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SUPERNOVAE VERSUS MODELS OF THE INTERSTELLAR MEDIUM AND THE GASEOUS HALO

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ABSTRACT

We consider the effects of the actual distribution of Type I and Type II supernovae (SNs) on the interstellar medium (ISM) and gaseous halo in spiral galaxies. Type I SNs are distributed as the Population I and old Population II starlight, in a thick disk decreasing exponentially with galactic radius. Some Type I SNs are located within the H I disk and disturb the interstellar medium in a random fashion. The others are located outside the disk and heat the gaseous halo. Type I SNs only dominate the ISM toward the galactic interior and may be responsible for the H I "holes" observed in the interior of many galaxies.

Type II SNs arise from young, massive progenitors that are born in stellar clusters and associations. They concentrate where star formation peaks, i.e., in the inner galaxy with the molecules. The stars in a cluster all act together to blow a cylindrical supercavity in the gaseous disk. The diffuse gas within this cylinder escapes into the gaseous halo. We calculate mass and energy injection rates into the halo and consider the possibility of a galactic wind from the halo, including the effect of the massive halo now known to be associated with spiral galaxies.

We predict that Type II SNs dominate the ISM by a large factor. This is in violent disagreement with galactic and extragalactic observational data. We also predict either that the mass injection rate into the halo is much larger than allowed by diffuse X-ray observations or that the Galaxy has a wind. The fundamental problem is that SNs eject more energy into the ISM or the halo, or both, than is allowed by observational data. Reconciliation requires major, fundamental changes in either the theory or the observational input data, and we mention a few possibilities.

Subject headings: galaxies: Milky Way — galaxies: structure — interstellar: matter — nebulae: supernovae remnants — stars: supernovae

I. INTRODUCTION

Some years ago, Field, Goldsmith, and Habing (1969, hereafter FGH) developed a theory of a quiescent interstellar medium (ISM) in hydrostatic equilibrium under the z-component of the Galactic gravitational field. Somewhat later, Cox and Smith (1974) pointed out that supernova (SN) explosions occur frequently enough to upset the quiescent equilibrium envisioned by FGH. McKee and Ostriker (1977, hereafter MO) then presented a detailed theory of a SN-dominated ISM, under the assumption that SNs occur randomly in time and space.

The MO model predicts that most of the interstellar volume is filled with a hot ionized medium (HIM) and that the cold H I gas exists predominantly in clouds. This picture is in accord with some observational data, specifically the diffuse X-ray emission and the O VI absorption line data. However, it disagrees with many aspects of H I data, especially the smooth distribution of H I and the relatively large abundance of warm H I. The author has never been able to fully accept the model because of these inconsistencies. A comprehensive application of the theory, including the best observational parameters known to the author, is the purpose of the present paper, which is a substantial modification of a previous paper (Heiles 1985, hereafter Paper I).

The MO prediction is strong because the SN rate in the Galaxy is so large. Below, we adopt a total Galactic SN rate of 0.05 yr^{-1} (interval 20 yr). If all the SNs explode within the H I disk of radius 10 kpc and full thickness 370 pc, MO's SN rate parameter $S = 4.3 \times 10^{-13}$ SN pc^{-3} yr^{-1}. For the interstellar parameters adopted below, the porosity parameter $Q = 4.4$. Formally, this means that the fraction of interstellar volume occupied by the H I gas is only $1/(1 + Q)$ (Shu 1982), or ~ 0.18.

In this paper we attempt to reconcile the theory and observations by using what is known about the distribution of SNs in space and time. Type II SNs are correlated in space and time because they arise from young stars, which are formed in clusters; Type I SNs occur randomly in time but are not distributed uniformly within the Galaxy. The discussion is oriented toward our own Galaxy but is sufficiently general so that it can be applied to spiral galaxies other than our own.

Section II discussed values of Q for correlated SNs; § III discusses observational values for model parameters. Section IV discusses Type I SNs and their effect on the ISM; the predictions and observations are compared in § IV*c*, and they appear to be consistent.

Type II SNs are discussed in § V. The predictions and observations are compared in § V*d*; we find very large discrepancies. Section VI discusses the effect of both types of SNs on the gaseous halo; again, the predictions and observations show large discrepancies. Section VII summarizes the results and suggests possible modifications in the observational input data and the theory to reconcile the discrepancies.

II. THE STRUCTURE OF THE ISM UNDER CORRELATED SNS

a) A Three-dimensional Medium

SNs create large, rarefied cavities of hot gas in an ISM. The quantity Q, the porosity parameter of Cox and Smith (1974) describes the fraction of interstellar volume occupied by such cavities, which is equal to $Q/(1+Q)$.[1] The quantity Q can be reliably calculated only if Q is itself small, because the theoretical models of SN cavities assume an ambient ISM whose temperature is much smaller than that inside the cavity. All of our expressions for Q are derived under this assumption; thus, their application requires $Q \lesssim 1$. For uncorrelated SNs and a three-dimensional medium, MO's equation (2) gives the convenient expression

$$Q_{3D} = 10^{-0.59} n_0{}^{-0.14} P_{04}{}^{-1.30} E_{51}{}^{0.28} (E_{51} S_{-13}) , \quad (1)$$

