Letters 3, 307 (1966)]; Yu. P. Raizer, Zh. Eksperim. i Teor. Fiz.—Pis'ma Redakt. 4, 3 (1966) [translation: JETP Letters 4, 1 (1966)]. In this context we also refer to a paper by K. A. Brueckner and S. Jorna, Phys. Rev. Letters 17, 78 (1966). There linearized hydromechanical equations of momentum and energy are derived, which predict instabilities in the propagation of a laser beam.

[3]E. Garmire, R. Y. Chiao, and C. H. Townes, Phys. Rev. Letters 16, 347 (1966).

[4]R. Y. Chiao and E. Garmire, Fourth International Quantum Electronics Conference, Phoenix, April, 1966 (to be published); R. Y. Chiao, M. A. Johnson, S. Krinsky, H. A. Smith, C. H. Townes, and E. Garmire, to be published.

[5]Ya. B. Zel'dovich and Ya. P. Raizer, Zh. Eksperim. i Teor. Fiz.—Pis'ma Redakt. 3, 137 (1966) [translation: JETP Letters 3, 86 (1966)]; Y. R. Shen, Phys. Letters 20, 378 (1966).

[6]H. Fröhlich, Proc. Phys. Soc. 87, 330 (1966).

[7]V. L. Ginzburg and L. D. Landau, Zh. Eksperim. i Teor. Fiz. 20, 1064 (1950).

[8]A. A. Abrikosov, Zh. Eksperim. i Teor. Fiz. 5, 1174 (1957) [translation: Soviet Phys—JETP 32, 1442 (1957)].

[9]R. Y. Chiao and P. A. Fleury, Proceedings of the International Conference on the Physics of Quantum Electronics, edited by P. L. Kelley, B. Lax, and P. E. Tannenwald (McGraw-Hill Book Company, Inc., New York, 1966).

[10]D. Hondros and P. Debye, Ann. Physik. 32, 465 (1910). See also A. Sommerfeld, Elektrodynamik (Akademische Verlagsgesellschaft, Leipzig, 1961), p. 177.

# MASER MODEL FOR INTERSTELLAR OH MICROWAVE EMISSION

M. M. Litvak, A. L. McWhorter, M. L. Meeks, and H. J. Zeiger
Lincoln Laboratory,* Massachusetts Institute of Technology, Lexington, Massachusetts
(Received 4 August 1966)

Recent observations[1-4] of 18-cm OH emission from the direction of certain H II regions have shown remarkable anomalies which have led a number of workers[2-4] to suggest that maser amplification is involved. We will present a maser amplifier model, based on ultraviolet continuum pumping from nearby stars, which seems capable of explaining most of the observations, including those of anomalous absorption. We find that sufficient gain is obtained only by taking into account selective absorption of the uv radiation as it penetrates into the OH cloud, a process leading to the weakening of the major anti-inverting uv transitions. Other authors[5] who have independently considered uv pumping but have neglected selective absorption are unable to explain the observed intensity or polarization of the OH emission. We explain the polarization properties in terms of the electromagnetic normal modes of the maser medium by including the contribution to the susceptibility tensor from both the OH molecules and the electrons and by considering the nonlinearities resulting from gain saturation. The intensity ratios of the different microwave transitions have not yet been accounted for, however.

Experimental situation.—The OH emission spectrum consists of four ground-state Λ-doublet transitions at 1612, 1665, 1667, and 1720 MHz, with the 1665-MHz transition generally strongest and the 1612- and 1720-MHz weakest. Each transition exhibits several narrow Doppler-shifted features (~1 kHz wide) with some anticorrelation between the features of different hyperfine transitions as a function of frequency. The features are all strongly polarized, predominantly circular (left or right), but frequently elliptical. The emission regions have apparent angular sizes <15 sec arc,[6] and are separated by 4-12 min arc from bright centers of continuum radio emission. The antenna temperatures of most of the features are many times the antenna temperature of the associated continuum source, and the brightness temperature of some features exceeds $2\times10^6$ °K.[6] In contrast, OH absorption Doppler widths indicate translational temperatures of ~100°K. However, anomalies in absorption have also been noted; in particular, OH regions in the direction of the galactic center show strong, decidedly nonthermal absorption that cannot be explained in terms of high optical depth.[3] Since anti-inversion is produced by the presence of some weak uv pumping, the thermal-equilibrium estimate[7] of one OH for every $10^4$ atoms in these anomalous regions may be too high.

