

4. THEORETICAL DESCRIPTION OF THE INTERSTELLAR MEDIUM

Introductory Report

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1. Introduction

As excellent reviews have appeared recently (Spitzer, 1968a; Pikel'ner, 1968a) we shall consider here only some special topics where progress is being made – thermal properties of H I and H II clouds, and shock waves. The components of the medium which are dynamically important are neutral atoms, charged ions and electrons, grains, cosmic rays, magnetic field, and light. In Table I we indicate the degree of

TABLE I
Momentum coupling between components
(S = Strong, M = Medium, W = Weak)

Component	Atoms	Ions	Grains	Cosmic rays	Light	Magnetic field
Atoms	S	S	M	W	W	W
Ions	S	S	M	W	W	S
Grains	M	M	W	W	M	M
Cosmic rays	W	W	W	W	W	S
Light	W	W	M	W	W	W
Magnetic field	W	S	M	S	W	-

momentum coupling among these components. The main mass of the medium, in ions and atoms, is well coupled through collisions, with a mean free path $\lambda \approx 10^{-3}$ pc, and may usually be treated hydrodynamically. Parker (1965) has shown that for many purposes the cosmic-rays can be considered as a separate gas, strongly coupled to the magnetic field. The coupling of the atom-ion gas to the magnetic field is through the magnetic forces on the ions, which transmit the forces to the neutral atoms by collisions unless the differential stresses on the two components are unusually large. Grains are not coupled strongly to any other component, as their stopping distances are of the order of parsecs and their gyro-periods of the order of 10^4 yr. The coupling of grains to radiation pressure can lead to interesting differential effects. Light is important in carrying away thermal energy but does not critically affect the momentum balance of the gas. A possible exception is the Ly- α radiation field discovered by Kurt and Sunyaev (1967) which has a pressure of 10^{-13} dyne cm^{-2} . If this phenomenon is common in interstellar space, it could have effects.

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Most work has been on the large-scale motion of the atom-ion gas, includes magnetic stresses, and assumes that the magnetic field is frozen in. Such studies include equilibrium configuration in the galactic magnetic field, cloud collisions, expansions of H II regions and supernova shells, and perturbation calculations aimed at discovering instabilities (gravitational, thermal, hydromagnetic). Many of these calculations have paid insufficient attention to the effects of cosmic-ray pressure. Similarly, insufficient work has been done on the dynamical effects of radiation pressure on grains.

2. Thermal Properties of H I Regions and Formation of Clouds

Spitzer calculated thermal equilibrium conditions in the 1940's. He concluded that the main heating is due to photo-electrons ejected from carbon atoms, and that the main cooling is by infrared emission following collisional excitation of the fine-structure levels of various ions. Spitzer's equilibrium temperatures (≈ 20 K) are too low to explain the 21-cm result that $\langle T^{-1} \rangle^{-1} \approx 100$ K, therefore Kahn (1955) suggested that the temperature may be maintained by shocks associated with supersonic cloud collisions. We shall see below that this mechanism appears inadequate to maintain the temperature. Hayakawa *et al.* (1961) showed that if low-energy (1 to 100 MeV) cosmic rays are present in sufficient numbers, the heating by the secondary electrons they produce would explain the temperatures of clouds. Field (1962) showed that the observed cosmic rays with $E > 1$ GeV would heat any low-density gas to 10^4 K or above, so that there could be pressure equilibrium between two phases: a rarefied, hot intercloud medium and dense, cool clouds. This explanation of the stability of clouds had been suggested by Spitzer (1951). Pikel'ner (1967) revived the discussion, showing that the heating by cosmic rays and the partial ionization of otherwise neutral H I explains various observations, including the existence of hot H I (21-cm) regions, Faraday rotation, and long-wavelength radio absorption. Pikel'ner included the heating due to Coulomb collisions of cosmic rays with ambient thermal electrons and discussed extensively the possibility and the consequences of the existence of three thermal phases in the interstellar gas: a stable, hot phase of low density ($n \lesssim 10^{-1} \text{ cm}^{-3}$, $T \gtrsim 6000$ K), a stable, cool phase of high density ($n \gtrsim 1 \text{ cm}^{-3}$, $T \lesssim 300$ K), and an intermediate, thermally unstable phase in which $dp/dq < 0$.

Spitzer and Tomasko (1968) investigated the efficiency of cosmic-ray heating by calculating the energy put into the secondary electrons and subtracting the energy lost by the secondaries through excitation of optical transitions. Later, Spitzer and Scott (1969) (to be called SS) extended these calculations to include elastic collisions between the secondary electrons and the ambient thermal electrons, and Coulomb collisions between cosmic rays and ambient electrons.

Field *et al.* (1969) (FGH) and Goldsmith *et al.* (1969) using a different method, also included the efficiency of the secondaries in their calculations and calculated equilibrium temperatures (heating=cooling) using cooling processes and cross-sections somewhat different from those of Pikel'ner and of Spitzer and Tomasko.

They found results qualitatively similar to Pikel'ner's and showed that Pikel'ner's model also explains the electron density observed by pulsar dispersion measurements. They concluded that if agreement with cloud temperatures is to be obtained, depletion of cooling elements on grains is to be postulated. The latter effect might then explain the low abundance of sodium and calcium as deduced from interstellar absorption-line observations. Hjellming *et al.*, (1969) (HGG), using somewhat different cooling processes, applied the efficiency calculations of Spitzer and Scott and obtained thermal equilibria in agreement with FGH. Silk and Werner (1969) have suggested that ionization by X-rays < 1 keV could supplement or replace low-energy cosmic rays as a heating and ionizing agent.

Figure 1 shows the amount of energy delivered to the gas and the total number of electrons produced per primary ionization as a function of n_e/n . The values of SS are the most reliable. Figure 2 shows the equilibrium temperatures derived. The results have all been normalized to a particular value of the primary ionization rate ζ of hydrogen found by FGH to be in good agreement with observation. (FGH showed that all processes are proportional to this function of the cosmic-ray spectrum, so that the detailed form of the latter, which is unknown, is irrelevant). Three of the calculations agree fairly well in view of the somewhat different assumptions about ζ_{eff} and

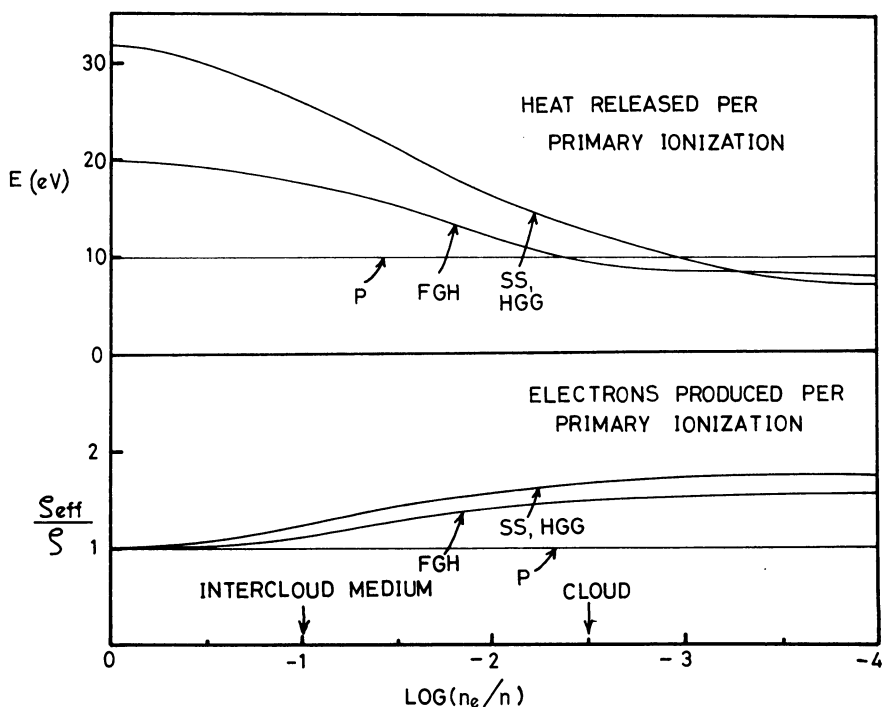


Fig 1. Heat released and number of electrons produced per primary ionization of a neutral hydrogen atom. P = Pikel'ner (1967), FGH = Field, Goldsmith, and Habing (1969), SS = Spitzer and Scott (1969), HGG = Hjellming, Gordon, and Gordon (1969). ζ is the rate of primary ionization per neutral hydrogen atom; ζ_{eff} includes in addition ionization by energetic secondary electrons.

