EMISSION AND ABSORPTION OF RADIO WAVES IN THE GALACTIC DISK

By R. J. Gould*

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Summary

Bremsstrahlung in electron collisions with helium atoms produces up to 0·04% of the radio emission continuum from the main part of interstellar HII regions, while, near the outer edge where hydrogen is about half-ionized, electron collisions with hydrogen atoms produce up to 2% of the normal electron–proton bremsstrahlung. Radio emission and absorption in electron–hydrogen and electron–helium collisions in HI regions are investigated and found to be negligible at the low temperatures therein. However, low energy cosmic rays may ionize HI regions to a sufficient degree to significantly attenuate radio waves by inverse bremsstrahlung (free–free absorption) in electron–proton Coulomb collisions. Future high resolution surveys of the cosmic radio brightness should be able to distinguish between HI and HII absorption.

The dependence of the mean absorption coefficient on height above the galactic plane is considered for attenuation in HI and HII regions. Assuming only absorption in HII regions and a temperature therein of 6000°K, from the low frequency turnover in the spectrum of Cas A the mean squared electron density in the galactic plane within 3 kpc of the Sun is evaluated to be $0 \cdot 06_3$ cm^{-6}, in agreement with the value derived from thermal radio emission measurements.

I. Introduction

It is now well established that the bulk of the radio emission from the galactic disk at high frequencies originates from bremsstrahlung in Coulomb collisions of free electrons with hydrogen and helium ions in interstellar HII regions. Moreover, the reverse physical process (inverse bremsstrahlung) attenuates low frequency radio waves, and both the cosmic radio background spectrum and the spectra of a number of discrete sources are observed to turn over at the low frequency end. The theoretical problem of Coulomb bremsstrahlung has also been treated by a number of authors from both the classical and quantum-electrodynamic approaches; references to this literature may be found in the review by Oster (1961).

The principal aim of the present paper is to emphasize the possibility that there may be appreciable absorption of radio waves in interstellar HI regions where a small amount (perhaps $\sim 0 \cdot 1$–1%, see Section V) of ionization is produced by cosmic rays. That there could be appreciable ionization by cosmic rays has become evident through recent experimental and theoretical work which indicates a high flux of low energy (nonrelativistic) cosmic rays outside the solar system (Gloeckler and Jokipii 1967). Actually, even without cosmic ray ionization, elements like carbon, which has an ionization energy less than 13·6 eV, would be ionized by stellar ultraviolet radiation,

* On leave to the School of Physics, University of Sydney, Sydney, N.S.W. 2006, from permanent address: Department of Physics, University of California, San Diego, California, U.S.A.

producing an electron density $n_e \sim 10^{-3} n_H$ in HI regions. However, as will be shown, only cosmic rays can ionize HI regions to a degree whereby radio wave absorption is comparable to that in HII regions.

In Section III of this paper soft photon (radio) bremsstrahlung in electron–atom (e–H and e–He) collisions is computed. This type of emission process has a temperature dependence different from that for electron–ion bremsstrahlung and is negligible at the low temperatures of HI regions. Results for these two processes are presented in a form (atomic units) convenient for comparison. For this reason and for the sake of completeness the features of the electron–ion bremsstrahlung calculation are outlined briefly in Section II. A short discussion of the range of validity of the classical and quantum mechanical expressions for radio Coulomb bremsstrahlung is also given in Section II.

The dependence of the mean absorption coefficient on z, the distance above the galactic plane, is discussed in Section VI. Application to the observed spectra of the cosmic background and of the discrete sources Cas A and Cyg A is outlined in Sections VII and VIII.