where E_{51} is the SN energy released to the ISM in units of 10^{51} ergs, S_{-13} is the SN rate in units of 10^{-13} pc^{-1} yr^{-1}, n_0 the ambient density of H atoms in units of cm^{-3}, and P_{04} the ambient pressure in units of $10^4 k$, where k is Boltzmann's constant. Here we have multiplied the MO value of Q by 0.5 to allow approximately for the fact that a SN shell expands with $R \propto t^\eta$, with η ranging from $\frac{3}{5}$ to $\frac{2}{5}$ as the shell expands, and after reaching its full extent is encroached upon by the ISM at velocity $v_{\rm rms}$ (see eq. [1] of MO). The velocity dispersion of the ambient intercloud gas, $v_{\rm rms}$, we take equal to the isothermal sound speed $(p/\rho)^{1/2}$.

We suppose that Type II SNs are correlated in space and time, with N stars per cluster. If the N stars were to explode simultaneously as a single large explosion, E and S would become NE and S/N, respectively. The product ES would remain unchanged and Q would differ from the purely random case by a factor of $N^{0.28}$. In fact, however, the N stars release energy continually over $\sim$30 Myr in the form of stellar winds and SNs, with a power output that is approximately independent of time (McCray and Kafatos 1987). The effect is equivalent to a very powerful stellar wind—a "superwind." In this case an equation similar to equation (1) applies, but it yields a higher value of Q because a wind is more efficient at transferring energy to the ambient medium than a single large explosion. Thus, both because winds are more efficient and a single large explosion is more effective, if a three-dimensional ISM is dominated by uncorrelated SNs (corresponding to $Q \gtrsim 1$), it will be dominated even more by correlated SNs.

b) *A Two-dimensional Medium*

i) *The Porosity Parameter Q*

In spiral galaxies, the interstellar matter is distributed in a thin disk. If the size of regions influenced by explosions or winds is smaller than the thickness of the disk, then the three-dimensional result above applies. However, as N increases this condition must eventually be violated, because the supercavities will break through the thin disk. This is the two-dimensional case.

The details of the "breakthrough" process have been treated theoretically by Tomisaka and Ikeuchi (1986) for sequential supernova explosions. Rather than incorporating these detailed calculations, we will simplify the problem by assuming the gas density is independent of $|z|$ up to height h and by assuming that upon breakthrough the internal pressure of the supercavity is dissipated into the halo of the galaxy, where the pressure is smaller than in the disk. At this point the internal supercavity pressure quickly drops to the value in the halo, which we take as being much smaller than disk's ambient pressure. The shell then expands with constant momentum (the "snowplow" phase) until the shell velocity equals $v_{\rm rms}$, carving out a cylindrical supercavity in the disk. Subsequently, the ambient gas encroaches into the supercavity with velocity $v_{\rm rms}$. We neglect the Galactic gravitational field. This is justified because shells break through the disk at velocities large compared to $v_{\rm rms}$—and since the disk height h is determined by the hydrostatic equilibrium between the gravitational field and $v_{\rm rms}$, the gravitational field can have little effect on shell dynamics before breakthrough.

Below, in § V*b*, we find that an appropriate unit for the superwind power is $L_{38} \equiv 10^{38}$ ergs s^{-1}. Stellar wind theory (Weaver *et al.* 1977) gives for the supershell radius R_s:

$$R_s = 66 L_{38}{}^{1/5} n_0{}^{-1/5} t_6{}^{3/5} \text{ pc} , \quad (2)$$

where t_6 is the supershell age in units of 10^6 yr (Myr). Alternatively, the shell velocity v_s is given by

$$v_s = 640 L_{38}{}^{1/3} n_0{}^{-1/3} R_s{}^{-2/3} \text{ km s}^{-1} . \quad (3)$$

The supershell breaks through when $R_s = h$, where h is the half-thickness of the disk; h_{100} is in units of 100 pc. At this point, the supershell velocity at breakthrough, v_b, is

$$v_b = 30 L_{38}{}^{1/3} n_0{}^{-1/3} h_{100}{}^{-2/3} \text{ km s}^{-1} . \quad (4)$$

We have implicitly assumed that the internal supercavity pressure greatly exceeds the external pressure. If this is not the case, then the shell expansion will quickly be retarded by the ambient medium. The formal condition that the internal pressure exceeds the external pressure at breakthrough can be written

$$L_{38} > 2.7 \times 10^{-2} p_{04}{}^{3/2} n_0{}^{1/2} h_{100}{}^2 . \quad (5)$$

For Galactic parameters, this means $L_{38} > 0.01$ or, alternatively, $N \gtrsim 1$ (where N is the number of SNs per cluster). However, for small N our approximation that a large number of explosions simulates a wind is invalid; instead, the solution for a single explosion must be used. We assume that the crossover point between these occurs when the energy of all N explosions, acting simultaneously, is just sufficient to break through (i.e., MO's $R_E > h$); this gives $N = 12$ for the Galactic parameters in Table 1.