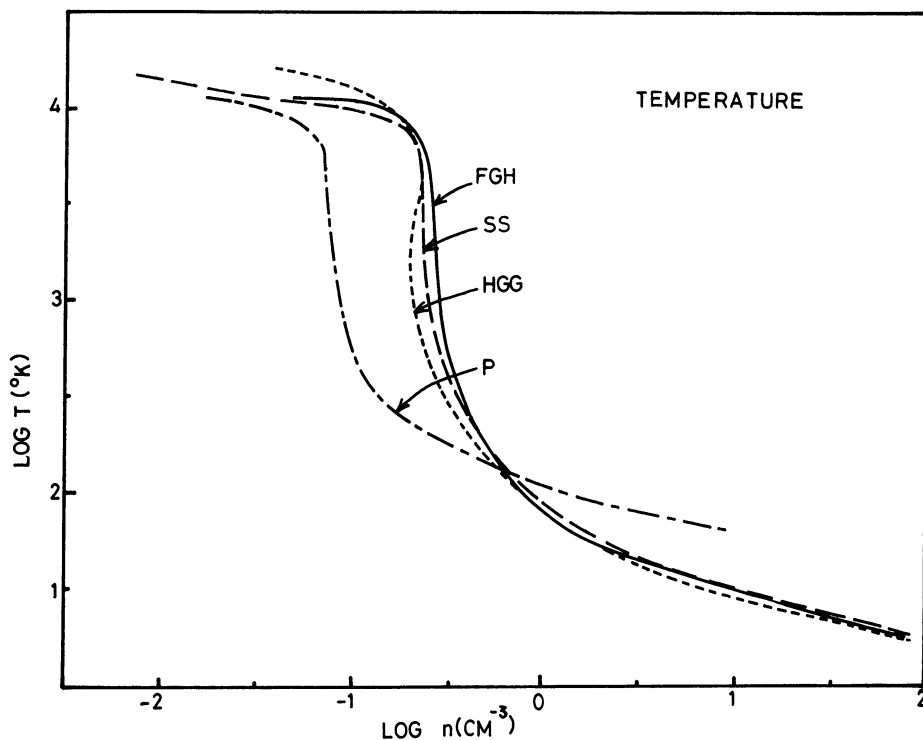


Fig. 2. Temperature for which the cosmic ray heating equals the cooling by inelastic, thermal collisions as a function of density n . Results of various authors are normalized to one value of $\zeta (= 4 \times 10^{-16} \text{ sec}^{-1})$. See Figure 1 for explanation on the names of the authors.

ΔE (Figure 1), and about the cooling mechanisms. The results of Pikel'ner agree qualitatively, and the quantitative differences may probably be explained by the fact that certain of his cross sections were only estimated.

The curves can be understood qualitatively as follow. As n drops, the cooling ($\propto n^2$) decreases faster than the heating ($\propto n$) and T , which influences only the cooling, must rise to compensate for the difference. In the range 10^3 to 5×10^3 K, fine-structure cooling is relatively insensitive to T , and T must rise rapidly to about 10^4 K, where Ly- α cooling becomes important. FGH included the $^3P-^1D$ transition of OI at 1.96 eV ($E/k=22\,800$ K), while SS and HGG included the $^3P-^1D$ transition of NII at 1.90 eV ($E/k=22\,100$ K) as well. The conclusion that NII should not contribute significantly (FGH) is strengthened by the agreement of HGG, SS, and FGH. All authors omitted what could possibly be an important mechanism for cooling the intercloud medium, i.e., electron excitation of a $^4F_{9/2}$ -level of Fe^+ with $E/k=2700$ K. Pottasch (1968) estimates $\Omega=15$ for this transition, and if $\text{Fe}/\text{O}=0.1$ (Pottasch, 1963; Swings, 1965; Gary *et al.*, 1969), rather than the 0.01 assumed by FGH, we find that this transition should be quite important in the range of several thousand degrees and would thus depress the temperature of the intercloud medium (see below).

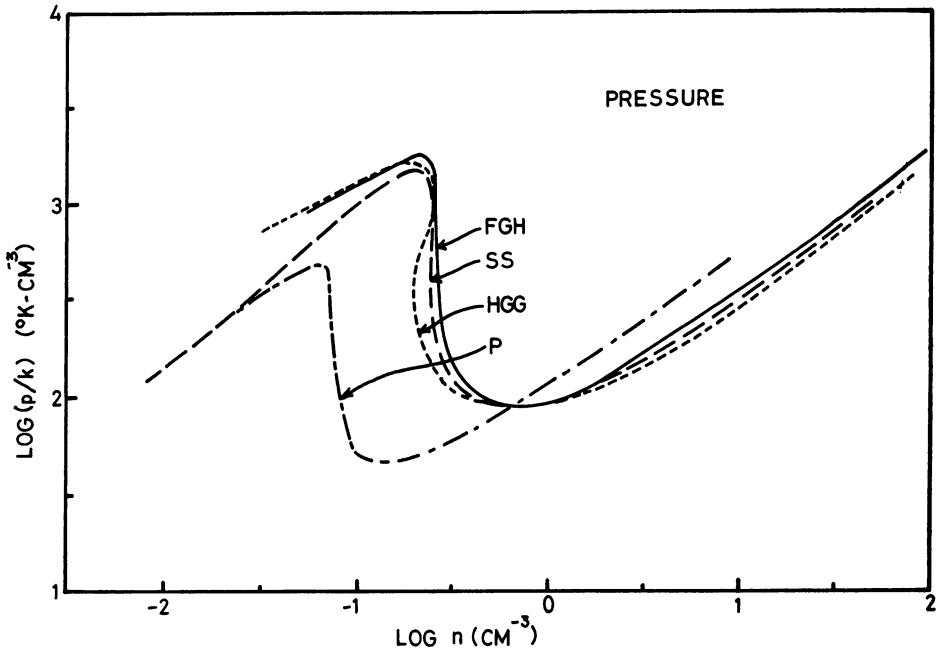


Fig. 3. The pressure of the interstellar medium as a function of total nuclear density. See Figures 1 and 2 for explanation.

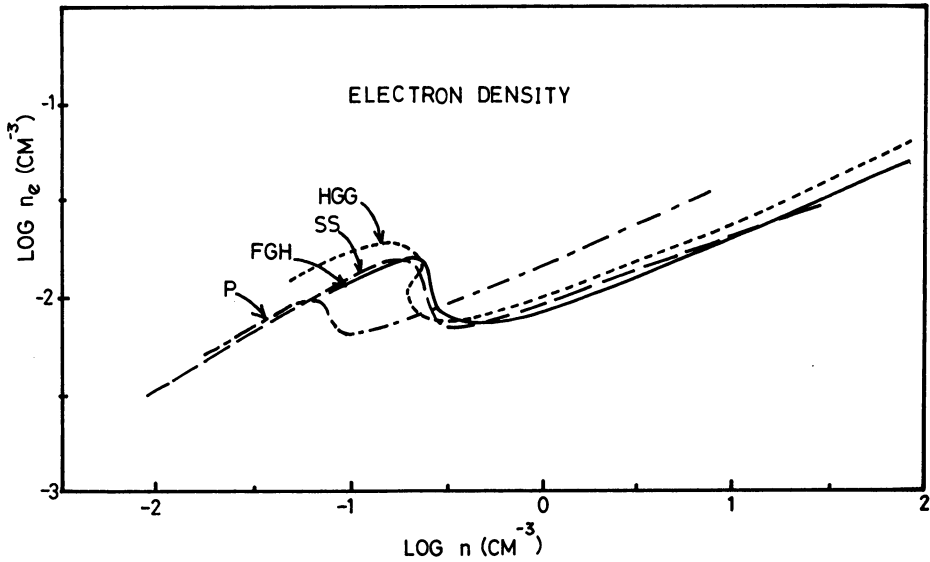


Fig. 4. Electron density in the interstellar medium as a function of total nuclear density. See Figures 1 and 2 for explanation.

Figure 3 shows the gas pressure $p/k=(n+n_e)T$; and Figure 4, the electron density n_e . The latter is about 10^{-2} cm^{-3} , in agreement with observations of pulsar dispersion and low-frequency radio absorption. In every case there is a pressure maximum; if gas in the high-temperature phase at the left of the maximum were to be slowly compressed, no equilibrium could be found beyond the critical point whose parameters (p_c , n_c , T_c) are listed in Table II. FGH interpret this to mean that in our local neigh-

TABLE II
Parameters of critical point
($\zeta = 4 \times 10^{-16} \text{ sec}^{-1}$)

Author	$n_c(\text{cm}^{-3})$	$T_c(\text{K})$	$p_c(\text{cm}^{-3} \text{K})$	$n_{ec}(\text{cm}^{-3})$
P	0.062	7000	430	0.009
FGH	0.21	7800	1800	0.016
SS	0.21	7000	1500	0.015
HGG	0.19	9000	1700	0.016

References: P = Pikel'ner (1967); FGH = Field, Goldsmith, and Habing (1969); SS = Spitzer and Scott (1969); HGG = Hjellming, Gordon, and Gordon (1969).

borhood gas must occur in the low-temperature phase. This gas forms the cool clouds that produce at 21 cm all of the absorption lines and part of the emission lines. The argument is as follows. Consider a particular tube of force, which on the Parker (1968a) model of the magnetic field (see below) loops out of the plane. Along the tube we must have hydrostatic equilibrium in the galactic gravitational field, K_z . At large z , the density is low and the gas is in the hot stable phase. If one can integrate the hydrostatic equilibrium equation to the lowest point on the flux tube without exceeding p_c , the entire tube could be filled with hot gas, but if the weight of the gas is too great, one will reach p_c first. Since $dp/dz < 0$ in hydrostatic equilibrium, a phase transition to the stable cool phase must occur and the gas will form normal interstellar clouds. FGH showed that the amount of gas observed in our local neighborhood is too great to be accommodated in the hot stable phase, so that clouds must occur. Pikel'ner (1967) also took this point of view, and pointed out that if the total density between arms is much lower, there may be no clouds there.