II. Coulomb Bremsstrahlung

The general quantum mechanical problem of Coulomb bremsstrahlung is extremely complicated mathematically but has been solved by Sommerfeld (1939), whose work was extended by Elwert (1939, 1948). The later work by Elwert (1948) is especially important, since in it the classical limit to the bremsstrahlung cross section is derived from the quantum mechanical formulation, thus indicating its range of validity. The classical result has been derived by a number of authors (see references in Oster 1961); one finds that the cross section for the emission of a soft bremsstrahlung photon of frequency within $d\omega$ (centred on ω) by an electron of velocity v_e incident on a charge ze is

$$d\sigma_{z(\mathrm{cl})} = \frac{16}{3} z^2 \alpha r_0^2 \frac{d\omega}{\omega} \frac{c^2}{v_e^2} \ln\left(\frac{2mv_e^3}{\Gamma z e^2 \omega}\right). \tag{1}$$

Here α is the fine structure constant, r_0 is the classical electron radius, and $\ln \Gamma = 0 \cdot 5772\ldots$; by "soft" is meant $\hbar\omega \ll \frac{1}{2}mv_e^2$. The bremsstrahlung energy emitted per unit time, volume, and frequency interval for a Maxwellian distribution of electrons incident on ions with number density n_z is

$$dE_z/dt\, d\omega\, dV = n_z \int \hbar\omega v_e (d\sigma_z/d\omega)\, dn_e. \tag{2}$$

It is convenient to now give the result of the integration (2) in atomic units (cf. Bethe and Salpeter 1957), which are functions of the *atomic* constants e, m, and $\hbar$, and in terms of the dimensionless *radiation* constant $\alpha = e^2/\hbar c$. We denote all physical quantities in terms of these units by Greek symbols. Thus, such quantities (and their units) as length (a_0 = Bohr radius), time ($\tau_0 = \hbar^3/me^4$), frequency (τ_0^{-1}), energy ($E_0 = e^2/a_0 = 2$ ry), energy density (E_0/a_0^3), number density (a_0^{-3}), and temperature (E_0/k) are denoted by λ, τ, ω, η, ε, ν, and θ respectively. One then readily obtains for the energy density emission rate per unit frequency (2), using the classical cross section (1)

$$d\varepsilon_{z(\mathrm{cl})}/d\tau\, d\omega = \nu_e \nu_z \epsilon_{z(\mathrm{cl})} ,$$

with

$$\epsilon_{z(\mathrm{cl})} = C_r z^2 \theta^{-\frac{1}{2}} \ln(2^{5/2}\theta^{3/2}/\Gamma^{5/2} z\eta) \tag{3}$$

and

$$C_r = \tfrac{16}{3}(2/\pi)^{\frac{1}{2}}\alpha^3 = 1\cdot 65351\times 10^{-6} .$$

Here η is the photon energy and $\epsilon_{z(\mathrm{cl})}$ is the basic parameter for the process. The advantages of presenting the formula in this notation are that the dependences of the result on density, temperature, and the "order" (third) of the radiative process are more clearly exhibited. Also, one can more readily discuss the range of validity of the formula and compare it with that for other processes (see Section III).

From the work of Elwert (1948) one concludes that the classical cross section (1) is valid as long as the momentum p_e of the incident electron is small compared with the Bohr momentum $p_0\,(= me^2/\hbar)$. When $p_e \gg p_0$ the exact quantum mechanical formula reduces to the Born approximation cross section which, in the low frequency limit, is identical with the classical expression (1) except that the argument of the logarithm is now $4E_e/\hbar\omega$, where E_e is the kinetic energy of the incident electron. Thus, especially for low frequencies where the arguments of the logarithms are large, the domains of applicability of the classical and quantum (Born) approximation formulae are complementary.*

Integrating the Born formula as in (2) we then have

$$\epsilon_{z(\mathrm{q})} = C_r z^2 \theta^{-\frac{1}{2}} \ln(4\theta/\Gamma\eta) . \tag{4}$$

Comparing the two formulae (3) and (4) we see that they are identical for a critical temperature $\theta_c = \frac{1}{2}z^2\Gamma^3 = 2\cdot 825\, z^2$, or for $T_c = 8\cdot 92\times 10^5\, z^2\,{}^\circ\mathrm{K}$. By inadvertently leaving out a factor Γ, Oster (1961) obtained an incorrect value $(5\cdot 5\times 10^5\, z^2)$ for this critical temperature. In summary, for $T < T_c$ the classical formula (3) should be used, while for $T > T_c$ the quantum (Born) expression is more accurate. It might be added that this criterion is independent of photon frequency ω except for the requirement that the photons be "soft". Thus for application to emission and absorption in normal interstellar matter the classical formula is valid; for considerations of high temperature gases, for example, stellar coronae or thermal X-ray sources, the Born formula should be used.