Pump mechanism.—The pumping of the OH by the uv occurs via the six uv transitions[8] shown in Fig. 1 which lead from the ground state $^2\Pi_{3/2}$ to the electronically-excited $^2\Sigma_{1/2}{}^+$. We will assume a uv flux of $10^{-17}$ W/m$^2$ Hz, correspond-

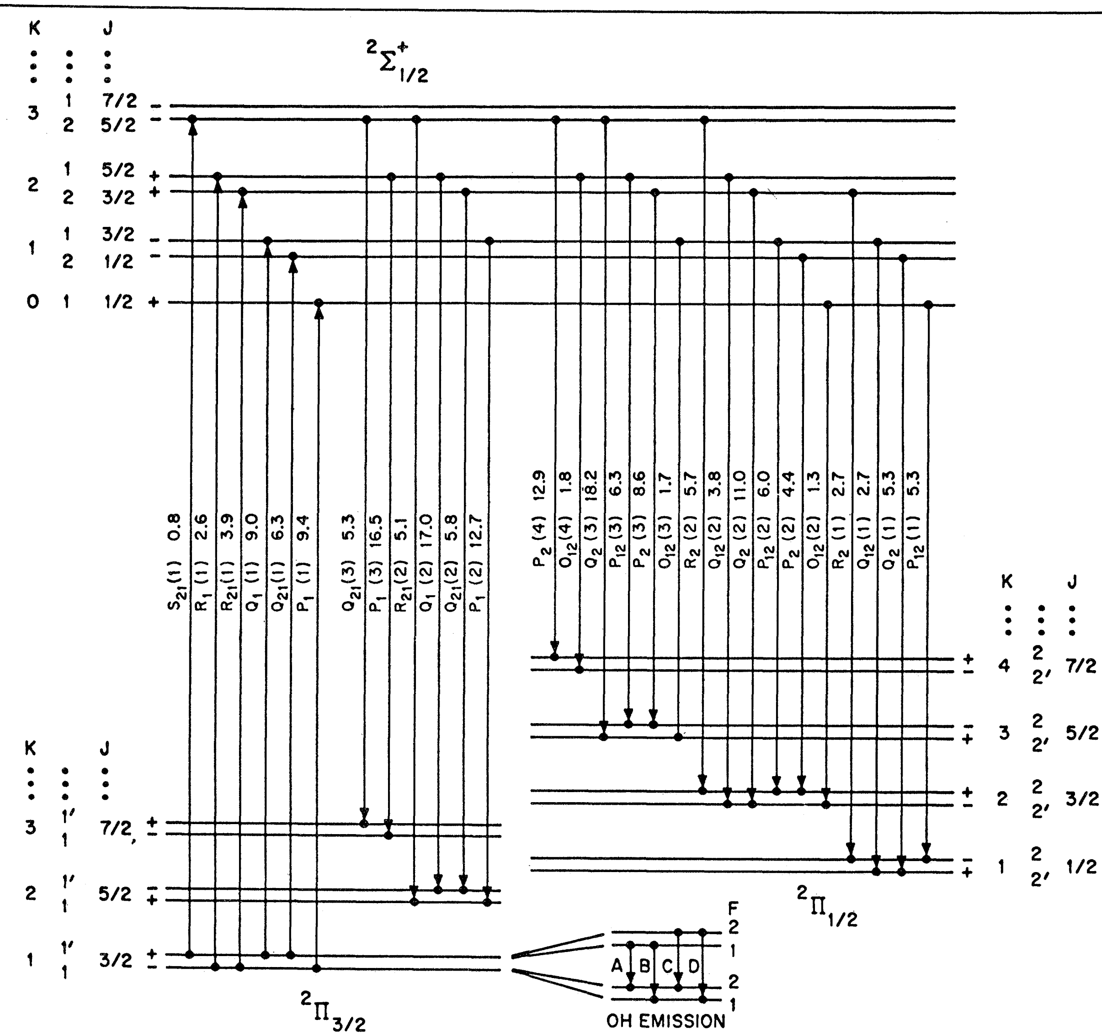


FIG. 1. Energy-level diagram (not to scale) of OH molecule showing rotational states of the lowest vibrational state for the $^2\Pi_{3/2}$, $^2\Pi_{1/2}$, and $^2\Sigma_{1/2}^+$ electronic states. Zeeman splittings are not shown; hyperfine splittings are indicated only for the ground-state $\Lambda$ doublet. The four microwave emission frequencies are 1612, 1665, 1667, and 1720 MHz, corresponding to the transitions labeled $A$, $B$, $C$, and $D$, respectively. Also indicated (using the notation of Ref. 8) are the pertinent uv transitions together with their relative strengths.

ing to an $O5$ star[9] 2 lt yr from the OH region.[10] Once a rotational state of $^2\Sigma$ is excited by the uv, which for the above flux will take about $3\times10^7$ sec,[11] decay by uv spontaneous emission back to the $^2\Pi$ states will occur in about $10^{-6}$ sec. Then by far-infrared spontaneous emission, the molecule will cascade down the $^2\Pi$ rotational states in about 1 sec[12] and return to one of the $\Lambda$-doublet ground states.[13]

When the uv pumping scheme is examined in detail, it is found that the strongest pumping routes involve excitation by $Q_1(1)$ and $P_1(1)$, which are anti-inverting and inverting transitions, respectively. At small optical depths, the uv pumping leads to an anti-inversion of the $\Lambda$ doublet when the populations are summed over all $M_F$ sublevels.[5] However, as the uv penetrates into the OH cloud, the $Q_1(1)$ and $P_1(1)$ wavelengths are absorbed most rapidly, followed by absorption of the anti-inverting $Q_{21}(1)$ radiation. The $R_{21}(1)$ and $R_1(1)$ transitions then become dominant and a sizable shell, of thickness inversely proportional to the OH concentration, can be inverted quite strongly.