There are several other implicit assumptions. First, we assume that molecular clouds within the supercavity do not act as significant centers of condensation for the hot gas; if they do, they will greatly reduce the fraction of SN energy available

TABLE 1

ADOPTED VALUES FOR PARAMETERS

Symbol	Value	Description	Units
E_{51}	1.0	Supernova energy	10^{51} ergs
N..........	40	Number of Supernova per cluster	...
n_0	0.24[a]	Intercloud gas density	cm^{-3}
h	185[b]	Exponential scale height of intercloud H I	pc
$h_{\rm SN\,I}$	325	Exponential scale height of Type I SNs	pc
P_{04}	0.40	Intercloud gas pressure	$10^4 k$
$v_{\rm rms}$	9.9	Intercloud rms velocity	km s^{-1}
s	2.1[c]	Correlation factor for Type II SNs	...

[a] Inside $R_{\rm Gal} = 11.5$ kpc. See eq. (15).
[b] Inside $R_{\rm Gal} = 11.5$ kpc. See eq. (14).
[c] Inside spiral arms; 0.7 outside spiral arms (see § V*a*).

[1] This result is based on percolation theory, and is not valid for large values of Q (Shull 1987). For large Q, the fraction is equal to $1 - f_{\rm clouds}$, where $f_{\rm clouds}$ is the fraction of volume occupied by cold clouds and their envelopes. There exists no theory to calculate $f_{\rm clouds}$. In this paper, we use the value $Q/(1+Q)$, even if $Q \gg 1$, for illustrative purposes.

for mechanical work in driving the supershell expansion. Second, we assume that the low-density, 480 pc scale height H I gas found by Lockman (1984) does not affect the breakthrough process. Third, we assume that there is no significant radiative cooling in the hot interior. According to McCray and Kafatos (1986), this occurs at a radius $R_c \approx 680 L_{38}{}^{0.4} n_0{}^{-0.6}$, which for Galactic parameters is much larger than h. Finally, we assume that energy loss from evaporation of H I from diffuse cloud surfaces is small compared to evaporation from the dense shell. This assumption should be valid for diffuse (but not molecular) clouds. Before the first SN explosion, diffuse clouds should have been destroyed by ionizing photons from the hot stars (Elmegreen 1976; McKee, van Buren, and Lazareff 1984). This should have occurred out to the Strömgren radius of the entire assembly of hot stars in the cluster, equal to $65 n_0{}^{-2/3} N^{1/3}$ pc for the Galactic initial mass function discussed below in § V*a*. For the Galactic parameters in Table 1, this Strömgren radius exceeds h by factor of 2.9.

As we shall see below, observations indicate that the free parameters are roughly equal to unity, so that v_b is rather small, ~ 25 km s^{-1}. The reason is that breakthrough occurs in only ~ 5 Myr, when only $\sim 20\%$ of the total energy of the N stars has been released. Since the internal pressure immediately dissipates into the halo, the subsequent expansion of the shell is dominated by the "snowplow" solution, with no internal pressure because the hot gas in the interior escapes rapidly to the halo. The snowplow geometry is cylindrical; thus $v_s \propto R_s{}^{-2}$. The shell then expands to its final radius, R_f, when v_s decreases to $v_{\rm rms}$:

$$R_f = h\left(\frac{v_b}{v_{\rm rms}}\right)^{1/2}$$

or

$$R_f = 550 h_{100}{}^{2/3} L_{38}{}^{1/6} n_0{}^{-1/6} v_{\rm rms}{}^{-1/2} \text{ pc} . \tag{6}$$

The ambient matter then penetrates the supercavity in a time of order

$$\tau_p \approx R_f / v_{\rm rms} . \tag{7}$$

It is appropriate to express results for the two-dimensional case in terms of the SN rate per unit area of the disk. Let σ denote the individual SN rate in units of kpc^{-2} Myr^{-1}; below, in § V*c*, we show that $\sigma = 100$ is a typical value in the Galaxy, The quantity σ/N is the formation rate of clusters that produce superwinds. As discussed below in § V*a*, we also include a factor s that describes the degree to which the superexplosions are concentrated in the spiral arms. Using the two-dimensional analog of equation (1) of MO, and assuming that most of the time for shell expansion is spent in the snowplow phase, $Q_{\rm 2D} = s\sigma N^{-1} \pi R_f{}^2 \tau_p$, i.e.,

$$Q_{\rm 2D} \approx 520 s\sigma N^{-1} v_{\rm rms}{}^{-5/2} h_{100}{}^2 L_{38}{}^{1/2} n_0{}^{-1/2} . \tag{8}$$

This estimate of $Q_{\rm 2D}$ is an overestimate if condensation onto molecular clouds is important. We note that $Q_{\rm 2D}$ is larger than derived in Paper I, where we assumed $R_f = h$; for the parameters in Table 1, it is 4.2 times larger.

ii) *Interaction with the Halo*

At breakthrough, $\sim 30\%$ of the cold shell (the fraction of solid angle occupied by the halo) escapes from the disk. The shell may break out in fragments because it is expanding into an atmosphere whose density decreases outwards. The vertical velocity at escape is $\sim v_b$ or less. For the Galactic parameters in Table 1, $v_b = 25$ km s^{-1} and the shell, or shell fragments, rise to z-heights of only ~ 390 pc. They then fall back to the plane, reaching maximum velocity 29 km s^{-1} and traveling ~ 600 pc from their point of origin in the Galactic plane. Shells that are breaking through can be seen in the H I photographs of Heiles (1979); these pictures also show that some of the shells have broken into fragments.

