and Raman processes has been observed previously (*4*).

In order to verify that additional nonlinear effects did not set in at higher powers, we raised the incident power on the dewar in small increments. The measured results continued to be the same up to 100 megawatts, at which point catastrophic failure of the helium dewar occurred.

We believe that the anomalous Brillouin shifts reported by Krivokhizha *et al.* can be attributed to stimulated Raman scattering which becomes quite strong at low temperatures. In fact, simultaneous observation of Raman and Brillouin lines on the same Fabry-Perot interferogram has enabled us to measure Raman shifts and their temperature dependences to an accuracy of 0.1 cm^{-1} (*5*).

P. E. TANNENWALD
J. B. THAXTER
Lincoln Laboratory, Massachusetts Institute of Technology, Lexington

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15 August 1966

Possibility of Maser Action in Cosmic Radio Sources

Abstract. *A mechanism is described for maser action in synchrotron sources. The process may contribute to the low-frequency radiation from 3C273B.*

The question raised here is whether a substantial part of the radiation from a cosmic radio source may be due to induced emission rather than spontaneous emission. Such a situation can arise when a population inversion exists in the energy source of the radiation and when the source is optically thick to its own radiation, as is the case with an ordinary maser.

It is widely believed that the radio power emitted from cosmic radio sources is synchrotron radiation—the classical radiation of relativistic electrons gyrating in a magnetic field (*1–5*). Most of the power is radiated by the electrons at high harmonics of their gyrofrequency, which is given by

$$\nu_g = \frac{\nu_0}{\gamma} = \frac{eH}{2\pi mc\gamma} \tag{1}$$

where γ is the Lorentz factor of the electron. Appreciable power is radiated at harmonics up to a critical frequency given by

$$\nu_c = \frac{3}{2}\,\nu_g\,\gamma^3 \tag{2}$$

It is to be expected that inhomogeneities in the magnetic field over the orbit of the electron will cause the line width of a given harmonic to be much larger than the spacing between harmonics. If so, the power spectrum radiated by a given electron will be a continuum, whose shape is given by (*6*)

$$p(\nu,\gamma) = \frac{\sqrt{3}e^3 H}{mc^2}\frac{\nu}{\nu_0}\int_{\frac{\nu}{\nu_c}}^{\infty} K_{5/3}(x)\mathrm{d}x \approx$$

$$\frac{\sqrt{3}e^3 H}{mc^2}\begin{cases} 2.15\,(\nu/\nu_c)^{\frac{1}{3}},\ \nu/\nu_c \ll 1 & \text{(a)} \\ 1.25(\nu/\nu_c)^{\frac{1}{2}}\,e^{-\nu/\nu_c},\ \nu/\nu_c \gg 1 & \text{(b)} \end{cases} \tag{3}$$

If a radio source is optically thick to its own synchrotron radiation, one has to calculate the induced emission and absorption of synchrotron radiation as well as the spontaneous emission given by Eq. 3. The rate of induced emission and absorption can be calculated from the rate of spontaneous emission by means of Einstein's relations. Ordinarily, absorption exceeds induced emission, so that a radiator which is optically thick can emit radiation only from a surface layer. But in certain situations, such as a maser, the induced emission can exceed the absorption. In this case, the combination of the two processes is called negative absorption.

Is it possible for a source of synchrotron radiation to display this phenomenon? This question was examined in a very interesting paper by Twiss (*7*), before the advent of quasars. He concluded that if the relativistic electrons have a population inversion, that is, if there are more electrons at higher energies than at lower energies, negative absorption is indeed possible. Unfortunately, as was shown later, Twiss had not included the statistical weights of the electrons correctly. A correct calculation (*8*) showed that under the circumstances described by Twiss, negative absorption is not possible. The conditions necessary for negative absorption have been reviewed in a very lucid fashion by Wild *et al.* (*9*).