Figure 2 shows the result of a detailed calculation of the $\Lambda$-doublet population inversion as a function of optical depth, based upon the uv line strengths as shown in Fig. 1 and the assumption of equal $M_F$ populations in each $\Lambda$-doublet state. Also shown in Fig. 2 is the decrease with distance of the sum of the six uv-induced transition rates between the $\Lambda$-doublet states; this quantity will be used to estimate the degree of microwave saturation. The effect of reradiated uv, electron collisions, and microwave saturation are not included in Fig. 2. These are being taken into account in a more detailed calculation now in progress, in which the $M_F$ populations are self-consistently determined as a function of depth[14] and microwave intensity.

Intensity and linewidth.—For purposes of making order-of-magnitude estimates, we will treat the OH maser in this section as if there were a single microwave transition between the $\Lambda$-doublet states, with all magnetic sublevels equally populated. In the presence of microwave saturation and electron collisions, the $\Lambda$-doublet population difference will be reduced from the value $\Delta n_0$ given in Fig. 2 to

$$\Delta n = \Delta n_0[1+(2W_m+2W_e)/W_{\text{uv}}]^{-1}, \tag{1}$$

where $W_m$, $W_e$, and $W_{\text{uv}}$ are the induced transition rates between the $\Lambda$-doublet states due to the microwave, electrons, and uv, respectively. $W_{\text{uv}}$ varies with depth as shown in Fig. 2, the maximum being about $3\times10^{-8}$ sec$^{-1}$ for the assumed source of uv. Following Seaton,[15] we calculate that $W_e \sim 7\times10^{-5} n_e T_e^{-1/2}$ sec$^{-1}$, where $n_e$ is the electron density in cm$^{-3}$ and $T_e$ is the electron temperature in °K. We now