The heated gas inside the cylindrical supercavity is at high pressure and escapes into the halo. This heated gas was evaporated from cool gas located either in the shell or in clouds within the supercavity. More evaporation will occur after breakthrough. This evaporated gas is a source of mass for the halo. We now estimate both the amount and temperature of the injected gas.

First, we assume that all of the evaporated gas comes from the shell. We use equation (34) of Weaver *et al.* (1977) to calculate the mass of evaporated gas before breakthrough, $M_{\rm ev,bb}$:

$$M_{\rm ev,bb} = 9.9 \times 10^3 L_{38}{}^{8/21} n_0{}^{1/3} h_{100}{}^{41/21} M_\odot . \tag{9}$$

After breakthrough, we use the mass evaporation rate given by equation (5) of Castor, McCray, and Weaver (1975), multiplied by the ratio of the surface area of the cylinder to that of the sphere. The evaporation rate is very sensitive to the interior temperature, varying as $T^{5/2}$. We calculate the interior temperature by using the specific heat at constant pressure, together with the assumption that all of the energy goes into either mechanical work or thermal energy of the evaporated gas. This yields the temperature of the evaporated gas in the snowplow phase, $T_{\rm ev,sn}$:

$$T_{\rm ev,sn} = 2.8 \times 10^6 L_{38}{}^{2/7} h_{100}{}^{-2/7} \text{ K} . \tag{10}$$

We now assume that the evaporation occurs over the entire length of time during which the N supernovae explode, $\tau_{\rm SN}$. This yields for the evaporated mass during the snowplow phase

$$M_{\rm ev,sn} = 1.65 \times 10^3 h_{100}{}^{2/7} L_{38}{}^{5/7} \tau_{\rm SN,6} M_\odot . \tag{11}$$

The total mass is the sum of the two contributions in equations (9) and (11). For Galactic parameters, $M_{\rm ev,sn} \approx 2 M_{\rm ev,bb}$. Both these equations overestimate the evaporated mass if condensation of hot gas onto molecular clouds in the supercavity is important.

Second, we treat the alternative possibility, that most of the evaporated gas comes from diffuse clouds that remain inside the supercavity. Diffuse clouds should have been destroyed by ionizing photons before the O stars become SNs, but under some circumstances, e.g., shielding by other clouds or molecular clouds, they may survive. In this case, the increased evaporated mass lowers the cavity temperature and makes drastic differences. Treating evaporation in detail is rather complicated and requires assuming a geometrical form for the clouds. Theoretical treatments assume either spherical (Cowie and McKee 1977) or ellipsoidal (Cowie and Songaila 1977) clouds. The theoretical papers find that evaporation is so efficient that conventional spherical clouds would be destroyed by evaporation during $\tau_{\rm SN}$ (McKee and Cowie 1977). However, real clouds are filamentary or sheetlike (Kalberla, Schwarz, and Goss 1985; Kulkarni and Heiles 1986*a*, *b*). Furthermore, the theories may not be correct, since they do not appear to explain the fact that O VI is observed in only $\sim 10\%$ of the cloud interfaces (de Jong 1980; Cowie and Songaila 1986). We proceed by considering the extreme case in which all the gas in clouds is evaporated; then the total mass evaporated from clouds is

$$M_{\mathrm{ev,cl}} = 2\pi R_f^{\,2} h\mu \langle n_{\mathrm{cl}} \rangle \;,$$

where $\langle n_{\mathrm{cl}} \rangle$ is the volume density in clouds reduced by the filling factor, i.e., the density if the cloud matter were spread uniformly throughout space. This can be rewritten

$$M_{\mathrm{ev,cl}} = 6.6 \times 10^6 \langle n_{\mathrm{cl}} \rangle h_{100}^{\;4/3} L_{38}^{\;1/3} n_0^{\;-1/3} v_{\mathrm{rms}}^{\;-1} \; M_\odot \quad (12)$$

This mass estimate is an upper limit for two reasons: not all the cloud matter will evaporate, and some condensation onto molecular clouds will occur. Assuming that this gas expands at constant pressure, we would have

$$T_{\mathrm{ev,cl}} = 700 L_{38}^{\;2/3} \langle n_{\mathrm{cl}} \rangle^{-1} h_{100}^{\;-4/3} n_0^{\;1/3} v_{\mathrm{rms}} \tau_{\mathrm{SN,6}} \; \mathrm{K} \; . \quad (13)$$

For realistic values of the parameters, $T_{\mathrm{ev,cl}} \approx 10^5$ K, which is so low that two additional considerations enter: first, the temperature is so low that the evaporation may not be fast enough to destroy all of the clouds; and second, the cooling rate at this temperature is rapid, so that the gas could not remain at this temperature for the duration of the superwind, τ_{SN}. In this case, the temperature would involve an equilibrium between energy input from the wind, cooling by heavy-element excitation in the gas, cooling by expansion, and cooling from putting energy into evaporating cold matter from cloud surfaces. Furthermore, the evaporation rate is affected to an unknown degree by possible inapplicability of the theory and by magnetic fields in the clouds. Evidently this is a complicated problem, far beyond the scope of the present paper. Furthermore, in a SN-dominated medium it is quite possible that clouds suitable for evaporation (diffuse H I clouds embedded in the hot intercloud medium) do not actually exist. Such clouds tend to be broken up and destroyed by SN shocks (Heathcote and Brand 1983) so that, with a sufficiently high SN rate, they will exist only if their rapid destruction is balanced by a corresponding formation process, as envisioned by MO. The small fraction, 2/13, of high-velocity UV absorption components toward OB associations (Cowie *et al.* 1981) suggests that clouds are ineffective in producing radiative cooling and is consistent with the idea that clouds are destroyed by ionizing radiation of the OB stars before supernovae begin to explode. Thus not only is the problem complicated, but clouds may exist only to a degree depending on the degree of supernova domination.