I have examined the question of negative absorption again, and have found that it is possible after all. This conclusion differs from previous ones because the effect of an ambient plasma on the synchrotron radiation has been included. In order to describe this effect, I shall proceed from the discussion of Wild *et al.* (*9*). There, one finds the derivation of a formula (Eq. 27 in *9*, p. 348) for the general coefficient of absorption of a radiating system. For the special case of synchrotron radiation, it becomes (*3*)

$$K(\nu) = \frac{1}{4\pi m\nu^2}\int_1^{\infty}\mathrm{d}\gamma\,\frac{N(\gamma)}{\gamma^2}\,\frac{\mathrm{d}}{\mathrm{d}\gamma}\times [\gamma^2 p(\nu,\gamma)]\ \mathrm{cm}^{-1} \tag{4}$$

where $N(\gamma)\mathrm{d}\gamma$ is the number of electrons per unit volume having Lorentz factors between γ and $\gamma + \mathrm{d}\gamma$.

If, as is the usual case, $K(\nu)$ is positive, the radiation from an optically thick body comes mainly from a surface layer of thickness K^{-1} and is less than that due to spontaneous emission alone. But if $K(\nu)$ is negative, the radiation from an optically thick body may exceed the spontaneous radiation by the factor (approximately)

$$\tau(\nu)^{-1}\,e^{\tau(\nu)} \tag{5}$$

where

$$\tau(\nu) = -\int_0^L K(\nu,x)\mathrm{d}x \tag{6}$$

is the integral of the absorption coefficient across the radiating system, whose depth is L.

Equation 4 gives the criterion for negative absorption. The relativistic electron distribution $N(\gamma)$ must be confined primarily to an energy region such that $(\mathrm{d}/\mathrm{d}\gamma)[\gamma^2 p(\nu,\gamma)]$ is negative. But the $p(\nu,\gamma)$ given by Eq. 3 never decreases faster than $\gamma^{-2/3}$ (for fixed ν); therefore, regardless of the electron distribution, negative absorption is not possible for such a system of synchrotron radiators.

Now let us consider how the situation is altered by the presence of an ambient plasma. In this case, the synchrotron power spectrum has a striking difference from that of Eq. 3, known to radio astronomers as the Razin effect (*10–12*). The detailed discussion of this effect is rather involved; for our purposes, a heuristic discussion due to Scheuer (*11*) will suffice.

It is well known that a simple plasma

has an index of refraction to transverse waves which is less than one:

$$n(\nu) = \left(1 - \frac{\nu_p^2}{\nu^2}\right)^{\frac{1}{2}} \tag{7}$$

so that the phase velocity of light is greater than c. ν_p is the plasma frequency:

$$\nu_p = \left(\frac{n_e e^2}{\pi m}\right)^{\frac{1}{2}} \tag{8}$$

where n_e is the number density of thermal electrons. Roughly speaking, the result of this change of phase velocity is that the Lorentz factor γ in Eq. 2 (but not the one in Eq. 3) should be replaced by a new γ' as follows:

$$\gamma = \left(1 - \frac{\nu^2}{c^2}\right)^{-\frac{1}{2}} \rightarrow \gamma'$$
$$= \left(1 - \frac{n(\nu)^2\nu^2}{c^2}\right)^{-\frac{1}{2}} \tag{9}$$

This replacement is the key to the Razin effect. Synchrotron radiation is the radiation at high harmonics of the gyrofrequency. But if the frequency is sufficiently low so that $n(\nu)$ differs appreciably from one, then

$$\gamma' \approx \frac{\nu}{\nu_p} \tag{10}$$

and

$$\nu_c \approx \frac{3eH \cdot \nu^3}{4\pi m c \gamma \nu_p^3} \tag{11}$$

If one substitutes this new value of ν_c into the asymptotic form (3b), one obtains the asymptotic behavior valid for low frequencies:

$$p(\nu,\gamma) \sim \frac{\nu_r}{\nu} \exp\left(-\frac{\nu_r^2}{\nu^2}\right) \tag{12}$$

where

$$\nu_r = \left(\frac{4\pi\gamma m c \nu_p^3}{3\,eH}\right)^{\frac{1}{2}} \tag{13}$$