This model predicts the existence of a rarefied, hot HI intercloud medium. Recent observations of the 21-cm line in absorption confirm this prediction and give $T > 1000 \text{ K}$ (Clark 1965; Radhakrishnan and Murray, 1969). Observations of the same line in emission also indicate the existence of hot HI gas (Heiles, 1967a; Habing, 1969; Grahl *et al.*, 1969). However, according to Heiles (1967a) the velocity dispersion of the line profile indicates $T \leq 4000 \text{ K}$, which is much less than that predicted by theory for the stable hot phase ($T > 8000 \text{ K}$, see Table II). Mechanisms, such as iron cooling, which might reduce T_c (the temperature at the critical point), are therefore of particular interest. HGG suggest that a way out of this problem is to take the intercloud

medium to be the intermediate unstable phase. We have two objections to this procedure. Away from the plane, where K_z is significant, hydrostatic equilibrium demands that $dp/dz < 0$, and since $(dp/d\rho) < 0$ in the intermediate phase, a region containing this phase would have $(d\rho/dz) > 0$, which is Rayleigh-Taylor unstable (Layzer, 1967) on a time scale $(dK_z/dz)^{-1/2} = 1.3 \times 10^7$ yr. Even if z is small, and $K_z \propto z$, gas in the intermediate phase is thermally unstable (Field, 1965). According to Pikel'ner (1967), intermediate-phase gas may be present while the long wavelengths required to make interstellar clouds are developing, since the time scale for development of long wavelengths is about $\lambda/c \gg$ thermal time scale (c is the sound speed). However, this overlooks the fact that shorter wavelengths will develop faster, namely on the thermal time scale, and it has been shown (Goldsmith, 1969) that this leads to a transition to stable phases in about 10^6 yr.

We now consider an alternative model proposed by Parker (1966). Parker (1968b, 1969) gives references to the many papers by him and Lerche on this model. In 1960 Hoyle and Ireland pointed out that differential rotation of the Galaxy will double the disk magnetic field in about 4×10^8 yr. Since $B^2/(8\pi\rho z)$ is comparable with K_z , the support of the gas against gravitation is partly due to magnetic forces, and the increase of B will lead to an increase in z . They argued that this process will not be uniform but will proceed fastest where the density is lowest. Further, the process is unstable (Hoyle and Harwit, 1962); since the gas will run from the elevated regions into the depressed regions to form clouds, the elevated regions will be able to expand even faster.

Parker also shows mathematically that, even if one starts with a relatively smooth field, constant in time, and parallel to the disk, the system is unstable toward perturbations in which the field acquires a sinusoidal z -component. The kinetic energy of the unstable gas derives from three sources: the gravitational energy of the gas sliding down the line of force, the magnetic energy released by the upwelling field lines, and the cosmic-ray energy released by the expansion of the field. The effect of gas pressure depends on the sign of $(\gamma - 1)$, where $\gamma = d(\log p)/d(\log \rho)$ characterizes the thermal behavior of the gas. $\gamma - 1 < 0$ implies destabilization.

Parker's results for the case $K_z = \text{constant}$ are easiest to relate, and he showed that they are approximately correct for the more realistic case in which $K_z \propto z$. He finds the following instability criterion:

$$\left[\frac{1}{2}\alpha + \beta + (\alpha + \beta)^2\right] - [(\gamma - 1)(1 + \frac{3}{2}\alpha + \beta)] - 2\alpha\gamma k^2 L^2 > 0, \quad (1)$$

where $\alpha = \text{magnetic pressure/gas pressure}$, $\beta = \text{cosmic-ray pressure/gas pressure}$, $\gamma = d(\log p)/d(\log \rho)$, $L = \text{scale height} = (dp/d\rho)(1 + \alpha + \beta)/|K_z|$, and $k = \text{horizontal wave number of the disturbance along the magnetic field}$. The first term shows that magnetic and cosmic-ray pressures are always destabilizing. The second shows that gas pressure is also destabilizing if $\gamma < 1$; otherwise it is stabilizing, because a hot gas will resist collapse down the field line more than a cool one. The last term shows that if the gas is thermally stable ($\gamma > 0$) short wavelengths tend to be more stable because of the restoring forces associated with strong gas-pressure gradients in short-wavelength

disturbances. If, on the other hand, the gas is thermally unstable the effect of gas pressure is destabilizing, whatever the wavelength.

Parker has discussed a model in which $\alpha = \beta$ (equipartition between cosmic rays and magnetic field, $\beta \approx 3.5 \mu\text{G}$). This model may be close to reality because Anand *et al.* (1968) have demonstrated that proper accounting for solar modulation of energetic electrons lowers the field estimated from the synchrotron emission of the Galaxy to $6 \mu\text{G}$ (from the old value, $20 \mu\text{G}$), in harmony with the values 3 to $6 \mu\text{G}$ from Zeeman and Faraday rotation measurements. Equation (1) then reveals that the medium is unstable for the conditions shown in Table III. Here k_c is a critical wave number

TABLE III
Conditions for Parker instability
($K_z = \text{const.}$, cosmic-ray = magnetic pressure)

$\gamma = d \log p / d \log \rho$	α	k
$\gamma \leq 0$	All	All
$0 < \gamma < 1$	All	$k < k_c$
$1 < \gamma$	$\alpha > \frac{2}{3}(\gamma - 1)$	$k < k_c$

obtained by replacing in Equation (1) the sign $>$ by $=$. We may apply to Parker's model the calculation of pressure given previously, using the results of FGH, SS, and HGG as an example. These curves are consistent with various observations to a factor of 2, provided one takes $B \approx 3 \mu\text{G}$ in the intercloud medium. From Figure 3 we see that $\gamma > 0$ above $n = 0.8$ nuclei cm^{-3} ($n_{\text{H}} = 0.7 \text{ cm}^{-3}$ when 10 percent helium is taken into account). Since this is close to the observed mean density of the gas, one might consider a model in which all the gas is in the stable cool phase, close to the pressure minimum, with magnetic pressure (5×10^{-13} dyne cm^{-2}) much greater than gas pressure (1.3×10^{-14} dyne cm^{-2}). This would correspond to $\alpha = 40$. Such a disk can be supported in equilibrium by the pressure of magnetic field and cosmic rays, as Parker showed, but Table III shows that the disk is unstable for all wavelengths $\lambda > \lambda_c \approx 2\pi L \sqrt{(2\alpha\gamma)/(1+2\alpha)}$. The value of γ is quite sensitive to the exact value of n , but if we take $\gamma = 0.1$ as an example, $\lambda_c = 33$ pc.

Parker analyzed the situation in which γ is constant. However, as the density changes the pressure will change in a much more complicated way. In fact, since the thermal relaxation time scale is short (10^5 to 10^6 yr) compared to the dynamical time scale (10^7 yr), the gas will either remain close to a stable phase, or, if led into an unstable regime, make a rapid phase transition out of it. While no detailed calculations have been done, one can imagine the following sequence of events as the dynamical instability develops.

Suppose we start with a horizontal flux tube containing a uniform gas in conditions close to that of the pressure minimum in Figure 3. If the flux tube begins to develop toward a configuration like that in Figure 5, its volume will expand. Since the gas cannot leave the tube, the mean density has to drop. This drop will bring the gas in the

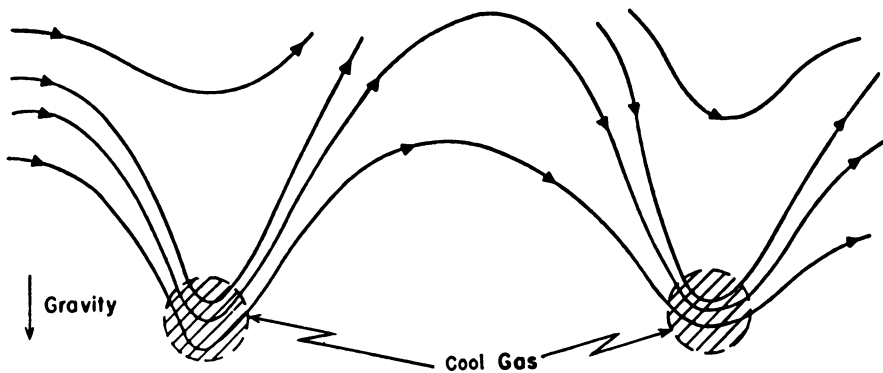


Fig. 5. Sketch of the local state of the lines of force of the interstellar magnetic field and interstellar gas-cloud configuration resulting from the intrinsic instability of a large-scale field along the galactic disk or arm when confined by the weight of the gas (Parker, 1966).

unstable phase. One may suppose that some part of the gas will expand and some part will contract. The contracting gas loses buoyancy and runs down the flux tube. In the expanding gas thermal instability develops rapidly (10^6 yr). Regions less dense than average (presumably those near the top of the flux tube) are below the equilibrium line of Figure 3, and gain heat from the cosmic rays faster than they can radiate it. This increases the pressure and the rarefied gas expands. If one assumes that the gas, which started contracting at the beginning of the expansion occupied about half the volume of the flux tube the density in the hot gas can only fall by a factor of about 2. The hot gas will therefore continue to increase in pressure, and stabilize only when it has reached the hot stable phase at $n \approx 0.2 \text{ cm}^{-3}$, $p \approx p_c = 2.5 \times 10^{-13} \text{ dyne cm}^{-2}$. As the flux tube continues to rise upward, hydrostatic equilibrium gives $p = p_0 \exp(-z/L)$ in the approximation $K_z = \text{constant}$. Here $L [= (dp/dQ)/K_z \approx 100 \text{ pc}]$ is the scale height in the hot stable phase. It seems likely that p_0 , the pressure at the bottom of the field line, remains close to p_c . This pressure, according to Figure 3, requires $n = 100 \text{ cm}^{-3}$, $T = 18 \text{ K}$. While some clouds do have such characteristics, notably those that stand out in 21-cm absorption observations, most clouds are probably hotter and less dense, say 120 K and 15 cm^{-3} . FGH showed that depletion of cooling elements by accretion on grains can account for the increased temperature. A theory of depletion by preferential sticking of charged species with high binding energy to a dielectric grain shows that this takes about 10^8 yr in a cloud having densities of the order of 10 cm^{-3} (Field and Zachariades, 1969).