Mass input to the halo must be balanced either by matter falling back to the galactic plane, as in the fountain models of Shapiro and Field (1976) and Bregman (1980*a*), or by a galactic wind. If there is no cooling or heating in the halo itself, Chevalier and Oegerle (1979) have shown that a wind will ensue if the gas is injected with temperature and velocity sufficiently large to enable it to climb out of the galactic potential well. Most of the internal energy of the hot supercavity gas is thermal, so that the temperature of the gas (plus the kinetic energy from the galactic rotation) is a good indicator of its ability to climb out of the gravitational well. Whether or not radiative cooling occurs depends on the temperature to which the gas is originally heated; we discuss this matter below in § VI.

III. VALUES OF PARAMETERS

a) The Gaseous Disk

To apply the above ideas to a galaxy we need to know the relevant parameters of the gaseous disk. These parameters are difficult to obtain for our own Galaxy and essentially impossible for external galaxies. Thus, we discuss the data for our Galaxy and adopt reasonable values for the relevant parameters as an example. The adopted value of all parameters are given in Table 1.

There is a serious question regarding the identity of the ambient medium in which a SN explodes. If $Q \lesssim 1$, then this medium can be taken as the classical intercloud H I, observational evidence for which abounds; see Kulkarni and Heiles (1986) for a review. However, if $Q \gtrsim 1$, then the HIM fills most of the volume and the randomly distributed Type I SNs explode mainly in the HIM, which produces only small increases in Q. However, correlated Type II SN precursors reside in the relatively uniform H II gas produced by the homogenizing effect on H I clouds by the OB stars' ionizing photons (Elmegreen 1976; McKee, van Buren, and Lazareff 1984); in this case the z-distribution of the ambient gas should mimic that of the clouds. The proper treatment of this question would make relatively minor quantitative changes but greatly increase the complication. Therefore, we simply concentrate on the case $Q \lesssim 1$, which appears to be favored on observational grounds for the solar circle and beyond.

The observed properties of the classical intercloud gas are difficult to determine because it is difficult to distinguish between the warm intercloud gas and the colder "cloud" gas. Results have also been confused by the fact that some authors have not properly distinguished between the intercloud and cloud gas, and also by an incorrect (too small) scale height reported by Baker and Burton (1975) and Burton (1976), which has been quoted in several subsequent papers by other authors (see Heiles 1980).

The total intercloud H I column densities in the z-direction (through the full thickness of the disk) quoted by most authors agree rather well at ~ 90 cm^{-3} pc. This amounts to $\sim 38\%$ of the total H I column density given by Heiles (1976). We assume this fraction applies everywhere in the Galaxy, even though its observational determination refers only to the solar neighborhood. The total H I column density (not just the intercloud component) is essentially constant from Galactocentric radii 5–20 kpc (Kulkarni, Blitz, and Heiles 1982; Lockman 1984); thus we take the intercloud column density to be constant, too. Velocity dispersions also agree rather well, ranging from ~ 8 to 9 km s^{-1}. The average heights, $\langle |z| \rangle$, also agree rather well and average ~ 185 pc in the solar neighborhood, the value given by Falgarone and Lequeux (1973). The thickness of the total H I layer is essentially constant inside the solar circle, and rises dramatically outside (Kulkarni *et al.* 1982; Lockman 1984). We assume the height of intercloud H I is everywhere proportional to that of all H I. This gives the variation with Galactocentric radius R_{Gal} as approximately

$$\langle |z| \rangle \approx \begin{cases} 185 \,, & R_{\mathrm{Gal}} < 11.5 \text{ kpc} \\ 185\left(1 + \dfrac{R_{\mathrm{Gal}} - 11.5}{3.8}\right), & R_{\mathrm{Gal}} \geq 11.5 \text{ kpc} \,. \end{cases} \quad (14)$$

We need suitable values of n_0, the intercloud density, and h, its z-height, for use in the equations given in § II. The theory behind those equations assumes n_0 is independent of distance from the center of explosion, which it clearly is not because of the gradual falloff of density with z. The shape of the actual z-distribution is uncertain. Given this, plus the absence of a detailed theory in any case, we adopt the simplest choice: the column density equals its observed value and h equals $\langle |z| \rangle$. This gives

$$n_0 = \frac{90}{2\langle |z| \rangle} \text{ cm}^{-3} . \quad (15)$$

Thus $n_0 = 0.24$ inside the solar circle. Note that because Type II SNs occur in spiral arms, while the values of intercloud density considered here are averaged between arm and interarm regions, the true value of n_0 in the spiral arms, where superexplosions occur, is probably larger than that given by equation (15).