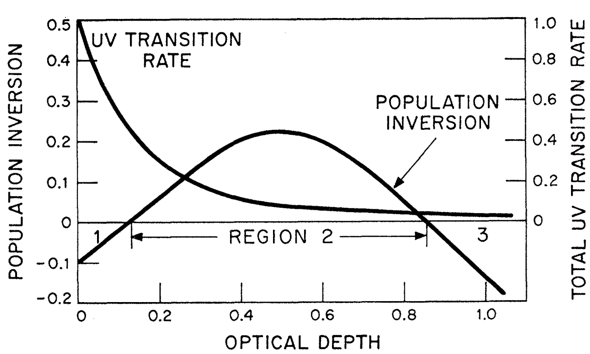


FIG. 2. Population inversion of OH $\Lambda$ doublet due to isotropic uv pumping as a function of optical depth into the OH cloud, showing the effect of selective uv absorption. Inversion is expressed as the ratio of the population difference between upper and lower $\Lambda$-doublet states to the number of OH molecules. Regions 1 and 3 are anti-inverted and would show anomalous absorption; region 2 is inverted. For an OH concentration of $10^{-2}$ cm$^{-3}$, unit optical depth is about 0.15 lt yr. The decrease of the total uv transition rate with optical depth is also shown, unit rate being about $3\times10^{-8}$ sec$^{-1}$ for the uv flux assumed.

show that for $T_e \gtrsim 100$°K, the maser will be saturated ($W_m > W_e, W_{\text{uv}}$) unless $n_e \gg 1$ cm$^{-3}$.

Let us first assume to the contrary that the maser is unsaturated. Then if $T_B$ is the maximum brightness temperature (as a function of angle) of a point on the surface of the OH cloud, we require for $kT_B \gg h\nu$ that

$$h\nu A n_b(e^{\alpha L}-1)/4\pi\alpha\delta\nu = 2kT_B/\lambda^2, \tag{2}$$

where the gain coefficient $\alpha$ is

$$\alpha = \lambda^2 A\Delta n/8\pi\delta\nu, \tag{3}$$

$A$ is the Einstein coefficient, $n_b$ the upper-state population, $L$ the path length of maser gain in the direction under consideration, and $\delta\nu$ the (full) linewidth of the Doppler feature. The experimental lower bound[6] of $2\times10^6$ °K for $T_B$ requires $\alpha L \gtrsim 16$, which implies[16] a solid angle of maser emission through the point of $\Omega_m \sim (\alpha L)^{-1} \sim 0.06$. Expressing the microwave spectral intensity $I_m$ as $2kT_B\Omega_m/\lambda^2$, we can then use the Einstein relation to determine that

$$W_m = BI_m = AkT_B\Omega_m/4\pi h\nu, \tag{4}$$

which yields $W_m \gtrsim 2\times10^{-5}$ sec$^{-1}$ for $A = 7\times10^{-11}$ sec$^{-1}$[17] and the above values of $T_B$ and $\Omega_m$. Therefore, $W_m \gg W_{\text{uv}}$, and unless $n_e$ is extremely large, $W_m > W_e$ as well, which means that the maser must, in fact, be saturated.

Under saturated gain conditions with $W_m \gg W_{\text{uv}}$ and $W_e$, the microwave intensity reaching the antenna is

$$\begin{aligned} I_A &= h\nu W_m \Delta n V/R^2\Omega_0\delta\nu_m \\ &= h\nu W_{\text{uv}}\Delta n_0 V/2R^2\Omega_0\delta\nu_m, \end{aligned} \tag{5}$$