Consider the spectral power $p(\nu,\gamma)$ radiated at a fixed frequency as a function of the electron Lorentz factor γ. This function is shown in Fig. 1 for two cases: (a) $\nu_p = 0$, so that there is no Razin effect; (b) $\nu_p \neq 0$. The cutoff at the lower γ comes from the critical frequency (Eq. 2). The cutoff of curve b at the higher γ is the Razin cutoff. The existence of this cutoff makes negative synchrotron absorption possible. If the electron distribution $N(\gamma)$ has a low energy cutoff which is in the high energy tail of the function $p(\nu,\gamma)$ for some frequency, $K(\nu)$ can be negative.

An estimate of this coefficient may be obtained by assuming that the electron spectrum has a low energy cutoff at γ_1, and that the function $p(\nu,\gamma)$ is well approximated by Eq. 12. Under these assumptions, Eq. 4 yields

$$K(\nu) \approx -.7 r_0 c\, N_e \frac{\nu_p^3}{\nu^4} (\beta\gamma_1)^{\frac{3}{2}} e^{-\beta\gamma_1} \times$$
$$\left[1 - \frac{3}{\beta\gamma_1} + 0(\beta\gamma_1)^{-2}\right] \tag{14}$$

where

$$\beta = \frac{2\nu_p^3}{3\nu_0\nu^2} \tag{15}$$

r_0 is the classical electron radius, and γ_1 is the low energy cutoff on the electron spectrum. N_e is the number density of relativistic electrons having Lorentz factors between γ_1 and $\gamma_1 + \beta^{-1}$:

$$N_e = \int_{\gamma_1}^{\gamma_1 + \beta^{-1}} d\gamma\ N(\gamma) \tag{16}$$

From Eq. 14, one obtains a criterion for the absorption to be negative, namely, $\beta\gamma_1 > 3$, therefore

$$N(\gamma) \rightarrow 0 \text{ for } \gamma \lesssim \frac{9\nu_0\nu^2}{2\nu_p^3} \tag{17}$$

(the exact value of the numerical coefficient depends somewhat on the shape of the low energy electron cutoff).

Figure 2 illustrates the kinds of situations which can arise if there is an ambient plasma and an inverted distribution of relativistic electrons with a low energy cutoff at γ_1. The dashed curve is the electron distribution. The solid curves a, b, and c are functions $p(\nu,\gamma)$ for three frequencies in order of increasing frequency. In case a, the curves $p(\nu,\gamma)$ and $N(\gamma)$ do not overlap, so that the absorption is negligible. In case b, the electron distribution satisfies criterion 17 (see above) and there is appreciable overlap, hence negative absorption. In case c, criterion 17 is not satisfied and the absorption is positive. The spike of the dashed curve at $\gamma = 1$ represents the background plasma, which is not coupled to the synchrotron radiation. For maser action to be significant, criterion 17 must be satisfied over a path length L such that $\tau > 1$. In that case, the enhancement factor is given by expression 5 (above) and synchrotron radiation at that frequency is much more efficient than usual.

If the source were strictly homogeneous, it would be possible to predict the spectral signature of the negative absorption process. At low frequency, the synchrotron radiation is strongly suppressed by the Razin effect. At frequencies slightly above the threshold of criterion 17 there is a marked increase in intensity. The enhancement factor decreases rapidly with increasing frequency. There is a range of frequencies where $K(\nu)$ is positive and the synchrotron radiation is self-absorbed. At still higher frequencies, the source becomes optically thin to synchrotron radiation and the radiation almost entirely spontaneous emission. However, if the medium is not homogeneous, the spectral characteristics may be more complicated.

Let us now investigate the applicability of the negative absorption process to the radiation from quasars, using as a prototype 3C273, which is the best known of such objects. We refer in particular to the radio source 3C273B, the one coincident with the optical source. This source has a very peculiar spectrum (*13–16*), which is reproduced in Fig. 3 (*17*).