Lockman (1984) and Lockman, Hobbs, and Shull (1986) have recently found that the z-distribution of the warm H I has a very low-density tail extending to high $|z|$. This high-z component may affect the details of the breakthrough process, perhaps invalidating our simple treatment of using the single parameter h to describe the thickness of the gas layer. This question deserves theoretical attention.

The pressure of the intercloud medium is very difficult to measure directly. The simplest procedure is to assume it is equal to the pressure in the clouds, which can be measured from UV observations of C I (Jenkins, Jura, and Loewenstein 1983). However, this neglects the reasonable possibility, suggested by Cox (1981), that the gas pressure in the intercloud medium is somewhat higher than in the clouds because the magnetic field is smaller. We adopt the simple approach of ignoring this possibility; thus our adopted intercloud pressure is likely to be too low. This may not be all that serious because the measured pressures are anything but constant. We adopt the "representative intermediate pressure" of Jenkins *et al.* (1983), $p_{04} = 0.40$. This, together with the adopted value $n_0 = 0.24$, gives $(p/\rho)^{1/2} = 9.9 \text{ km s}^{-1}$, close to the observed value of the intercloud H I velocity dispersion.

Consider for a moment the variation of the ISM parameters with Galactocentric radius R_{Gal}. The z-component of the gravitational field increases toward the inner Galaxy because the mass of the stellar disk decreases exponentially with R_{Gal} (see § IV*a*). Between the solar circle at 10 kpc and the inner edge of the H I disk at 5 kpc, the disk stellar density increases by a factor of ~ 3.1 and the z-column density of H I is constant. The increased gravitational force toward the interior raises the gas pressure at $z = 0$. This would lead to a thinner disk for all H I (cloud + intercloud). However, *observationally* the total H I z-distribution is independent of R_{Gal} (Lockman 1984). This can be reconciled with theory only if the total interstellar pressure increases toward the interior, so as to combat the increased z-component of gravitational force. If the extra pressure arises solely from gas pressure, v_{rms} would have to increase by a factor of 1.8. This is not inconceivable and violates no observational constraint known to the author. Alternatively, the extra pressure might be produced by magnetic fields or cosmic rays, or a combination of both, with gas pressure. In any case, it is remarkable that all these conspire to keep the total thickness of the H I disk independent of R_{Gal}. It therefore seems inevitable that v_{rms} must increase inside the solar circle, by a factor of 1.8 or less. We neglect this dependence in the ensuing discussion, but it is worth keeping in mind because it introduces yet more quantitative uncertainties.

b) SN Explosion Energy

Estimates given in the literature for the SN explosion energy, E, are based on analysis of observed remnants and depend on the details of theoretical interpretation and the assumed distances. The distance uncertainty is removed by considering extragalactic SN remnants, but the interpretive difficulties remain. Blair, Kirshner, and Chevalier (1981) find that the derived energies for a large sample of extragalactic SN tend to increase with the observed remnant diameter; this clearly indicates an interpretive difficulty. The derived energies should be more reliable for larger remnants for a variety of reasons, and the derived energies for the largest remnants are $\sim 10^{51}$ ergs. Thus we adopt $E_{51} = 1$, both for Type I and Type II SNs. There is considerable uncertainty in this number, however. Some determinations argue for energies ~ 5 times lower (e.g., Long and Helfand 1979), and determinations for the same SN remnant in our Galaxy can differ by an order of magnitude (see Chevalier, Kirshner, and Raymond 1980; Seward, Gorenstein, and Tucker 1983). Some astronomical lore says that Type I energies are smaller than Type II energies. The SN energy values need to be resolved definitively.

IV. UNCORRELATED SNS VERSUS THE ISM

a) Type I SNs

Tammann (1982) has reviewed the statistical properties of SNs, and the statistical properties used herein are derived from his discussion. His SN rates for the Galaxy are derived from a combination of rates in external galaxies and historical SNs in the Galaxy. The former is uncertain because of the uncertain degree of selection effects. For example, the extragalactic SN rates are multiplied by a factor of 2.8 to account for obscuration in inclined spiral galaxies, and there is some controversy concerning this correction. The Galactic rate is uncertain because of small-number statistics.

Historically, Type I SNs have been connected with Population II stars. However, in recent years the connection has become less clear. Type I SNs occur in all types of galaxies, including elliptical galaxies in which their distribution follows the light. This implies that they arise from an old stellar population. In spiral galaxies Type I SNs concentrate toward the disk and also toward the center (Tammann 1977*a*). The Type I SN rate increases with the presence of gas, young objects, and blue color in spiral galaxies. This implies that they arise from Population I stars. However, they are not concentrated in spiral arms (Maza and van den Bergh 1976). The conflicting and contradictory evidence has been well summarized by Trimble (1982), who concludes that we have little certain knowledge concerning the progenitors of Type I SNs.

We give high weight to the fact that Type I SNs are distributed as the light in E galaxies. We thus assume that Type I SNs arise from an old stellar population and are distributed as the Population I and old Population II stars in spiral galaxies. They are uncorrelated in space and time. The distribution of these uncorrelated Type I SNs within a galaxy is important: if enough of them explode within the gaseous disk, the disk will be SN-dominated as in the uncorrelated SN models of MO and Cox.