III. Emission in Electron–Atom Collisions

In investigations of the absorption and emission of radiation from interstellar HI regions it is necessary to consider collisions of electrons with hydrogen and helium atoms (and possibly with hydrogen molecules), since the atomic species are more abundant than their ions. Calculation of the electron–atom bremsstrahlung cross section appropriate for our purposes is simplified for two reasons. Firstly, since we

* The statement by Bethe and Salpeter (1957) that "... the Born approximation holds very close to the low-frequency limit even for very low initial momenta ..." is a little misleading and should more properly be stated that the Born formula is only *asymptotic* to the correct result as $\omega \to 0$. For small but finite ω and small p_e the classical expression is more accurate than the Born formula.

are again concerned with collisions of electrons of energies $\sim kT_{\rm II} \sim 1$ eV (HII regions) or $\sim kT_{\rm I} \sim 10^{-2}$ eV (HI regions) and the emission or absorption of radio waves, we can consider only soft photons; thus, the bremsstrahlung cross sections can be found essentially from the cross sections for *elastic* scattering (see below). Secondly, these electron energies correspond to momenta small compared with the Bohr momentum; for such low energy collisions the scattering off atomic systems is isotropic. Also the scattering cross section can be expected to be independent of incident energy. Experimental work described in the article by Moiseiwitsch (see Bates 1962) bears out these latter assertions.

To compute the bremsstrahlung cross sections from the elastic scattering cross sections we make use of the fundamental formula for the probability of emission of a soft photon of frequency within $\mathrm{d}\omega$ accompanying a change $\Delta \mathbf{v}_{\rm e}$ in the velocity of an electron:

$$\mathrm{d}w_{\omega}(\Delta \mathbf{v}_{\rm e}) = \frac{2}{3\pi}\alpha \frac{(\Delta \mathbf{v}_{\rm e})^2}{c^2}\frac{\mathrm{d}\omega}{\omega}; \tag{5}$$

here, again, α is the fine structure constant. The formula (5) can be derived from both classical (cf. Jackson 1962) and quantum electrodynamics (cf. Jauch and Rohrlich 1955). For elastic scattering of incident electrons of velocity $v_{\rm e}$ through an angle $\theta_{\rm sc}$, $(\Delta \mathbf{v}_{\rm e})^2 = 2v_{\rm e}^2(1-\cos\theta_{\rm sc})$. We can then compute the differential bremsstrahlung cross section for electron–atom collisions from the angular differential cross section for elastic scattering:

$$\mathrm{d}\sigma_{\rm a}(\omega) = \int \mathrm{d}\sigma_{\rm el}(\theta_{\rm sc})\, \mathrm{d}w_{\omega}(\theta_{\rm sc})\,. \tag{6}$$

For isotropic scattering $\mathrm{d}\sigma_{\rm el}(\theta) = (\sigma_{\rm el}/4\pi)\mathrm{d}\Omega$ and we get, setting $\sigma_{\rm el} = \pi R_{\rm a}^2$,

$$\mathrm{d}\sigma_{\rm a}(\omega) = \frac{4}{3}\alpha R_{\rm a}^2 \frac{v_{\rm e}^2}{c^2}\frac{\mathrm{d}\omega}{\omega}. \tag{7}$$

It might be mentioned that this method can be used to compute very easily both the classical and quantum mechanical Coulomb bremsstrahlung cross sections to within a factor ~ 1 in the argument of the logarithm (see Jackson 1962).

By integrating the result (7) over a Maxwellian velocity distribution, we get the energy radiated per unit volume, time, and frequency interval in electron–atom collisions. We find, again in atomic units,

$$\mathrm{d}\varepsilon_{\rm a}/\mathrm{d}\tau\,\mathrm{d}\omega = \nu_{\rm e}\nu_{\rm a}\epsilon_{\rm a}\,, \qquad \epsilon_{\rm a} = C_r \rho_{\rm a}^2 \theta^{3/2}\,, \tag{8}$$

where $\rho_{\rm a} = R_{\rm a}/a_0$ and C_r is as in equation (3). Comparing this result with that for Coulomb bremsstrahlung, we see that the latter would always dominate at lower temperatures (remember $\theta = T(^{\circ}\mathrm{K})/316\,000$).