In order to provide the context for the negative absorption process, some of the salient features of this remarkable object are listed here briefly:

1) Lines in the optical spectrum re-

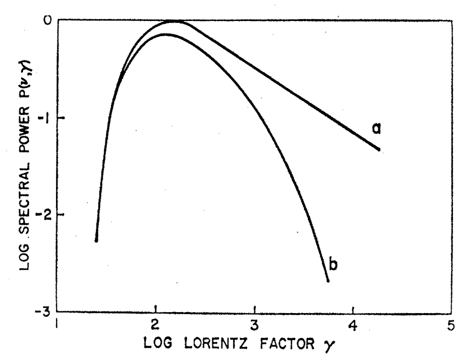

Fig. 1. Synchrotron power radiated at a given frequency versus electron Lorentz factor (arbitrary normalization). (a) No Razin effect ($\nu p = 0$). (b) $\nu p \neq 0$.

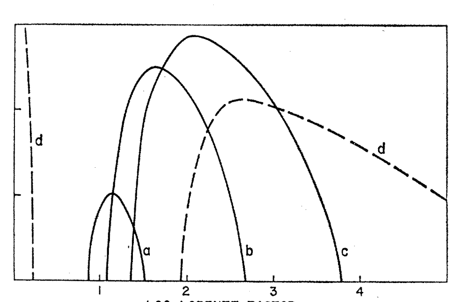

Fig. 2. The condition for negative absorption solid curves: same as Fig. 1, curve b (ordinate equals log spectral power) for three successively higher frequencies: $\nu(c) = 4\ \nu(b) = 16\ \nu(a)$. Dashed curve: a hypothetical electron spectrum (ordinate equals log $N_e(\gamma)$ which would result in negative absorption for $\nu(b)$.

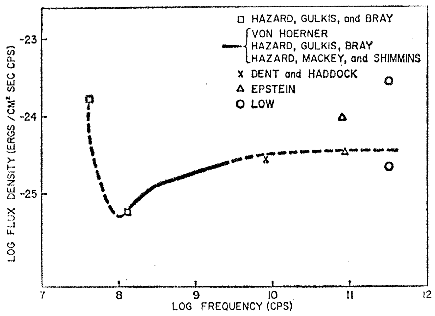

Fig. 3. The radio spectrum of 3C273B.

veal a red shift of $z = \Delta\lambda/\lambda = 0.158$. Gravitational red shift has been ruled out. Therefore, one is faced with the choice that the red shift is either an ordinary doppler shift, or that it is cosmological in origin, in which case a distance of 470×10^6 parsecs is assigned. The observation by Koehler (*18*) of a spectral absorption feature to the red of 21 cm, attributed to resonant absorption by neutral hydrogen in the Virgo cluster, argues for the cosmological hypothesis.

2) The source varies in intensity at both optical and radio frequencies. The optical variations consist of fluctuations of order 40 percent with characteristic times of order a few years and flare-like events of order 100 percent with characteristic times of order weeks (*19*). Radio variations of order 15 percent per year have been observed at frequencies of 1420 Mcy/sec and higher (*20–22*). The flux at lower frequencies remains nearly constant.

3) High resolution studies by lunar occultation (*13–15*) and by long baseline interferometry (*23*) reveal very complex structure in source B. It seems that the object has an angular size of order 1.0″ at 2695 and 1420 Mcy/sec, which decreases to less than 0.15″ at 430 Mcy/sec. There is a suggestion that the source is double or has a core-halo structure at lower frequencies (*14, 15*).

Although in most discussions of quasars it is assumed that the radio emission is due to electron synchrotron radiation (*24*), attempts to construct synchrotron models of radio source at cosmological distances encounter serious theoretical difficulties (*5, 25–27*). For homogeneous models, the absence of strong synchrotron self-absorption at frequencies greater than 400 Mcy/sec implies an angular size of $\theta > 0.05''$ (*20, 25, 28*) which is just barely below the resolved size. The observed secular variations compound the problem. It is difficult to construct a model of angular diameter $\theta > 0.005''$ which gives the observed radio variations. It is possible to resolve the angular size problem with inhomogeneous models, but then other difficulties arise. Such models must have total particle energies of order 10^{60} ergs in order to produce the observed radiation, and as Hoyle *et al.* (*27*) have pointed out, the energy loss rate of the electrons due to inverse Compton collisions is so great that the particle energy must be replenished in a very short time.