The final appearance of the system is as shown in Figure 5 (Parker, 1966). The cloud is supported against galactic gravitation by magnetic stresses much as are solar filaments in the theory of Kippenhahn and Schlüter (1957). The detailed magnetic configuration within the cloud is of interest. Pikel'ner (1967) argued that the curvature of the field lines is small, the radius of curvature being some 150 pc . However, if his argument is modified for smaller B and larger density, we obtain a radius of curvature of only a few pc, and this seems too small for the following reason. The scale height of

the cool stable phase in the galactic gravitational field at $z = 100$ pc is only 1 pc, and in the 10 pc radius of a cloud, one would encounter a Δz of about 3 scale heights. The intercloud pressure would then be only 1.2×10^{-14} dyne cm^{-2} , $n_{\text{H}} < 10^{-2}$ cm^{-3} , and $n_{\text{e}} < 5 \times 10^{-3}$ cm^{-3} . There would be significant disagreement with pulsar observations ($n_{\text{e}} = 2 \times 10^{-2}$ cm^{-3}) and observations of hot H I regions ($n_{\text{H}} \approx 0.2$ cm^{-3}) which agree with a model in which $p \approx p_c$.

Finally, we note that the presence of regimes where thermal instability occurs adds to the Parker instability by making $\gamma < 0$ in Equation (1). In the final state, however, if one has only stable phases, the effect of gas pressure is stabilizing. Not only is $\gamma \approx 1$ in both phases but also the gas pressure is no longer completely trivial: $\alpha = 2$ rather than 40. In Equation (1) this leads to $\lambda_c = 300$ pc, which is considerably larger than the mean length of a field line between clouds (≈ 120 pc); so the intercloud medium is stable against further activity.

Instreaming of gas has been observed at both poles of the Galaxy (Dieter, 1969). Dieter (1964) has suggested that this could be cool material condensing from hot gas at high altitude, perhaps by thermal instability. Within the framework of Parker's theory, such gas would stream down along the lines of force looping out of the disk, much like 'coronal rain' in the solar atmosphere. If the gas is much heavier than its surroundings, it will acquire free-fall velocity equal to 10, 20, 30, and 50 km sec^{-1} for origin at 100, 200, 400, and 1000 pc, respectively. As these velocities are of the order of those observed at the pole, the mechanism deserves further attention (Heiles, 1967b). The relative gas density $n(z)/n(0)$ is 0.6, 0.15, and 0.004 at $z = 100, 200,$ and 400 pc, respectively, and is negligible at larger z even if T is as high as $T_c = 9000$ K. Two mechanisms could initiate the infall. Parker (1965) has argued that the slow build-up of cosmic-ray pressure due to constant injection can initiate a local rapid expansion of the field. This could cause the relative density at the top of the field line to drop so low (≈ 0.05) that a stable thermal equilibrium is impossible, and thermal instability commences. Alternatively, a finite amplitude disturbance such as a shock wave could propagate along the field line from a supernova, for example. In this case Goldsmith (1969) has shown that a phase transition can be induced (essentially because $p > p_c$), with the result that clumps of cool gas fall down the field line as required. Such mechanisms might account for activity originating up to 400 pc altitude ($v = 30$ km sec^{-1}), but it seems unlikely that activity at 1000 pc (50 km sec^{-1}) would be observable because of the negligible density there on the present model. Of course a hot corona (Spitzer, 1956) would provide the density, but at present there are doubts about the stability of a phase at 10^6 K. In the present picture, gas is not easily visible as it rises because it does so very slowly; it becomes visible only when descending at measurable velocities.

3. Shock Waves and Cloud Collisions

Kahn (1955) proposed that clouds are heated by mutual supersonic collisions. Various authors have treated the cooling of clouds subsequent to such heating, with the sim-

plifying assumption that the gas density is constant. The results depend on certain parameters appropriate for the precise cooling mechanisms used, notably on the density of H_2 ; values of the cooling time vary from 10^5 to 10^7 yr. If the cooling time is of the same order as the estimated time between successive cloud collisions (10^7 yr), such collisions would be helpful in maintaining the high temperatures (> 100 K) observed for some clouds.

Zimmermann (1967, 1968a, b) described the dynamics of cloud collisions in more detail, neglecting the effects of magnetic field and cosmic rays. First, two clouds come together at velocities ≈ 10 km sec $^{-1}$. Detailed study was made of the central collision of two spherical clouds with a relative velocity of 18 km sec $^{-1}$. The initial parameters were $n_H = 10$ cm $^{-3}$ and $T = 100$ K. Second, upon first contact, a shock wave forms and proceeds into each cloud. The conditions immediately behind the shock are adiabatic, and the temperature rises to 3300 K. Third, after a certain time has elapsed, the gas sufficiently far behind the shock has cooled down to 20 K. The density in the cool regions rises in order to keep the pressure constant behind the shock. Fourth, the shock reaches the outer boundary of the cloud. At this point, all the gas in the two clouds has been compressed into a thin cloud of high density. The compression occurs along the line of collision, motion in the transverse dimension being neglected. The gas at the edge of the cloud begins to expand, and an expansion wave proceeds into the cloud at about the speed of sound (1 km sec $^{-1}$). This expan-

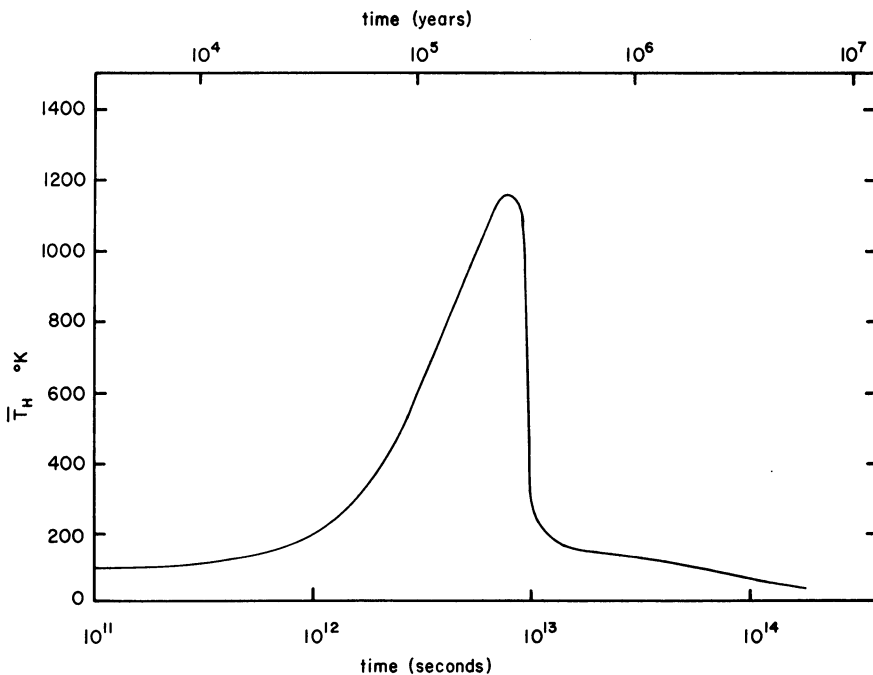


Fig. 6. The arithmetic mean temperature as a function of time of an $H\text{I}$ cloud after it starts colliding with another cloud (Zimmermann, 1968b).

sion proceeds until, fifth, the pressure in the cloud has dropped to the pressure in the intercloud medium. This requires an expansion of about a factor 15 if the expansion is adiabatic (a doubtful assumption); correspondingly, the gas temperature falls to 3 K. Cosmic-ray heating, which would increase the temperature obtained during the expansion phase, was neglected.

Figure 6 shows the mean temperature of the whole cloud (solid curve) as a function of time. One sees first the increase in temperature as more of the cloud has gone through the shock, and then the rapid drop as the shock reaches the opposite side and cooling is completed. Since a large fraction of the cloud is at a very low temperature in the postshock region, the results would presumably be quite different if the harmonic mean temperature were shown in Figure 6, instead of the arithmetic mean temperature. Such a curve would be useful for comparing with 21-cm measurements.