where $V$ is the effective volume of the OH, $R$ is the distance to the antenna, and $\Omega_0$ is the beam angle of the maser emission. We assume that a single feature arises from perhaps several regions of OH having essentially the same velocity, with a total effective path length $L$, width $w$ transverse to the uv flux, and an inverted layer of thickness $d$. Unit optical depth in Fig. 2 is about $1.5\times10^{21} n_{\text{OH}}^{-1}(\delta\nu/\nu)_{\text{uv}}$ cm. Therefore, if we take the average $W_{\text{uv}} \sim 0.1$ of the maximum $W_{\text{uv}}$ to account for absorption, the effective value of $d\Delta n_0$ is about $2\times10^{20}\cos\theta_0 \times(\delta\nu/\nu)_{\text{uv}}$ cm$^{-2}$, where $\theta_0$ is the angle between the uv flux and the surface normal. With $\langle W_{\text{uv}}\rangle = 3\times10^{-9}$ sec$^{-1}$, $(\delta\nu/\nu)_{\text{uv}} \approx (\delta\nu/\nu)_m$, and $R = 4000$ lt yr and $\cos\theta_0 \sim 1$, we find that to explain the observed value of $I_A \approx 5\times10^{-24}$ W/m² Hz of the brightest feature of the 1665-MHz line of the OH-emission region near W3, we require $wL/\Omega_0 \approx 50$ (lt yr)²/sr. A calculation of $\Omega_0$ under nonlinear saturation conditions has not been made. However, if we assume that the beam angle is determined purely geometrically, we obtain $\Omega_0 \approx 4dw/L^2$. Taking $(\delta\nu/\nu)_m = 10^{-6}$ and $\Delta n_0/n_{\rm OH} \approx 0.2$ from Fig. 2, we can then rewrite the above requirement on $wL/\Omega_0$ in the form $n_{\rm OH} \approx 0.2L^{-3}$ cm$^{-3}$, where $L$ is in lt yr. This condition does not appear at all unreasonable. The corresponding value of $d$ would be $5\times10^{-3}L^3$ lt yr. If the maser is strongly saturated, the microwave intensity at the output is

$$I_m = h\nu W_{\rm uv}\Delta n_0 L/2\delta\nu_m, \tag{6}$$

which yields $W_m = 4\times10^{-5}L^{-2}$ sec$^{-1}$ and $T_B = 2\times10^7 w^{-1}L^{-3}$ °K (where $w$ and $L$ are in lt yr) upon using the above parameter values.

From (1) and (3) we find that in order to reach saturation levels, we need a path length $L_0$ of exponential growth of about $1\times10^{14}n_{\rm OH}^{-1}(1+2W_e/W_{\rm uv})$ cm, which for $n_e \lesssim 0.1$ cm$^{-3}$ and for the previous values of the other parameters gives $L_0 \lesssim 0.3L^3$ lt yr. It should be remembered that this relation as well as those immediately above depend on the assumption about $\Omega_0$.

During growth in the unsaturated region the maser linewidth will narrow by the usual factor $(\ln 2/\alpha L_0)^{1/2}$, about 3 or 4 in this case. Normally, the line would become broad again during the stage of saturated growth.[18] However, collisions with hydrogen might make the saturation sufficiently homogeneous over the Doppler line by transferring OH molecules across the Doppler feature that little broadening would occur except for the most intense features. A hydrogen concentration of $10^4$ cm$^{-3}$ would give a collision rate of about $10^{-6}$ sec$^{-1}$, comparable to $W_m$.

Polarization.—We shall first show by a general consideration of the electron and OH contributions to the dielectric susceptibility tensor $\chi$ that elliptically polarized maser emission can easily occur. However, it is more difficult to explain the gain preference of one sense of polarization over the other in a single Doppler feature without invoking circular polarization in the uv pump. As we shall discuss, the gain preference could result from asymmetric features and saturation effects.

In a coordinate system with the dc magnetic field $\vec{H}$ along the $z$ axis, $\chi$ has the form

$$\chi = \begin{bmatrix} \chi_{11} & \chi_{12} & 0 \\ -\chi_{12} & \chi_{11} & 0 \\ 0 & 0 & \chi_{33} \end{bmatrix}$$

regardless of the uv pump direction or the presence of microwave saturation, since the Zeeman-splitting frequencies are much larger than the inverse lifetimes of the OH ground-state $\Lambda$ doublet. It is convenient to define $\eta = 2\chi_{12}\cos\theta/\chi_{11}$ and $\zeta = (\chi_{33}-\chi_{11})\sin^2\theta/\chi_{11}$, where $\theta$ is the angle between $\vec{H}$ and the wave vector $\vec{k}$. Then $k = \omega/c + \Delta k$, where

$$\Delta k = [2+\zeta\pm(\zeta^2-\eta^2)^{1/2}](\pi\omega/c)\chi_{11}. \tag{7}$$

The gain coefficient is $\alpha = 2\,{\rm Im}\Delta k$ and the ratio of the electric field components transverse to $\vec{k}$ is

$$E_x/E_y = -\eta^{-1}[\zeta\pm(\zeta^2-\eta^2)^{1/2}] \tag{8}$$

in the coordinate system with $\hat{y} = \hat{H}\times\hat{k}$ and $\hat{x} = \hat{y}\times\hat{k}$.