We need, therefore, to establish the radial distribution and scale height of the disk population of Type I SNs. This cannot be done directly from Type I SNs because statistics are limited. In our own Galaxy, there are only two SNs that are generally considered Type I: SN 1572 (Tycho) and SN 1604 (Kepler), with z's of 98 and 474 pc, respectively (Tammann 1982). The exponential scale height $h_{\text{SN I}}$ from this small sample of two would be just the average, 286 pc, but the uncertainty is extreme, to say the least. Furthermore, Kepler has a z-velocity of ~ 100 km s^{-1} (van den Bergh and Kamper 1977, with Tammann's [1982] distance of 4 kpc), which would make its time-average z-height considerably higher. In external galaxies, Tammann (1977*b*) finds the average z-distances of Type I SNs in edge-on spirals to be 700 pc; however, this result is from a sample of only six and, furthermore, is actually an upper limit because of selection effects.

We take the straightforward approach of assuming that Type I SNs are distributed as the light of the disk population and adopt the distribution of old Population I and Population II stars given by Bahcall and Soneira (1980). This has an exponential scale height of 325 pc, ~ 1.8 times the scale height of H I in our Galaxy. If $h_{\rm SN\,I}$ is really larger, then Type I SNs have a smaller effect on the ISM of the Galaxy than calculated below, and a larger heating effect on the gaseous halo. We regard the Type I SN scale height as very uncertain.

For the radial distribution of Type I SNs in spiral galaxies, the surface brightness of the stellar disk population decreases exponentially with radius. Bahcall and Soneira (1980) adopt a radial exponential scale length of 3.5 kpc and a solar distance of 8.0 kpc. We multiply both by 1.25 to yield 4.4 and 10 kpc, respectively, since 10 kpc is used for the solar Galactocentric radius throughout this paper. This scale length is fully consistent with values quoted by others (e.g., de Vaucouleurs and Pence 1978; Tammann 1982).

Our Galaxy has a "hole" in the H I distribution, with a precipitious drop in surface density within 5 kpc Galactocentric radius. This is the region where many Type I SNs occur; with our adopted stellar mass distribution, 31% of the Type I SNs occur within radius 5 kpc. The rest explode randomly outside radius 5 kpc, and only 1/1.8 of these explode within the intercloud H I layer. We suppose that the others, which explode outside the layer, dissipate their energy in the halo and do not affect the H I. Thus, only 38% of the Type I SNs affect the H I disk. The remainder act as a heat source for the gas in the halo.

b) *Runaway O Stars*

Below, in § V*b*, we discuss runaway O stars. There we conclude that fewer than one-fourth of all OB stars are runaways moving too fast to be cluster members when they explode. These runaways add to the uncorrelated SN rate, by an amount $\sim \sigma/4$ or less, where σ is the Type II SN rate; σ is derived in § V and presented in Figure 1. These are the highest velocity portion of the runaways, with velocities ~ 60 km s^{-1}. During their ~ 7 Myr lifetimes, the stars typically travel a total distance of ~ 400 pc and a z-distance of ~ 300 pc. With $h = 185$ pc, they explode outside the H I layer. Thus they do not contribute appreciably to Q_{3D} as uncorrelated SNs. Rather, their energy goes directly into heating the halo. This effect is discussed in § VI.

c) *Type I SNs versus the* H I *Disk*

In our own Galaxy, Tammann (1982) estimates the total frequency of Type I SNs to be 0.027 per year (interval 36 yr). This normalizes the functional dependences discussed above and gives

$$S_{-13} = 3.4 \frac{325 \text{ pc}}{h_{\rm SN\,I}} \exp\left(-\frac{R_{\rm Gal}}{4.4}\right) \tag{16}$$

and

$$Q_{3D} = 3.5 \frac{325 \text{ pc}}{h_{\rm SN\,I}} \exp\left(-\frac{R_{\rm Gal}}{4.4}\right). \tag{17}$$

Thus, for example, $Q_{3D} = 1.1$ at $R_{\rm Gal} = 5$ kpc, and 0.36 at the solar circle (assumed $R_{\rm Gal} = 10$ kpc). In the solar vicinity, the