Steady shocks in H I clouds were studied by Wentzel (1967) and by Field *et al.* (1968), including the effect of a magnetic field parallel to the shock front. Neither paper considered the expansion phase described by Zimmermann. Field *et al.* included a variety of cooling mechanisms; in particular, fine-structure excitation of O I by hydrogen atoms, and excitation of the fine-structure levels of C⁺, Si⁺, and Fe⁺ by collisions with hydrogen atoms, in addition to those considered by Zimmermann. Figure 7 summarizes the results of a number of calculations of cooling time (defined as the time for the gas to drop back to 100 K from the postshock temperatures) for various shock velocities and magnetic field strengths. This figure is based on $H_2/H = 10^{-3}$ but the results for $H_2/H = 10^{-5}$ are very similar; the latter value is probably more

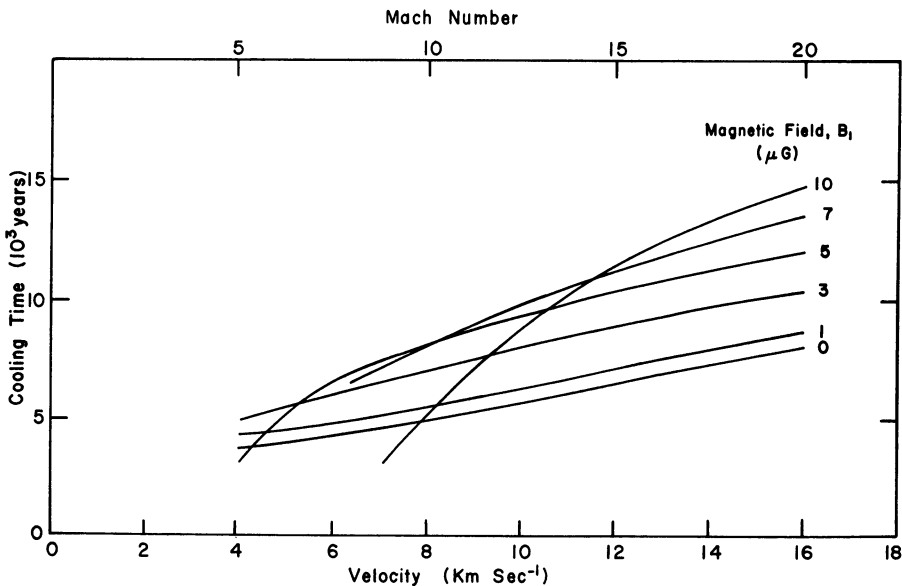


Fig. 7. Cooling time for H I gas after passage through a shock front. The magnetic field is parallel to the front. $H_2/H = 10^{-3}$, but a much lower value would give approximately the same results (Field *et al.*, 1968).

realistic (Stecher and Williams, 1967). The magnetic field has two compensating effects on the cooling time: first a decrease in time by reduction of the original post-shock temperature, and second, an increase in time by prevention of a density rise in the cool regions. As a result, the cooling times are relatively insensitive to the value of B over the plausible range 0 to $10 \mu\text{G}$. It appears that for shock velocities to be encountered in cloud collisions (half the relative velocity, about 10 km sec^{-1}) the cooling time should be less than 10^4 yr – a significant reduction from the 10^6 yr found by Zimmermann for the case $\text{H}_2/\text{H}=0$. The difference is attributed to his neglect of O I cooling, which dominates in every case.

If we could be confident of the present results, we could state definitely that cloud collisions cannot maintain cloud temperatures at the observed values, since 10^4 yr is such a small fraction of the time it takes (10^6 yr) for a 10 km sec^{-1} shock to cross a 10 pc cloud. After the first 10^4 yr has elapsed, a cloud collision actually decreases the cloud temperature by compressing a large fraction of the cloud to such a high density that cosmic rays cannot keep it hot.

Unfortunately, serious revisions are required before we can be certain of this:

(i) The cross section for excitation of O by H should be reduced a factor 3 (Bahcall and Wolf, 1968).

(ii) The free electrons from cosmic-ray ionization should be included in electron-ion cooling.

(iii) Account should be taken of depletion on grains.

(iv) The excitation of the a^4F term of Fe^+ should be included (Pottasch, private communication).

Of these, (i) and (iii) will increase T and (ii) and (iv) will decrease it, so the outcome is hard to predict. In addition to these revisions, study should be directed to the effect of an oblique magnetic field, the coupling of the magnetic field to the neutrals through the ions (Wentzel, private communication) and the dynamical effects of cosmic-ray pressure (Hudson, 1965, 1968).

A more general problem is whether cloud collisions actually occur as described. As Pikel'ner (1967) points out, the simple concept of a cloud collision will be modified by the presence of a magnetic field. He estimates that a typical cloud will oscillate perpendicular to the field about its mean position with an amplitude of 10 to 15 pc. Since this distance is considerably less than the mean free path for collisions, the number of collisions will be reduced. Furthermore, collisions perpendicular to the field will tend to be elastic as shown by Pikel'ner (1957) and Kogure (1965); we have mentioned evidence of this above, in the reduced heating by shock waves normal to the field. In summary, collisions across the field will be relatively ineffective in heating the gas.

Pikel'ner has argued that motions along the field are unimpeded, because the depth of the wells in which they settle is small – only a few pc – so that little displacement of the lines of force is necessary for motion to occur. This agrees with the observational fact that random cloud motions seem to occur equally parallel and perpendicular to the field. At first sight it might appear that deep wells (such as might occur in the

Parker model) would keep gas from moving along the field. It is true that energy must be put into the expanded magnetic field in front of the cloud to compress it to the value within the cloud. However, as much of this energy is regained again upon emerging from the rear of the cloud, the net loss of kinetic energy is small. We note that the speed of the cloud is of the same order ($\approx 10 \text{ km sec}^{-1}$) as the sound speed in the intercloud medium, so there is probably little dissipation by a shock wave there. We tentatively conclude that collisions along the field can occur in a one-dimensional way, although much analysis remains to be done here.

If the collision event takes place only between clouds (or parts of clouds) on the same lines of force, transverse motion is inhibited, and the description of Zimmermann is quite accurate. While the speed of cooling behind the shocks is still uncertain, we can be confident that it is somewhat faster than the collision itself, so that collisions along the lines of force are quite inelastic. Each cloud therefore grows in mass, and the trend toward larger clouds will continue until some other process intervenes. According to Spitzer (1968a, p. 221), gravitational collapse for masses larger than the Ebert-Bonnor critical mass is one such process. For a pressure in the intercloud medium $p = 2.5 \times 10^{-13} \text{ dyne cm}^{-2}$ ($nT = 1800$), this critical mass is $\approx 10^3 M_{\odot}$; Spitzer argues, on the basis of the virial theorem (and more detailed arguments by Field, 1969, confirm this), that a field of $3 \mu\text{G}$ will increase this critical mass only by a factor of 2 or 3. Presumably fragmentation and star formation is the normal outcome of this process. According to Oort (1954), after formation, the more massive stars will create HII regions in the remaining gas of the cloud and accelerate it outward, accounting for the kinetic energy retained by the clouds in spite of their inelastic collisions. We shall return to the problem of HII regions below, but here we note that in the Parker model, all this would occur in a massive cloud at the bottom of a magnetic well, and kinetic energy would preferentially be transmitted by clouds thrown out along the field.

Field and Saslaw (1965), Field and Hutchins (1968), and Penston *et al.* (1969) proposed and developed a statistical model of the cloud mass spectrum, assuming that it is fed by small units from expanding HII regions and is drained by collapse of massive clouds. They found a quasi-equilibrium spectrum $M^{-3/2}$, which is in fair agreement with observation. The latter paper shows that under certain reasonable assumptions the velocity distribution of clouds should obey equipartition and be Maxwellian. All of this work should be revised to take account of additional phenomena, such as the dynamical interactions with intercloud gas, the generation of new clouds by instabilities, and the constraint of one-dimensional motion imposed by the magnetic field.

All of the work on cloud collisions suffers from inadequate contact with observation. While radio astronomers have given us much data on mean values of temperatures, density, and magnetic field, I am not aware of observations of an individual cloud collision or even an individual shock wave (in the 21-cm line). It is dangerous to proceed with theoretical models without such observational checks. Field *et al.* (1968) estimated the fluxes of infrared lines between 4 and 156μ , which would be emitted by

the post-shock gas, and concluded that these are marginally observable with advanced equipment. Pottasch (1968) and Osterbrock (1969) have performed similar calculations for H I gas under other conditions. If such lines can be observed, information about temperature, density, composition, and spatial configuration can be achieved. Since the emission should be particularly intense from shock waves and other regions where dissipation is occurring, infrared measurements should tell us about the small-scale dynamics of the medium.

4. H II Regions; Expansion and Inhomogeneities

It has long been realized that the high temperature of H II regions ($\lesssim 10\,000\text{ K}$) and a similar density with respect to neighboring H I regions provide a pressure gradient which may be dynamically important. In the years since the last Symposium, considerable calculation has clarified this situation.