If we assume that the uv is somewhat directional but not circularly polarized, then for unsaturated conditions $\chi_{11}({\rm OH}) \neq \chi_{33}({\rm OH})$ but $\chi_{12}({\rm OH}) = 0$; i.e., the $|M_F|$ populations may be unequal but there is no skewing to positive or negative $M_F$. However, self-saturation by circularly polarized microwave emission or cross saturation by circularly polarized emission of the same frequency coming from other regions could produce a $\chi_{12}({\rm OH})$. For sufficiently large values of $H$ and $n_e$, $\chi_{12}(e) \approx i\omega_p^2\omega_c/4\pi\omega^3$ is important; but the electron anisotropy $\chi_{11}(e)-\chi_{33}(e)$, smaller by a factor of $\omega_c/\omega \sim 10^{-8}$ for $H \sim 10^{-5}$ G, should always be negligible. As usual, $\omega_c = eH/mc$ and $\omega_p^2 = 4\pi n_e e^2/m$. The isotropic part $-(\omega_p^2/4\pi\omega^2)$ of $\chi(e)$ will be included in the wave velocity $c$ rather than in $\Delta k$.

Two limiting cases arise in connection with Eqs. (7) and (8): (i) $|\eta/\zeta| \ll 1$. In this OH-dominated regime the two modes are very nearly linearly polarized, but with gain coefficients that differ by $\Delta\alpha = (4\pi\omega/c)\,{\rm Im}(\zeta\chi_{11})$. (ii) $|\eta/\zeta| \gg 1$. If $\chi_{12}({\rm OH})$ is negligible then the electrons control the polarization. Both modes are nearly circular and have almost the same unsaturated gain. However, by expanding the square root in (7) and keeping the lowest term in $(\zeta/\eta)^2$, we find that the peak gains for right and left polarized modes are split apart in frequency by approximately $\epsilon\delta\nu$ and in magnitude by about

$2\alpha\epsilon \operatorname{Re}\chi_{11}/|\chi_{11}|$, where $\epsilon = |\zeta^2/\eta(2+\zeta)|$. The above quantities are evaluated at the frequency for which $\operatorname{Im}\chi_{11}$ has a maximum. If the Doppler features do not have a symmetric shape, perhaps because of the overlap of several features associated with interpenetrating OH streams, $\operatorname{Re}\chi_{11} \neq 0$ where $\operatorname{Im}\chi_{11}$ is maximum. Both the frequency split and amplitude difference of the two gain peaks could amount to several percent. While this alone does not explain the high degree of polarization observed, this effect could well initiate nonlinear suppression of one sense of polarization over the other, especially with fast relaxation across the Doppler feature by hydrogen collisions, which would mainly conserve the magnetic and other quantum numbers. Similar phenomena occur in optical masers in magnetic fields, but even for these laboratory cases the analyses are difficult and incomplete.[19] If $\chi_{12}(\mathrm{OH}) \neq 0$ because of saturation or pumping by somewhat circularly polarized uv, the gain difference between right and left polarizations might be considerably larger than the value given above.

In order to be in regime (ii) we need $|\eta/\zeta| > 1$ on line center, which requires $n_e > \zeta\Delta nAmc^3/16\delta\nu\omega_c e^2\cos\theta$ if $\chi_{12}(\mathrm{OH}) = 0$. Here we have used (3) and the fact that $\alpha \approx (4\pi\omega/c)\operatorname{Im}\chi_{11}$. As $\Delta n$ decreases because of saturation this requirement becomes easier to meet. Taking $H = 10^{-5}$ G and the previous values for the other parameters, we find with the aid of (1) and (6) that we must have $n_e > (10\zeta/L^3\cos\theta)^{1/2}$, where $L$ is in lt yr, if regime (ii) holds in the unsaturated growth region or $n_e > 2\zeta/L\cos\theta$ if crossover from regime (i) to (ii) occurs in the saturated region. The latter condition is actually independent of $W_{\mathrm{uv}}$ and $\Omega_0$. Since $\zeta$ is expected to be small because of strong electron collisions, values of $n_e \sim 0.1\ \mathrm{cm}^{-3}$ should suffice for the existence of circular polarization. If $\chi_{12}(\mathrm{OH}) \neq 0$, the condition to be in regime (ii) may be even less stringent.