The situation in light of these difficulties has been succinctly summarized by Hoyle and Burbidge (*5*): "We arrive therefore at the conclusion that one or other of the following three possibilities represents the true situation: (i) the cosmological theory is incorrect, (ii) some process other than synchrotron radiation is responsible for the radio emission, (iii) a radically different model of the synchrotron radiation from that used above [the homogeneous synchrotron model] must be found." The observation of Koehler (*18*) and the absence of observed proper motion (*29*) argue against (i), although this possibility deserves further consideration, as does (ii). As for item (iii), it remains to be shown whether the negative absorption process can help to clear up the difficulties.

It is not possible to estimate the quantity $K(\nu)$ of Eq. 14 for 3C273B because it depends sensitively on unknown parameters such as the magnetic field, the plasma frequency, the number of relativistic electrons per unit volume (hence, the size of the object), and the low energy cutoff of the synchrotron electrons. It can only be demonstrated that it is possible that the enhancement is large. To do this, it is assumed that the object is at the cosmological distance, and, for simplicity, that it is homogeneous. The angular diameter of the object is taken to be $\theta = 0.01''$, which corresponds to a linear diameter approximately $D = 7.2 \times 10^{19}$ cm. The magnetic field strength is very uncertain. Most authors have suggested values between 10^{-2} and 10^{-4} gauss, based on equipartition arguments. Here the value $B = 10^{-2}$ gauss is adopted.

Given the volume of the object and the value of the magnetic field, one can now find the energy spectrum of relativistic electrons required to fit the observed spectrum of Fig. 3. Indeed, the shape of the spectrum for frequencies $\nu > 400$ Mcy/sec suggests an inversion in the electron spectrum, if it is assumed that the object is optically thin for those frequencies. Using the standard formulas (*2, 3*) one obtains the following electron spectrum:

$$N(\gamma) = \begin{cases} 0, \ \gamma \lesssim 3 \times 10^{2} \\ 10^{-2}\,\gamma^{-1}, \ \gamma \gtrsim 3 \times 10^{2} \end{cases} \mathrm{cm}^{-3} \tag{18}$$

The low energy cutoff at $\gamma \approx 3 \times 10^2$ is required to reproduce the behavior $p(\nu) \sim \nu^{\frac{1}{3}}$ in the frequency range 400 Mcy/sec $< \nu <$ 3000 Mcy/sec.

The electron spectrum inferred from the observed radio spectrum in this model satisfies one of the criteria for negative absorption. Now consider the spectrum for frequencies $\nu < 400$ Mcy/sec. In the range 100 Mcy/sec $< \nu <$ 400 Mcy/sec, the spectral index is greater than one-third. This kind of spectrum cannot be due to synchrotron radiation regardless of the electron spectrum, unless there is absorption or some suppression mechanism operating. Assume synchrotron self-absorption is responsible. In that case, another criterion for negative absorption is satisfied: the object is optically thick to its own radiation. Perhaps, then, the surprisingly large flux at 41.7 Mcy/sec is a manifestation of this process.

The only quantity which has not been specified for the evaluation of the $K(\nu)$ of Eq. 14 is the plasma frequency, or the density of the thermal plasma. Since this is unknown, we let it remain a free parameter and inquire whether τ can be large for *any* value of the plasma frequency. The optimum value of $K(\nu)$ will occur for $\beta\gamma_1 \approx 4$. This choice requires $\nu_p \approx 10^6$ cy/sec, hence a plasma density $n_e \approx 10^4$ cm^{-3}. Now the enhancement factor can be evaluated. From Eqs. 14–16 one obtains $K(\nu) \approx 2.7 \times 10^{-19}$ cm^{-1}, and taking the path length equal to the diameter, $\tau = 20$. The enhancement factor (see Eq. 15) is large.