Spitzer and Savedoff (1950) calculated equilibrium temperatures of H II regions, balancing the gains by photoionization of hydrogen against collisional losses due to excitation of both fine-structure and optical levels of impurity atoms. They obtained values in the range $(5\text{ to }13) \times 10^3\text{ K}$ and pointed out that the observed strength in many nebulae of $\text{O}^+ \lambda 3727$ ($E/k = 39\,000\text{ K}$) requires temperatures of this order. Burbidge *et al.* (1963) revised these calculations, emphasizing the importance of fine-structure transitions such as $^2P_{3/2} - ^2P_{1/2}$ of Ne^+ at $12.8\ \mu$. Taking $\text{Ne}/\text{H} = 5 \times 10^{-4}$ and $\Omega = 1$ (est.), they found $T = 3500\text{ K}$ if the central star has a T of $30\,000\text{ K}$. Osterbrock (1965) showed that Ω was only 0.3, so that T increased to 6000 K for the same type of star. He also called attention to the fact that collisional de-excitation of O^{++} , N^+ , and C^+ for $n_e > 10^4\text{ cm}^{-3}$ reduces the cooling rate. Hjellming (1966, 1968) did a complete calculation, taking into account the detailed ionization equilibrium of the elements responsible for cooling, the transfer of Lyman continuum radiation, the hardening of the radiation by absorption as it proceeds through the nebula, the effects of collisional de-excitation, and the detailed flux distribution of the stellar atmosphere. He used modern Ω 's (Saraph *et al.*, 1969). This work was extended to non-uniform densities by Rubin (1968). For a main-sequence star of $30\,000\text{ K}$ Hjellming obtained 5800 K for a low-density region near the star. This temperature drops by about 800 K in the outer parts of the nebulae because of the Lyman-continuum diffuse radiation field. If one raises n_e to $\sim 10^4\text{ cm}^{-3}$ then T will go up by about 2000 K , whereas if one increases the optical depth τ at the Lyman limit to 2, T goes up by about 1000 K . Hjellming concludes that 6000 K is a typical temperature, with values up to 9000 K being possible in dense or optically thick regions, and values down to 4000 K in low-density outer regions. Inspection of his results indicates that fine-structure transitions often play only a minor role. This would be especially true if the abundance of neon is only 6×10^{-5} (Flower, 1969; Pagel, 1969) rather than 4×10^{-4} as he assumed. We note that in addition to the temperature equilibrium calculations described here, Hjellming computed the structure of various types of ionization fronts.

There is some controversy as to whether theory and observations yield different

temperatures. Peimbert (1967) gives optical results for three nebulae, based on the ratios of auroral-to-nebular lines of heavy ions and line-to-continuum intensities of hydrogen. He obtains values between 6000 and 14 000 K in the center of Orion, depending on the method. Kaler (1968) reviews other optical data with similar results. Measurements based on the ratio of radio recombination lines to free-free continuum seem to give systematically lower values (Dieter, 1967; Mezger, 1968) but are unreliable because of non-LTE effects. Radio measurements based on fitting the shape of the free-free continuum give very low values in two cases (Terzian *et al.*, 1968). These authors obtained 3000 (+2000; -1000) K for Orion. Since the method is sensitive primarily to the temperature of the outer parts, they suggest that the discrepancy between optical and radio measurements reflects a strong inward temperature gradient. This may be in qualitative agreement with the calculations of Hjellming (1968), because the diffuse radiation field tends to decrease T in the outer parts, while collisional de-excitation tends to increase it in the dense inner parts.

One of the aspects of H II regions of interest here is the dynamical expansion caused by high thermal pressure. The above discussion indicates that the theoretical situation

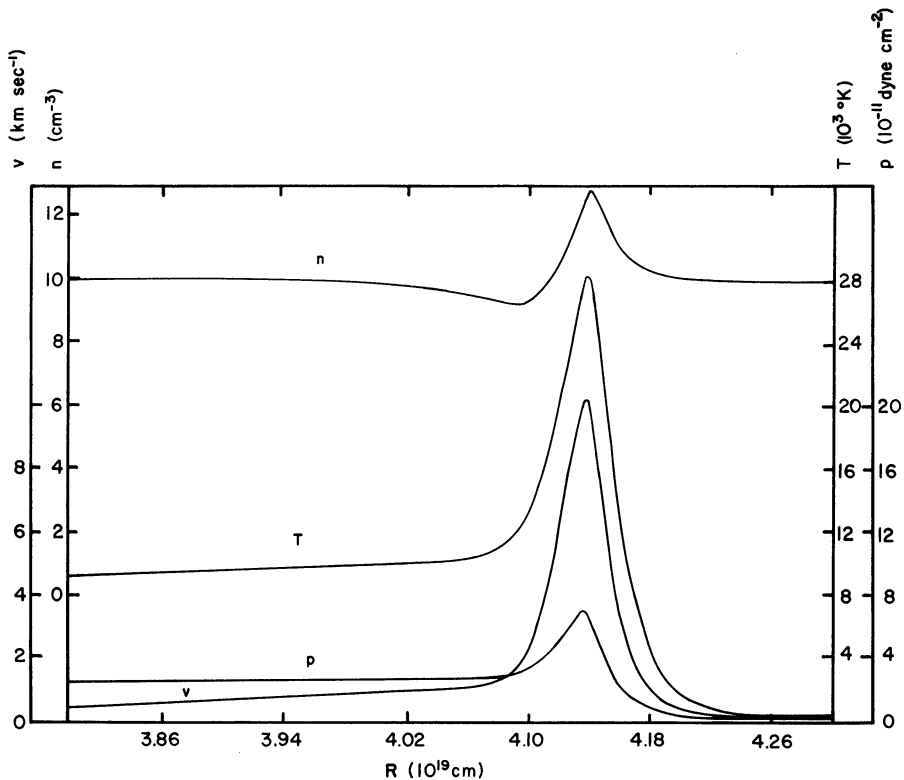


Fig. 8. Variation of velocity v , number density n , temperature T , and pressure p as a function of radius R from the star in the vicinity of an ionization front moving away from an O star with $T = 42\,000\text{ K}$ after $3 \times 10^4\text{ yr}$ (Mathews, 1965).

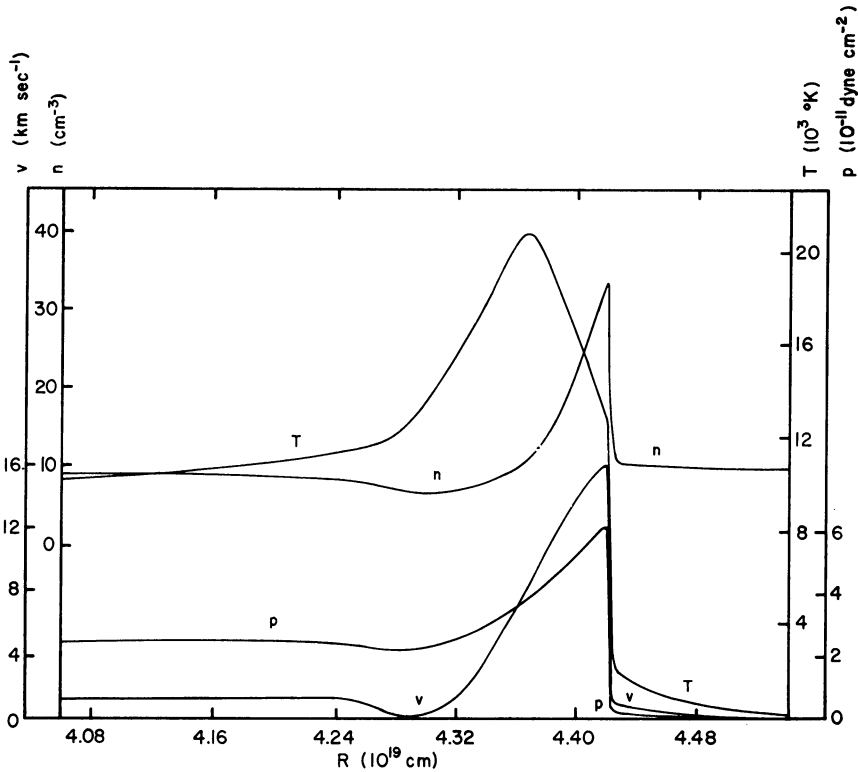


Fig. 9. The same as in Figure 8, but after 6×10^4 yr.

regarding the temperature is well in hand and that there are no glaring disagreements with observation. We now turn to two recent discussions of the expansion problem, by Mathews (1965) and Lasker (1966a, b). (A review of these and early studies of the problem has been given by Mathews and O'Dell, 1969). Earlier work on this problem had been restricted to similarity solutions but the papers of Mathews and of Lasker are based on numerical solution of the partial differential equations involved for specified initial values and use finite-difference methods.

Mathews (1965) treats the evolution of an H I region with $n_{\text{H}} = 10 \text{ cm}^{-3}$ and $T = 100 \text{ K}$ as it is ionized by a $30 M_{\odot}$ star, whose main-sequence temperature is 42 000 K. Mathews treats in detail the luminosity of the star during its approach to the main sequence, since the time of approach (6000 yr) is comparable to $1.3/\alpha n = 10^4$ yr, the classical time of formation of the H II region. As predicted from earlier work, an R-type ionization front moves very rapidly into the gas, averaging 1000 km sec^{-1} in the first 10^4 yr. Velocities of the ionized gas remain less than 0.2 km sec^{-1} during this phase, in accordance with prediction for steady R-type fronts. At 6×10^4 yr, the buildup of the pressure wave behind the ionization front has developed into a shock front which leads the ionization front by 3 per cent of the radius. Although the com-

putations were terminated at this point, the much-discussed detachment of a shock front which permits the ionization front to be D-type by virtue of its compression of the neutral hydrogen was demonstrated by Mathews (Figures 8 and 9).