Intensity ratios.—Since the line strengths of the 1612-, 1665-, 1667-, and 1720-MHz transitions are in the ratio 1:5:9:1, the generally observed greater intensity of the 1665-MHz transition has been quite surprising. We cannot yet tell whether the uv pumping model presented here is able to account for the observed ratios because the detailed calculations including electron collisions and saturation effects are not completed. However, we have determined that saturation of the 1667-MHz line will not prevent strong exponential growth of the other lines, and that, in fact, the 1665-MHz line can grow to an intensity comparable to that of 1667 MHz after the latter has saturated. Subsequent absorption in the 9:5 ratio as the radiation passes through noninverted OH regions could then result in a relatively higher 1665-MHz intensity. Alternatively, the saturated growth process may be responsible for the higher intensity at 1665 MHz, since with sufficient inequality of the $M_F$ populations the saturated gain for 1665 MHz could exceed that for 1667 MHz. If this is so, then in cases where the emission emerges from the inverted region with little or no saturated growth, the 1667-MHz feature could be dominant but would generally be less intense than other features at 1665 MHz which benefited from saturated growth.

It is worth remarking that although uv pumping appears adequate to account for the OH observations, some of the discussion in the paper (particularly that on polarization) would remain valid if a still stronger pumping mechanism were discovered. Conversely, the problem of the intensity ratios is not peculiar to uv pumping, but would be shared by most pumping schemes.

We thank W. Liller for discussions on the uv flux estimates, M. A. Gordon for discussion on magnetoplasma polarization effects, and F. Perkins for communicating some unpublished calculations on the rotational transition probabilities to check against our own.

---

*Operated with support from the U. S. Air Force.

[1]H. Weaver, D. R. W. Williams, N. H. Dieter, and W. T. Lum, Nature 208, 29 (1965).

[2]S. Weinreb, M. L. Meeks, J. C. Carter, A. H. Barrett, and A. E. E. Rogers, Nature 208, 440 (1965).

[3]R. X. McGee, B. J. Robinson, F. F. Gardner, and J. G. Bolton, Nature 208, 1193 (1965).

[4]A. H. Barrett and A. E. E. Rogers, Nature 210, 188 (1966).

[5]F. Perkins, T. Gold, and E. E. Salpeter, Astrophys. J. 145, 361 (1966), report a calculation in which, for the uv directed along the magnetic field, the sublevels with large $|M_F|$ values can be inverted slightly, and weak amplification of linearly polarized emission is predicted. We disagree with these authors on several points, as supported by explicit calculation: (1) The effect of the hyperfine splittings on the $M_F$ populations cannot be deduced simply by the application of Clebsch-Gordan coefficients to the $M_J$ population calculation, which first neglects the hyperfine interaction. (2) Even using the authors' method of calculation we disagree

with the steady-state populations given in their Table I, and find no population inversion.

[6]Based on measurements of OH-emission region near W3. A. E. E. Rogers, J. M. Moran, P. P. Crowther, B. F. Burke, M. L. Meeks, J. A. Ball, and G. M. Hyde, Phys. Rev. Letters 17, 450 (1966); D. D. Cudaback, R. B. Read, and G. W. Rougoor, Phys. Rev. Letters 17, 452 (1966).

[7]B. J. Robinson, F. F. Gardner, K. J. Van Damme, and J. G. Bolton, Nature 202, 989 (1964).

[8]G. H. Dieke and H. M. Crosswhite, J. Quant. Spectr. Radiative Transfer 2, 97 (1962).

[9]C. W. Allen, Astrophysical Quantities (The Athlone Press, The University of London, London, England, 1963), p. 201. The near-uv color temperature was taken as approximately that in the visible, 70 000°K.