The values of $\rho_{\rm a}$ for e–H, e–He, and e–H_2 collisions may be taken from the data given in the article by Moiseiwitsch (Bates 1962). From the differential cross section measurements of Gilbody, Stebbings, and Fite (Bates 1962, p. 316) at $3\cdot 8$ eV giving $\mathrm{d}\sigma_{\rm el}/\mathrm{d}\Omega \simeq 1\cdot 5\,\pi a_0^2$ (isotropic) for e–H collisions we infer $\rho_{\rm H}^2 \simeq 19$; from their work we also get $\rho_{\rm H_2}^2 \simeq 14$. For e–He collisions, the best data are from electrical conductivity measurements on free electrons in helium gas; these suggest $\rho_{\rm He}^2 \simeq 7$. The accuracy of these parameters for H, H_2, and He is probably about 30%.

We can apply these results to estimate the contribution of e–He bremsstrahlung to the total radio continuum from gaseous nebulae. For a gas in which hydrogen is ionized but helium (about one-sixth as abundant) is neutral and the temperature is about 6000°K, e–He collisions produce only about 0·04% of the total emission, which is due mainly to electron–proton Coulomb bremsstrahlung. Consider, however, an HII region near its outer boundary where hydrogen is about 50% ionized. There are theoretical reasons (Aller, Baker, and Menzel 1939; Burbidge, Gould, and Pottasch 1963; Mathews 1965) and observational evidence (Mezger and Ellis 1968) that the temperature of an HII region increases near its boundary. If $T = 10^4$ °K, e–H collisions produce about 2% of the total emission per unit volume.

IV. Absorption Coefficients

By a straightforward application of the principle of detailed balance, the absorption coefficient or reciprocal mean free path may be found from the emission coefficient. In reciprocal atomic length units (a_0^{-1}) we obtain, for each of the processes described in Sections II and III,

$$\kappa_j = \tfrac{1}{4}\lambda^2\theta^{-1}\nu_e\,\nu_j\,\epsilon_j\,. \tag{9}$$

In equation (9) all quantities are in atomic units; λ is the photon wavelength and j refers to the species scattering the electron. The most significant feature of the result is the dependence on temperature: for absorption during electron–ion Coulomb scattering, essentially $\kappa_z \propto \theta^{-3/2}$; for absorption during electron–atom collisions $\kappa_a \propto \theta^{\frac{1}{2}}$. Coulomb collisions are increasingly important at low temperatures, and this result provides the principal motivation for considering the possibility of absorption in cool interstellar HI regions. Although the ionization fraction in HI regions is small, the low temperature of the region can compensate to give a large value to the absorption coefficient. Cosmic rays can ionize the hydrogen to a degree sufficient to produce this effect.

V. Cosmic Ray Ionization of HI Regions

The work of Gloeckler and Jokipii (1967) in attempting to determine the *de*modulated cosmic ray spectrum has emphasized the importance of low energy (nonrelativistic) cosmic rays and their influence on interstellar gas. Indeed, Balasubrahmanyan *et al.* (1967) and Habing and Pottasch (1967) have estimated the ionization rate constant associated with this cosmic ray flux (see also Spitzer and Tomasko 1968). Unfortunately this rate γ_i (sec^{-1}) cannot be estimated with any accuracy, for most of the ionizations are due to the cosmic rays at the low energy end of the spectrum, which is especially uncertain. To get an idea of the range of values that γ_i can take one can assume that this spectrum is of the form

$$\begin{aligned} dJ &= K_m T^{-m}\,dT \qquad &&\text{for} \qquad T_l < T < T_u\,, \\ &= 0 && \qquad\qquad\text{otherwise,} \end{aligned}$$

where dJ is the differential flux, $T_u =$ (say) 600 MeV, and T_l and m are adjustable parameters. Moreover, let us normalize the spectrum to fit Gloeckler and Jokipii's value $dJ/dT \simeq 10^{-3}$ protons cm^{-2} sec^{-1} sr^{-1} MeV^{-1} at $T_0 = 200$ MeV. The ionization cross section $\sigma_i(T)$ is essentially proportional to T^{-1} so that the ionization rate