Of course, this sample calculation is not to be taken as a serious model of 3C273B. It is only meant to demonstrate that in a model for this radio source with physical conditions rather similar to those in models previously proposed, negative absorption may play an important role. The assumption that it does raises a number of unanswered questions—for instance, how does the relativistic electron inversion arise? At this time, it is not at all certain that negative absorption is occurring in quasars. The answer to this question must await further high-resolution observations. On the other hand, the mechanism has some very attractive features which offer the hope of resolving some of the difficulties associated with the construction of models for quasars.

Perhaps the most dramatic feature of the mechanism is that there can be secular variations in the radiation from the entire source on a time scale considerably shorter than the light travel time through the source. This is a consequence of the geometric increase of radiation across the source. If the emission along a small segment of the path length varies by a certain percentage due to a change in the physical conditions there, the radiation from the entire source varies by a similar percentage, even though conditions remain unchanged elsewhere. Therefore, the radio variations do not imply that the angular diameter of the source must be less than 0.055″. In fact, if the 41.7-Mcy/sec flux is evidence of negative absorption, one might expect it to have large secular variations.

In the sample calculation above, it was assumed that the object is optically thin to radiation at frequencies $\nu > 400$ Mcy/sec, in order to derive the density of relativistic electrons. This may not be the case. If the object were optically thick, owing to either synchrotron self-absorption or free-free absorption, the density of relativistic electrons might be much greater than the value of Eq. 18. Indeed, such a model has been proposed (*21*, *30*). In that case the radiation at frequencies greater than 41.7 Mcy/sec might be partly due to negative absorption, and the observed secular variations might be explained by this mechanism.

The addition of negative absorption to the theory of synchrotron radiation opens some interesting possibilities for the construction of detailed models of 3C273B. One that immediately comes to mind is a core-halo structure, already indicated by observation (*14*). Perhaps the high-frequency radiation is due to negative absorption in the core, which is not radiating at lower frequencies owing to the Razin effect. The low-frequency radiation may be due to negative absorption in the halo, which is optically thin to the high-frequency radiation. Then one might see a very pronounced "limb brightening" at low frequencies. Furthermore, there is no reason to expect the physical conditions in quasars to be perfectly isotropic. A slight anisotropy in physical conditions will yield great anisotropy in the amplified radiation. If the radiation is a highly directional beam, the energy requirement of the source is greatly reduced (*31*). There is a possibility that the negative absorption may be marginal over a broad band of frequencies; for, if self-absorption, positive or negative, is very strong at a given frequency, there is a tendency for the electron spectrum to rearrange itself in order to level out the radiation spectrum. Finally, there is a hope of understanding the secular variations as arising from the high sensitivity of the absorption coefficient to the physical parameters, such as the magnetic field strength or the plasma density (*32*).

RICHARD MCCRAY

Space Science Center, University of California, Los Angeles 90024

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32. Supported by an NSF predoctoral fellowship. I thank Dr. Peter Goldreich for stimulating my interest in this subject.

10 October 1966

Kodiak Seamount Not Flat-Topped

Abstract. *Earlier surveys in the Aleutian Trench southeast of Kodiak Island, Alaska, indicated that Kodiak Seamount had a flat top and was a tablemount or guyot. This seamount is of special significance because it has been supposed that its surface was eroded at the same time as those of a line of guyots to the southeast. If so, its present position in the axis of the Aleutian Trench indicates that the line of guyots was formed before the trench. A two-part survey in 1965 showed that Kodiak Seamount is not flat-topped, and should be eliminated from the category of guyots. Reflection profiling records indicate that the seamount was formed before the adjacent sediments were deposited, and that the small trough, or moat, on the south side is a depositional feature probably formed by a scouring effect or by the acceleration of turbidity currents around the base of the mount.*

Kodiak Seamount has been the subject of unusual interest and significance in studies of the geology of the Gulf of Alaska and the Aleutian Trench. Earlier surveys of Kodiak Seamount indicated that it had a flat top and could be designated a tablemount or guyot (*1*). Flat-topped seamounts are