Lasker (1966a) assumed that his star starts to shine instantaneously; his initial conditions are similar to those of Mathews. Unlike Mathews, who followed the Lyman-continuum radiation field in detail, he approximates the ionization front by a discontinuity through which atoms flow and are instantly ionized and heated to 10 000 K by the ultraviolet photons which escape absorption in the HII region; hence his results are unreliable in the immediate vicinity of the ionization front. On the other hand, in certain models, approximate allowance is made for cooling behind the shock (in the HI). His computations extend up to 2×10^6 yr, much longer than those of Mathews, and one can therefore examine the later stages of development which are probably more pertinent to observed HII regions. Figure 10 shows his results for a case with no HI cooling. One sees the detachment of the shock at about 8×10^4 yr, and the progressive acceleration of ever-increasing amounts of HI as time goes on. At about 9×10^5 yr an expansion wave caused by the acceleration of the outer HII has travelled to the center; thereafter the central density of HII decreases as the region expands.

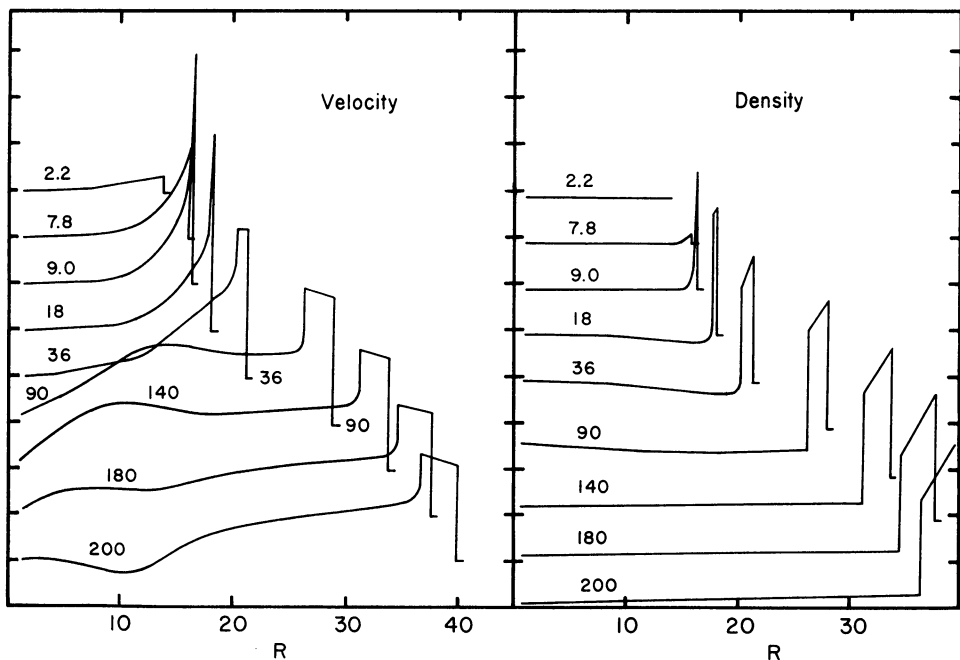


Fig. 10. Variation of velocity V and density n as a function of radius R from the star for one of the cases of expanding HII regions considered by Lasker (1966a). To each curve is added the time after beginning of the expansion expressed in units of 10^4 yr. The velocity unit is 3.2 km sec^{-1} , the density unit 6.4 cm^{-3} .

The thicknesses Lasker obtained for the H I shock seem unrealistically small. The cooling time of H I (10^4 yr) is small compared to the evolution time (2×10^6 yr). Lasker therefore considered an additional model with an exponential cooling $e^{-t/\tau}$, with $\tau = 5 \times 10^4$ yr, and found that the evolution of the H II region is changed hardly at all. On the other hand, Δr , the thickness of the compressed H I region, decreases from about 10 per cent of the radius of the shock, to only 1 per cent, with a corresponding increase in density. Apparently the I-front acts like a piston which forces a shock at about 10 km sec^{-1} to compress the H I up to about the same pressure as the H II. It does not matter much whether this pressure is due to a high temperature (no cooling) or a high density (cooling). In the absence of magnetic fields, however, one would expect cooling to be effective in bringing about compression, and consequently, that a very thin but dense H I shell should be found around old H II regions. The available

TABLE IV
H II regions with observed H I shells

Nebula	Radius of I-front (pc)	Outer radius (pc)	$\Delta R/R$	Expansion velocity (km sec^{-1})	References
Rosette	20	50	0.6	15	Raimond, 1966
NGC 6910	1.5	5.5	0.7	10	Davies and Tovmassian, 1963
λ Ori	13	30	0.6	8	Wade, 1958
I Mon	40	90	0.7	23	Girnstein and Rohlfs, 1964
NGC 5146	1.5	3.3	0.5	< 6	Riegel, 1967
NGC 281	7.5	15	0.5	< 8	Riegel, 1967
DR 21	0.4	—	—	25	Thompson, Colvin and Hughes, 1969

observations seem at first glance to disagree with this. In Table IV we note some 21-cm observations of expanding H I shells around H II regions. In every case $\Delta R/R$ is much larger, not only than the value of 10^{-2} predicted for H I cooling, but also than the value of 10^{-1} predicted for purely adiabatic shocks. It is interesting that both λ Ori and NGC 5146 have dust shells around them of the order of a few per cent of the radius (Wade, 1958; Riegel, 1967). Possibly this dust (which in places has several magnitudes of extinction) marks the actual compressed region. From this conclusion it would follow that the 21-cm antenna beams are too large to resolve the expanding shells (Lasker, 1966a; Riegel, 1967). We note that infrared observations at the wavelengths of cooling transitions should be able to locate the H I shock unambiguously with high resolution (Field *et al.*, 1968).

Among the mechanisms suggested for increasing the width of the compressed region are turbulent or magnetic pressures in the H I region, a decrease in driving pressure as the star evolves, or expansion of the outer part of the cloud into the inter-

cloud medium. Lasker (1966b) investigated the effects of magnetic fields on I-fronts and concluded that if $Q = B^2/8\pi p \lesssim 1$ ahead of the front, there is little effect (for either R- or D-type fronts). In fact, $Q = 1.5$ if $B = 3 \mu\text{G}$ in a cloud at the critical pressure found by Field *et al.* (1969), $2.5 \times 10^{-13} \text{ dyne cm}^{-2}$. After adiabatic compression by a 10 km sec^{-1} shock, Q falls to 0.33, so that an I-front moving into such a region is relatively unaffected. However, cooling behind the shock increases Q considerably (by reducing p). Lasker (1966a) discusses this situation for the two extreme cases: (i) the angle between magnetic field and radius from the star is 0° and (ii) the angle is 90° . Along the magnetic field (zero angle) there is no effect so the compressed H I region remains thin, while at an angle of 90° simple equations shows that the compressed region is about 10 times thicker. As Lasker notes, however, at intermediate angles the thicknesses may not be intermediate, since there can be discontinuous transitions from one branch of solutions to another (corresponding to the fast and slow hydromagnetic waves) as the angle between \mathbf{B} and the radius vector varies from 0° to 90° . More work on this problem is recommended, both observationally (by seeking thin compressed regions) and theoretically (by verifying the correct branches of solutions for the shock and I-fronts in the context of the actual initial-value problem).

Lasker (1967) summarized numerical results on H II-region evolution in a simple way. The results indicate that (i) the radius r_i of the ionized region very quickly (about $10^4/n \text{ yr}$) approaches the value $s_0 = (3S/4 \alpha n_0^2)^{1/3}$ predicted by static Strömgen theory (α is the recombination coefficient, S the rate of emission of ionizing photons from the exciting star), (ii) thereafter r_i grows slowly as the I-front moves into the shocked H I, always satisfying the Strömgen relation for the actual density, n , in the H II region, (iii) n remains rather uniform and the velocities within the H II region remain $< \frac{1}{4}c$ (the isothermal sound speed in the H II region), (iv) the pressure is constant through the I-front after the early formation period, (v) the shock front never moves far ahead of the I-front. The last fact means that $dr_i/dt \approx v_s$ where the shock velocity v_s for strong shocks is given by $(4p'/3p_0)^{1/2} c_0$, where p_0 and c_0 refer to the unshocked H I and p' to the shocked H I. Since $p' = p_{\text{II}} = \rho_{\text{II}} c^2$ from (iv), where c is the sound speed in the H II region, and $p_0 = \rho_0 c_0^2$, it follows that $v_s = c_1 (n/n_0)^{1/2}$, where $c_1 = 2c/\sqrt{3}$. From (ii), $n/n_0 = (r_i/s_0)^{-3/2}$, so that

$$dr_i/dt \approx v_s = c_1 (r_i/s_0)^{-3/4}, \quad (2)$$

which has the solution

$$r_i = s_0 (1 + 7c_1 t/4s_0)^{4/7}, \quad (3)$$

where t is the time elapsed since the Strömgen region formed ($r_i = s_0$). Hence

$$n = n_0 (1 + 7c_1 t/4s_0)^{-6/7}. \quad (4)$$

These expressions are found to represent the numerical results quite well. The I-front can be considered to be a piston exerting a pressure ρc^2 . As it moves, it does work on the H I ahead. Of the energy imparted by the strong shock, half goes into heat and

half into kinetic energy. Hence the kinetic energy of the accelerated H I is

$$E_k = \int p \, dV = 2\pi\varrho_0 c^2 s_0^3 \int_{s_0}^{r_i} \left(\frac{n}{n_0}\right) \left(\frac{r_i}{s_0}\right)^2 \frac{dr_i}{s_0}, \quad (5)$$

which from Equations (3) and (4) is equal to

$$E_K = (4\pi/3) \varrho_0 c^2 s_0^3 [(r_i/s_0)^{3/2} - 1]. \quad (6)$$

This is readily shown to be $\frac{2}{3}[1 - (r_i/s_0)^{-3/2}]$ times the thermal energy in the H II region. It is important to compare E_K to the energy available from ultraviolet photons, $E_L = St \Delta E$, where ΔE is the average energy going into heat per ionization ($\Delta E \approx kT_*$). By using the Strömgen relation one can show that

$$\varepsilon = \text{efficiency} = \frac{E_K}{E_L} = \frac{2kT}{\Delta E} \frac{1}{\alpha nt} \left(1 - \frac{n}{n_0}\right). \quad (7)$$