$$\gamma_i = 4\pi \int \sigma_i(T)\, dJ$$

is proportional to $m^{-1}(T_0/T_l)^m$ when the spectrum is normalized in the manner described. We see that γ_i is a strong function of the two very uncertain parameters m (index) and T_l (low energy cutoff). Assuming values for m and T_l, γ_i can readily be computed using Bethe's expression for σ_i (cf. Mott and Massey 1965). To account for further ionization by the ejected "secondary" electron one should multiply the result by a factor of about two to get the total ionization rate. These calculations were done for the values $m = 1, 1 \cdot 5, 2$, and $2 \cdot 5$ and $T_0 = 1$, 10, and 100 MeV, which gave values of γ_i ranging from 1×10^{-19} to 1×10^{-15} sec^{-1}, the latter value corresponding to $m = 2 \cdot 5$ and $T_0 = 1$ MeV. Values of T_0 less than 1 MeV would be rather unlikely on the basis of particle energy loss and lifetime considerations. However, the absolute intensity of the cosmic ray flux could indeed be larger, especially in other parts of the Galaxy. Finally, it might be mentioned that for these spectra the cosmic ray energy density varied from $0 \cdot 1$ to 5 eV cm^{-3}.

The hydrogen ionization fraction x_{I} of an HI cloud with total number density n_{H} on the basis of a steady-state balance of cosmic ray ionization and radiative recombination is found from

$$x_{\mathrm{I}}^2 = \gamma_i / n_{\mathrm{H}}\, \alpha_{\mathrm{rec}} \,. \tag{10}$$

At the temperatures ($\simeq 60$°K) of HI regions, $\alpha_{\mathrm{rec}} \simeq 10^{-11}$ cm^3 sec^{-1}, and for a density $n_{\mathrm{H}} = 10$ cm^{-3}, x_{I}^2 ranges from 10^{-9} to 10^{-5} with even larger values (perhaps 10^{-4}) possible (see also Spitzer and Tomasko 1968).

VI. Magnitude and Galactic Latitude Dependence of Absorption in HI and HII Regions

For purposes of application, the numerical value of the absorption coefficient is more appropriately expressed in CGS units. Also, since the temperature of HII regions is typically around 6000°K and that of HI regions around 60°K, we write $T_{6000} = T/6000$ °K; moreover we express the frequency f in megahertz. Then for absorption in Coulomb scattering of electrons by ions of charge ze, the absorption coefficient κ (cm^{-1}) is

$$\kappa_z = 4 \cdot 839 \times 10^{-20}\, n_e\, n_z\, z^2 f^{-2}\, T_{6000}^{-3/2}\, L_{10} \,, \tag{11}$$

where

$$L_{10} = 7 \cdot 363 + \tfrac{3}{2} \log_{10}(T_{6000}) - \log_{10}(zf) \,.$$

Since helium is likely to be either neutral or singly ionized in most HII regions, it will contribute only about 10% of the absorption therein and also in HI regions; therefore we shall ignore its presence.

We shall assume that the interstellar gas is distributed in clouds in which the density is $n_c = 10$ atoms cm^{-3}, and that the matter between the clouds is negligible. The average smeared-out gas density varies with z (= distance above the galactic plane, not to be confused hereinafter with ionic charge), but we shall assume that this is due only to a variation in the number density of *clouds*, i.e. we assume $n_c \neq n_c(z)$. Since the mean smeared-out density of atomic hydrogen in the solar neighbourhood is $\langle n_{\rm H}(z=0)\rangle \equiv \langle n_{\rm H}\rangle_0 = 0{\cdot}7\ \rm cm^{-3}$, this model would correspond to a mean squared number density of atoms in the galactic plane of (see Gould 1966) $\langle n_{\rm H}^2\rangle_0 = 7\ \rm cm^{-6}$. In the galactic plane the amount of gas that is ionized by ultraviolet radiation from hot stars (i.e. the amount that is HII) can be found from observations of high frequency thermal radio emission. In this manner Wilson (1963) derived for the solar neighbourhood a value $\langle n_e n_p\rangle_{\rm II0} = 0{\cdot}08\ \rm cm^{-6}$, assuming $T = 10^4\,°\rm K$; for $T = 6000°\rm K$, this value would be lowered slightly to $\langle n_e n_p\rangle_{\rm II0} \to 0{\cdot}07\ \rm cm^{-6}$. It should be noted that this value is appreciably smaller than that determined previously by Westerhout (1958, see also Gould 1966). The corresponding mean product for HI regions would be $\langle n_e n_p\rangle_{\rm I0} = x_{\rm I}^2\langle n_{\rm H}^2\rangle_0$ and would have a value somewhere between 7×10^{-9} and $7\times10^{-4}\ \rm cm^{-6}$ (see previous section). These parameters would determine the mean absorption coefficient in the galactic plane.

