the waveguide is

$$P=\frac{1}{2}\int\mathbf{E}\times\mathbf{H}^*\cdot d\mathbf{S}=j[(ab)/(4\mu\omega)]\gamma^* e^{-2\alpha z}, \tag{6}$$

where a is the height of the waveguide, and the factor j is a phase factor. On using Eq. (3), Eq. (6) is of the form,

$$P=A\alpha e^{-2\alpha z}+Be^{-2\alpha z}, \tag{7}$$

in which A and B are constants not dependent on time.

It is known that the decay of a surplus minority carrier concentration in a semiconductor can be described under certain restrictive conditions of low surplus carrier concentration in terms of a lifetime τ associated with an exponential time decay.[2-4] For very large surplus carrier densities the decay is no longer of exponential form. However, the concept of carrier lifetime is sufficiently useful to justify its retention as an approximate description of the surplus carrier density decay. Consequently, the conductivity-dependent attenuation constant α is assumed to have the time dependence,

$$\alpha=\alpha_0[1+c\exp(-t/\tau)], \tag{8}$$

in which c is a measure of the initial surplus carrier concentration, and τ is the carrier lifetime. On using Eq. (8), Eq. (7) has the form

$$P=A\alpha_0[1+c\exp(-t/\tau)]\exp\{-2z\alpha_0[1+c\exp(-t/\tau)]\}+B\exp\{-2z\alpha_0[1+c\exp(-t/\tau)]\}. \tag{9}$$

In general, $c<1$ and if the sample thickness z is made very small, Eq. (9) becomes

$$P=A\alpha_0[1+c\exp(-t/\tau)]+B \tag{9a}$$

and the time dependence of the transmitted power is exponential with a time constant equal to the carrier lifetime τ.

However, a very small sample thickness is not practically desirable and it also increases the relative importance of surface recombination, leading to an incorrect result for bulk lifetime. An alternative experimental restriction is that of low initial surplus carrier density; $c\ll 1$. In this approximation, the exponential factors in Eq. (9) may be expanded. On neglecting terms in c^2, the transmitted power then has the form,

$$P=F+cG\exp(-t/\tau), \tag{9b}$$

where F and G are constants. Hence it is seen that with the restriction of low carrier injection level, the transmitted microwave power decays exponentially with time constant τ.

The transmitted power Eq. (9) may be expressed in a third alternative form if with the values of material constants σ and ϵ cited above, α in Eq. (3) is approximately an order of magnitude smaller than β. The logarithm of power in Eq. (6) has the form,

$$\ln P=\ln K+\ln\gamma^*-2\alpha z, \tag{10}$$

where K is a constant. After expanding the $\ln\gamma^*$ term,

$$\ln P=\ln[(K\beta)/j]-\alpha(2z+j/\beta). \tag{11}$$

Therefore, when α/β is small, the logarithm of the transmitted power exhibits a transient variation similar in character to that of the attenuation constant α, and hence to that of the conductivity of the semiconductor medium.

In summary, it has been shown that the time constant of the microwave power attenuation transient curve can furnish a