[10]Obscuration by interstellar dust prevents optical determination of the stars in the vicinity of the OH regions, but the existence of nearby *O* stars, or even clusters of early-type stars, is of reasonable probability. Because a high concentration of H would be expected in these regions, there should be an HII region between the stars and the OH emission point to absorb the detrimental ionizing uv flux. Such an HII region should emit microwave continuum radiation observable with an antenna having sufficient resolution and sensitivity. (If the OH region has an angular size <0.5 sr as viewed from the *O*5 star, the microwave antenna temperature would be <0.5°K for a 120-ft antenna like Haystack.) The 3080-Å Balmer continuum radiation produced at the OH by the HII layer would be about an order of magnitude less than the direct uv flux from the star.

[11]D. M. Golden, F. P. Del Greco, and F. Kaufman, J. Chem. Phys. 39, 3034 (1963). Total band oscillator strength for $^2\Pi_{3/2} \rightarrow {}^2\Sigma_{1/2} = 8\times 10^{-4}$.

[12]Based on an electric dipole moment of 1.54 Debye units. P. F. Wacker, M. Mizushima, J. D. Peterson, and J. R. Ballard, Microwave Spectral Tables—Diatomic Molecules (U. S. Government Printing Office, Washingto, D. C., 1964), NBS Monograph 70, Vol. 1, p. 31.

[13]Because OH is intermediate between Hund's coupling cases *a* and *b*, the transition rates between $^2\Pi_{1/2}$ and $^2\Pi_{3/2}$ are comparable to those between the rotational states of $^2\Pi_{1/2}$ and cannot be ignored. The transition probabilities were calculated using formulas of E. Hill and J. H. Van Vleck, Phys. Rev. 32, 250 (1928). Typographical errors in Eq. (9) for the matrix elements were corrected. Estimates of the far-infrared flux ($\sim 10^{-19}$ W m$^{-2}$ Hz$^{-1}$) from the neighboring star and HII regions indicate an absorption rate per OH molecule of $\sim 10^{-8}$ sec$^{-1}$. Because of the rapid decay directly back to the ground states, there is negligible effect of the IR on the population inversion caused by the uv and there is negligible population in the excited rotational states.

[14]The quantization axis is taken along the magnetic field. Because of the large population inversion allowed by selective absorption, the angle of uv pumping with respect to the magnetic field is not very important, contrary to cases of weak inversion similar to that in Ref. 5. The skewing of the $|M_F|$ populations by unpolarized but anisotropic uv radiation has a simple pattern. Consider first uv directed primarily along the magnetic field $H$. Then one finds that $F \rightarrow F$ uv excitations affect most strongly the sublevels with small $|M_F|$ values and $F \rightarrow F \pm 1$ excitations affect mostly the large $|M_F|$ values. Hence in region 1, where there is over-all anti-inversion, the $P_1(1)$ transition can invert the large $|M_F|$ sublevels; while at the beginning of region 2 the $Q_{21}(1)$ transition can anti-invert the large $|M_F|$ sublevels despite the over-all inversion. Further into region 2 there is inversion of all $|M_F|$ values due to the combined effect of $R_1(1)$ and $R_{21}(1)$. For uv radiation in the plane perpendicular to $H$, the effect on the large and small $|M_F|$ values is reversed.

[15]M. J. Seaton, in Atomic and Molecular Processes, edited by D. R. Bates (Academic Press, Inc., New York, 1962), p. 414.

[16]For a long, thin geometry $\Omega_m$ would be less than $(\alpha L)^{-1}$ but the corresponding $T_B$ would probably be greater than $2\times 10^6$ °K.

[17]The early estimates of $A$ by A. H. Barrett, IEEE Trans. Mil. Electron. MIL-8, 156 (1964), have been revised several times: by W. M. Goss and H. Spinrad, Astrophys. J. 143, 989 (1966), and most recently by B. Turner, private communication. The latest values are (7.7 and 7.1)$\times 10^{-11}$ sec$^{-1}$ for 1667 and 1665 MHz, respectively.

[18]D. F. Hotz, Appl. Opt. 4, 527 (1965).

[19]C. V. Heer and R. D. Graft, Phys. Rev. 140, A1088 (1965); W. Culshaw and J. Kannelaud, Phys. Rev. 145, 257 (1966).