Above the galactic plane we can expect, following the arguments of Gould (1966), that for HII regions the mean absorption coefficient should vary in proportion to the O- and B-star luminosity function $\phi(z)$:

$$\kappa_{\rm II}(z) = \kappa_{\rm II}(0)\,\phi(z)\,, \qquad \phi(z) = \exp(-z^2/h_0^2)\,. \tag{12}$$

(We neglect the dependence on distance from the galactic centre. Within about 3 kpc of the Sun this is not a bad assumption but it would be if we were considering absorption paths nearer the galactic centre (cf. Wilson 1963).) The parameter h_0 is related to the mean absolute height of O- and B-stars by $\pi^{-\frac{1}{2}}h_0 = \langle|z|\rangle = 50$ pc.

For HI regions the mean absorption coefficient should vary essentially as the mean smeared-out gas density $\langle n_{\rm H}(z)\rangle \equiv a(z)\langle n_{\rm H}(z=0)\rangle$:

$$\kappa_{\rm I}(z) = \kappa_{\rm I}(0)\,a(z)\,. \tag{13}$$

(We neglect the dependence of $a(\mathbf{r})$ on distance from the galactic centre. Again, this would be a bad assumption for positions nearer the centre.) The function $a(z)$ can be approximated by (see Schmidt 1957; Gould, Gold, and Salpeter 1963)

$$\left.\begin{aligned} a(z) &\simeq \exp(-z^2/h_{\rm g}^2) && z < 110\ \text{pc}\,,\\ &\simeq 1{\cdot}93\exp(-z/h_{\rm e}) && z > 110\ \text{pc}\,, \end{aligned}\right\} \tag{14}$$

where $h_{\rm g} = 130$ pc, $h_{\rm e} = 80$ pc.

For a radio source at a galactic latitude $b_{\rm s}$ and at a distance $z_{\rm s}$ above the plane, the absorption optical depth would be (since we are situated approximately in the plane)

$$\tau(b_{\rm s}, z_{\rm s}) = \operatorname{cosec} b_{\rm s}\int_0^{z_{\rm s}}\kappa(z)\,{\rm d}z = \operatorname{cosec} b_{\rm s}\,\kappa(0)\,I(z_{\rm s})\,, \tag{15}$$

where the integral I is, for HI or HII absorption,

$$\begin{bmatrix} I_{\rm I}(z_{\rm s}) \\ I_{\rm II}(z_{\rm s}) \end{bmatrix} = \int_0^{z_{\rm s}} \begin{bmatrix} a(z) \\ \phi(z) \end{bmatrix} {\rm d}z\,. \tag{16}$$

The functions I_{I} and I_{II} are plotted in Figure 1. For sources outside the galactic disk the asymptotic values $I_{\mathrm{I}}(\infty) = 127 \cdot 6$ pc and $I_{\mathrm{II}}(\infty) = 78 \cdot 5$ pc are appropriate (in these values only the first two digits can be considered significant).

It should be emphasized that the expressions (15) and (16) give a measure of the *average* absorption optical depth. In directions where there is more or less absorbing ionized matter the optical depth could be quite different.

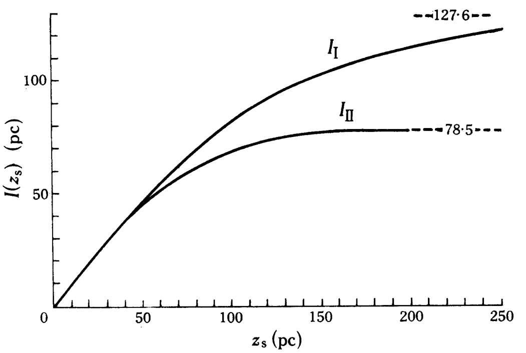

Fig. 1.—Graphs of the functions $I_{\mathrm{I}}(z_{\mathrm{s}})$ and $I_{\mathrm{II}}(z_{\mathrm{s}})$.