For late phases of evolution ($n \ll n_0$) this has the simple interpretation that the total number of recombinations per proton is αnt , and since each is balanced by an ionization involving the expenditure of ΔE , the total energy put into the H II is $\alpha nt \Delta E$ erg per proton. Almost all of this energy from the star goes into keeping the gas hot, only the fraction $2kT$ being the pressure (which can do work) available per proton. Lasker shows that when t = main sequence lifetime of the star and n = the corresponding final density starting with $n_0 = 10 \text{ cm}^{-3}$, ε varies from 0.002 for spectral type O5 to 0.07 for B1. When weighted by the numbers and fluxes of stars of various spectral types, the average efficiency is $\varepsilon = 0.0024$ (as O-stars supply most of the energy). A higher efficiency ($\varepsilon' = 0.014$) is required (Spitzer, 1968a) to account for the motions of interstellar clouds by the kinetic energy of H I shells accelerated by H II regions (Oort-Spitzer mechanism). As the discrepancy is only a factor 6, our tentative conclusion is that a significant fraction of the energy dissipated in cloud collisions, but probably not all, originates in H II regions. The alternatives include inflation of the magnetic field with resulting downward motion of gas clouds (Hoyle and Ireland, 1960; Parker, 1966), infall of intergalactic matter (Oort, 1969), electromagnetic radiation from neutron stars (Hoyle *et al.*, 1964), and expansion of supernova shells (Shklovskii, 1962; Kahn and Dyson, 1965; Kahn and Woltjer, 1967; Spitzer, 1968b).

Kahn and Woltjer (1967) make quantitative estimates of the efficiency of cloud acceleration by supernova shells. They adopt a snowplow model appropriate for rapid cooling of shocked gas, and consider two cases of shell speed. If the speed is relativistic, the efficiency is $\varepsilon_R = \frac{1}{2}(v/c)^{1/2}$, while for non-relativistic shells it is $\varepsilon_{NR} = v/V$, where $v = 10 \text{ km sec}^{-1}$ is the speed of interstellar clouds and $V = 7000 \text{ km sec}^{-1}$ is an estimated shell velocity. Hence $\varepsilon_R = 0.003$ and $\varepsilon_{NR} = 0.0015$, of the same order as found by Lasker (1967) for H II regions. Using the data for dissipation in cloud collisions used above, adopting the non-relativistic model, and taking the supernova rate as 1 per 100 yr, we find that the required kinetic energy of the shell is 3×10^{51}

erg, compared to the observational estimate of 10^{51} erg (Minkowski, 1968). HII regions contribute $\frac{1}{2}$ and supernovae contribute $\frac{1}{3}$ of what is required, so together they may well account for the motions of clouds.

When the acceleration is complete, the radius of the supernova shell is about 50 pc according to Kahn and Woltjer (1967). Shane and Katgert (1967) have tentatively identified such a feature observationally. The large radius means that many clouds are actually members of expanding shells. From the model of Kahn and Woltjer one can estimate that the fraction is roughly 20 per cent. It is interesting that according to Spitzer (1968b), about 10 per cent of the calcium clouds observed are of relatively high velocity and show an excess of negative velocities as if they had been recently accelerated away from the vicinity of the early-type stars against which they are seen.

In closing our discussion of the expansion of HII regions, reference is made to the effects of stellar winds. Mathews (1966) has shown that mass loss of the order $10^{-5} M_{\odot} \text{ yr}^{-1}$ can, by the pressure it exerts on the interior of the gas, cause the observed hole at the center of the Rosette Nebula, while Pikel'ner (1968b) has constructed a model based on stellar winds to explain the high internal velocities observed in some HII regions (Shcheglov, 1968). Mathews (1967) has also considered the effect of radiation pressure on grains within the HII region. He finds that the electrostatic drag on the charged grains can transmit the pressure effectively to the gas, with the result that the inner part of the Rosette Nebula can be evacuated.

We now turn to our final topic, inhomogeneities in HII regions. There is ample evidence that such inhomogeneities exist: determination of $\langle n_e \rangle$ by OII $\lambda 3727$ yields lower values than determination of $\langle n_e \rangle^{1/2}$ from hydrogen emission; the forbidden lines yield higher temperatures than the hydrogen lines and the radio continuum; the emission lines are broadened by motions comparable to the thermal velocity and, in some cases, are split in local regions. According to Kahn (1967), one can possibly interpret the discrepancy between the ages of the stars in Orion and the calculated expansion age of the Nebula, if the Nebula is actually made up of many dense globules of neutral matter which are being slowly ionized at present. This model can also account for the observed line splitting, by streams of ionized gas emerging from the globules and colliding to form shock waves. No doubt it would also account qualitatively for the variations in n_e and T observed, and possibly for the presence of OI $\lambda 6300$, which has always been hard to understand in a completely ionized region.

Some divergence exists on the theoretical treatment of globules. T Tau stars are often seen imbedded in elephant-trunk or comet-tail structures at the edges of HII regions and Dibai (1963) suggested that this could be explained as the result of a converging shock wave, initiated by the high pressure in the surrounding HII. Later Dibai and Kaplan (1964) studied such a shock wave by the method of similarity solutions and obtained velocities up to two times the speed of sound in HII regions and density increases (after passage of first a converging and then a reflected shock) up to 10^5 .

Dyson (1968), on the other hand, analyses the situation by assuming that the globule is in hydrostatic equilibrium under self-gravitation, internal pressure, and a boundary

pressure supplied by the H II streaming away from its surface in a D-critical I-front. No shocks are included, and a relation is derived between the mass and radius of the globule for a given external ionizing radiation field. This corresponds to the fact that the I-front (assumed to be D-critical) has a definite value of $p + \rho v^2$ which must be balanced by the internal gas pressure. The theory of truncated isothermal spheres then leads to the mass-radius relation. As the globule evaporates, gravitational binding usually becomes less and less important; and confinement by external pressure, more important. Gravitational instability does not occur, and the whole globule is evaporated (on a time scale of 2 to 8×10^4 yr per solar mass). On the other hand, globules rather close to instability initially can be set into collapse as time goes on.

The question of whether the front is D-critical was analyzed by Kahn (1969). He finds rather that the front is strong D-type, progressing subsonically into the globule and ejecting ionized gas supersonically into the H II region. He mentions the possibility that the detailed heating and cooling processes within the I-front may require a pressure wave to propagate ahead of the I-front, but otherwise omits deviations from hydrostatic equilibrium within the globule. Mendis (1968, 1969) analyzed the situation further, taking into account the fact that the globule is probably composed of H₂ rather than H at the high densities contemplated (10^6 cm^{-3}). He concludes that the front is strong D-type for all stars earlier than O9.

Considerable effort has been expended on discovering whether elephant trunks (which may later become globules) can arise as a result of instability. The present consensus appears to be that instabilities (Rayleigh-Taylor; ionization front) are present but either yield parameters in disagreement with observation or are too weak (Mathews and O'Dell, 1969). Rather, it seems more natural to suppose that large density fluctuations are already present in the H I region before it is overrun by the I-front. Two possibilities suggest themselves. The first is that dense concentrations (10^4 cm^{-3}) are present in the cloud before the hot stars begin to shine. As their internal pressure must greatly exceed that of the normal intercloud medium, they must be gravitationally bound. In this case, it is hard to see why they do not collapse in the free-fall time (10^6 yr), since isothermal spheres of this kind are unstable. Possibly they are actually collapsing, and represent unstable fragments of the original cloud which would have formed stars if the ultraviolet radiation field had not intervened.

Alternatively, one might identify them with the small cloudlets (1 to $30 M_{\odot}$) discovered by Heiles (1967a) in the interstellar medium. While the masses roughly agree, the densities do not. To account for this, we note that by the time the I-front makes the dense clouds visible, they have already been overrun by a shock front which will certainly compress them. The degree of compression can be estimated as $2T_2/T_1 \approx 20000/T_1$, where T_2 is the H II temperature and T_1 is the H I temperature. Ordinarily T_1 is taken to be 100 K, so the compression is a factor of 200. In regions which are already denser than average, however, cooling processes are particularly effective, and T_1 might drop as low as 10 K as a result of cooling by C^o. The resulting compression of a factor of 2000, would result in a final density of $2 \times 10^4 \text{ cm}^{-3}$.

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