At this point we might compare the relative magnitudes of the absorption in HI and HII regions. For absorption through the galactic disk the ratio is, on the average,

$$\frac{\tau_{\mathrm{I}}}{\tau_{\mathrm{II}}} = \frac{\kappa_{\mathrm{I}}(0)}{\kappa_{\mathrm{II}}(0)} \frac{I_{\mathrm{I}}(\infty)}{I_{\mathrm{II}}(\infty)}$$

$$= \frac{\langle n_{\mathrm{e}}\, n_{\mathrm{p}} \rangle_{\mathrm{I0}}}{\langle n_{\mathrm{e}}\, n_{\mathrm{p}} \rangle_{\mathrm{II0}}} \left(\frac{T_{\mathrm{II}}}{T_{\mathrm{I}}}\right)^{3/2} \frac{L_{10}(\mathrm{I})}{L_{10}(\mathrm{II})} \frac{I_{\mathrm{I}}(\infty)}{I_{\mathrm{II}}(\infty)}. \tag{17}$$

Taking $T_{\mathrm{II}}/T_{\mathrm{I}} = 10^2$ and $f = 1$ MHz, we obtain $\tau_{\mathrm{I}}/\tau_{\mathrm{II}} \simeq 1$ for $x_{\mathrm{I}}^2 = 10^{-5}$; this would indicate a moderate amount of absorption in HI regions. However, x_{I}^2 could easily be 10^{-4}, which would give $\tau_{\mathrm{I}}/\tau_{\mathrm{II}} \simeq 10$. The main point to be made here is that the effects of the low energy cosmic rays may possibly be observed indirectly through radio observations of interstellar absorption. We shall discuss the present observational evidence pertaining to this problem in the following sections.

VII. Observation of Absorption in Cosmic Background Spectrum

Ellis and Hamilton (1966, see also Hoyle and Ellis 1963) have analyzed their low frequency ($2 \cdot 1$, $4 \cdot 7$, and $9 \cdot 6$ MHz) surveys of the southern sky, interpreting the turnover as an effect of absorption in the ionized interstellar gas. The angular resolution in the $2 \cdot 1$ MHz survey, where absorption effects are greatest, was 8°. Assuming a temperature of 10^4 °K and a total thickness $2L_0 = 400$ pc for the ionized layer in the galactic disk, Ellis and Hamilton determined the r.m.s. electron density therein to range from $\langle n_{\mathrm{e}}^2 \rangle^{\frac{1}{2}} = 0 \cdot 2$–$0 \cdot 27$ cm^{-3} for galactic latitudes $b^{\mathrm{II}} = 60°$ and 5° respectively. On the basis of the discussion in the previous section, it would seem that the half-thickness chosen ($L_0 = 200$ pc) may be too large by about a factor of two. However, the correction to a temperature of 6000°K would essentially compensate this. Thus, the mean squared electron density $\langle n_e^2 \rangle = 0\cdot04$–$0\cdot07\ \mathrm{cm}^{-6}$ can be compared with the value (see previous section) $0\cdot07\ \mathrm{cm}^{-6}$ determined from radio *emission* measurements.

The smaller value determined from absorption measurements at the higher latitude may be due to the irregular behaviour of the absorption coefficient on direction in the sky. Owing to the finite angular resolution of the radio receiver, there would be a tendency to record radiation coming through the galactic disk along paths where the absorption optical depth τ is a minimum. The effect would be more pronounced if τ attained values equal to or greater than unity in some directions in the receiver beam. This "fluctuation" effect would be minimized when the absorption path length is larger, as it would be in the observations at $b^{\mathrm{II}} = 5°$. At any rate, it is understandable how $\langle n_e^2 \rangle$ could be *underestimated* due to this effect, and this may explain the smaller value determined for $b^{\mathrm{II}} = 60°$. Essentially this same effect has been suggested by G. Field (personal communication) to come into play for the attenuation of cosmic X-rays in the galactic disk. X-ray absorption would, of course, occur mainly in interstellar HI regions that have a more regular distribution than HII regions. In fact, if radio absorption is indeed taking place primarily in HII regions the effect would seem to be greater than that indicated by the present observations. The question of whether radio waves are absorbed mainly in HI or HII regions could undoubtedly be answered by low frequency observations with better resolution.

The southern survey work by Shain and his collaborators (Shain, Komesaroff, and Higgins 1961) should be mentioned here. This was a survey at $19\cdot7$ MHz with a beamwidth of $1°\cdot4$. Generally, these observations showed a definite correlation in the positions of radio "darkness" with locations of optically observed HII regions. However, in some places where the radio observations indicate absorption, e.g. at longitude $l^{\mathrm{I}} = 5°$, no HII regions are observed. Although Komesaroff (1961) has interpreted this as being due to a lack of emission from this direction, the overall evidence presented by him could rather be taken to suggest that the radio darkness is indeed an absorption effect. In fact, consultation of the results of the 21 cm surveys (Schmidt 1957) reveals the presence of a dense HI concentration in this direction at a distance of about 2 kpc.

VIII. Absorption Effects in Spectra of Discrete Sources

The difficulty mentioned at the end of the previous section is overcome if one considers the attenuation of the radio flux from point sources. For although the path to the source may happen to pass an unusually large or small amount of ionized matter, it is essentially a randomly chosen path. In fact, if data were available on attenuation in the spectra of a number of sources, the mean squared density of ionized matter in the solar neighborhood could be determined accurately.

Several source spectra show absorption effects at low frequency. Cyg A has a low frequency turnover, but the measurements are not accurate enough to determine a good value of the absorption optical depth τ_f; moreover, it may be that the absorption is taking place within the source. The best data exist for Cas A which has a very accurately determined spectrum even at low frequencies. Cas A is at a distance of 3400 pc (Minkowski 1964); its galactic latitude and longitude are $b^{\mathrm{II}} = -2\cdot1°$, $l^{\mathrm{II}} = 111\cdot7°$, so it lies at a distance $z_{\mathrm{s}} = 124$ pc below the galactic plane. Since it does lie near the plane the absorption path traverses a large amount of interstellar matter.

From the observational data of Bridle (1967) and Parker (1968) a fairly accurate value of τ_{f} can be determined for the Cas A spectrum. Assuming a form of the spectrum

$$S(f) = Kf^{-\alpha}\exp(-\tau_{\mathrm{f}}), \tag{18}$$

with α ($= 0\cdot770 \pm 0\cdot006$) determined from the spectrum at higher frequencies where $\tau_{\mathrm{f}} \ll 1$, τ_{f} can be found from the observed spectrum at low frequencies. Here the basic assumption is that the *emission* spectrum of the source retains its power-law form at the low frequency end. Parker (1968) has summarized the data for Cas A and made time evolution corrections to a common epoch. From the accurately determined observed fluxes at $10\cdot05$ and $152\cdot0$ MHz, the value $\tau(10\cdot05\ \mathrm{MHz}) = 1\cdot2_1$ is readily obtained (τ at $152\cdot0$ MHz is very small; about $0\cdot005$). From equations (11) and (15), assuming $T = 6000°$K (absorption only in HII regions), the mean value of n_{e}^2 in the galactic plane between the Sun and the point above Cas A can be determined. More precisely the value determined is the mean of $n_{\mathrm{e}}n_z z^2$, summed over the ions z (mainly hydrogen and helium). In this manner, taking $I_{\mathrm{II}} = 75$ pc (see Fig. 1), we obtain

$$\sum_z \langle n_{\mathrm{e}} n_z z^2 \rangle_0 = 0\cdot06_3\,\mathrm{cm}^{-6}, \tag{19}$$

a value in agreement with that found from thermal radio emission (see Section VI) and with that found from the cosmic background spectra *at low latitudes* (Ellis and Hamilton 1966; see Section VII).

Because of the agreement mentioned above, there appears to be no real need for hypothesizing absorption in HI regions although an appreciable attenuation therein still cannot be ruled out. Low frequency, high resolution studies of the cosmic background distribution should settle this question, however.

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