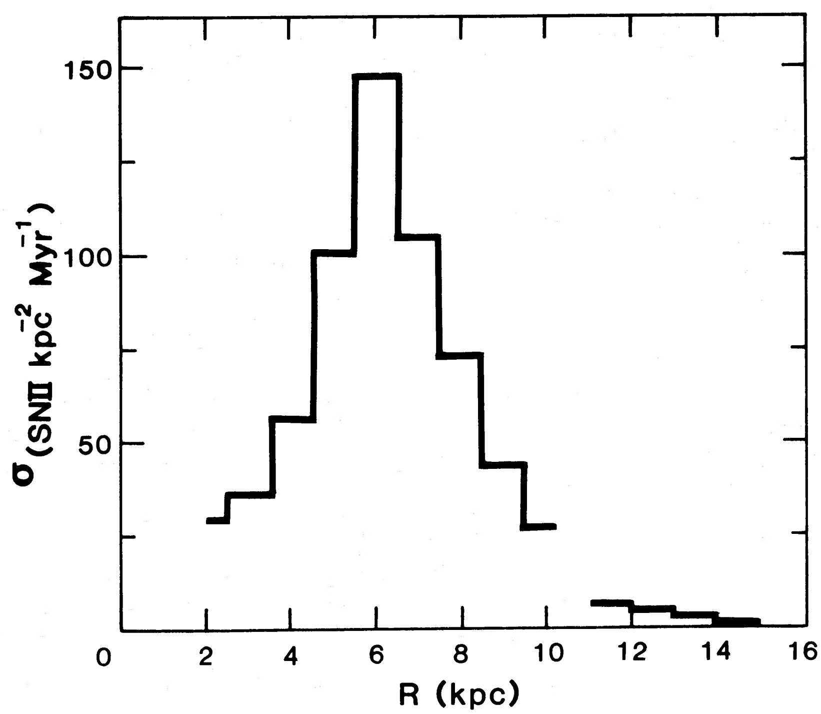

FIG. 1.—Frequency of the azimuthally averaged areal density of individual Type II SNs vs. Galactocentric radius

fraction of interstellar volume occupied by the HIM is 0.26. This is in reasonable agreement with the value of ~ 0.5 estimated from observational data by Kulkarni and Heiles (1986).

The fact that Type I SNs make Q increase toward the Galactic interior is intriguing. It suggests that Type I SNs offer a natural explanation for the H I "hole" inside 5 kpc—that outside 5 kpc the H I layer exists precisely *because* the disk is *not* dominated by Type I SNs. Inside, the H I disappears because SNs do dominate. Additional energetic agents that are concentrated toward the center disturb the H I inside, too—specifically, the SN-heated wind from the stars in the central bulge (Bregman 1978, 1980*b*).

The increase of Type I SNs toward the center of spiral galaxies, and the central bulge of stars, are universal characteristics. Thus, if H I is destroyed in a highly SN-dominated ISM, then decreases in H I surface density should be observed toward the centers of many spiral galaxies. Such decreases are, in fact, observed, particularly when the H I surface density is compared to the total mass density, and particularly for spirals with large central bulges (see Bosma 1981*a*, *b*, and quoted references).

However, the *total* gas distribution (including H_2) does *not* tend to decrease toward the center in most spirals. For example, in IC 342, NGC 6946, and M51 (Young and Scoville 1982*a*; Scoville and Young 1983), the H_2 surface density (derived from CO observations) follows the exponential increase of the light toward the center. However, in these galaxies the H I surface density *does* decrease toward the center, just as it does in our Galaxy. This circumstance is precisely what would be expected for a Type I SN–dominated ISM: the independent, random individual explosions would have little effect on the dense molecular clouds (Shull 1980; Wheeler, Mazurek, and Sivaramakrishnan 1980), but would heat the more diffuse H I and produce a HIM–dominated medium in the spirit of MO and Cox.

V. CORRELATED TYPE II SNS AND THE FORMATION OF SUPERCAVITIES

a) The Correlation of Type II SNs in Space and Time

Again, we adopt Tammann's (1982) SN statistics, realizing that they are uncertain and controversial. Type II SNs occur only in Sab or later galaxies having young stellar populations and concentrate (but not exclusively) along spiral arms (Maza and van den Bergh 1976); thus, they arise from Population I stars. The concentration to spiral arms shows observationally that stars explode before enough time has elapsed for them to pass through spiral arms. This time is of order 40 Myr for the Galaxy. It is fully consistent with deductions from stellar evolution theory, which imply a lower mass limit for the progenitors of $\sim 8\ M_\odot$ (Wheeler 1981); such stars have lifetimes of less than ~ 30 Myr (26 Myr for 9 $M_\odot$; Iben 1966), smaller than 40 Myr.

We adopt $\tau_{\rm SN} = 30$ Myr (defined just after eq. [10]). Within this time, the random velocities of stars within associations cause the stars to occupy a large region. Random velocities are typically quite small; 5 km s^{-1} may be an overestimate for a typical cluster (Blaauw 1964). At 5 km s^{-1}, the stars occupy a region 150 pc in radius after 30 Myr; 150 pc is smaller than R_f, the final radius of the cylindrical supercavity. Thus, it is valid to consider the N individual SNs to act in concert as a single superwind. A fraction of the Type II SNs are runaways and are not correlated with clusters; as discussed above in § IV*b*, they heat the gaseous halo.

The concentration of Type II SNs to spiral arms means that superwinds are not distributed randomly in the disk but rather are themselves concentrated to spiral arms. The parameter s, in equation (8), accounts for this concentration. From Maza and van den Bergh (1976), we estimate that 43% of all Type II SNs occur within spiral arms. Spiral arms occupy perhaps 20% of the area of a typical galactic disk. For these values, $s = 2.1$ inside spiral arms and 0.7 outside. Thus, for individual pieces of gas the SN rate is time dependent. The quantity $Q_{\rm 2D}$, given by equation (8), was derived assuming time independence. Matter entering a spiral arm is subject to a sudden increase in effective SN rate, and it will take several τ_p's (eq. [7]), or ~ 60 Myr, to reach a steady state. This is a bit larger than the residence time in a spiral arm, so the assumption is not quite valid. Nevertheless, we assume that equation (8) still applies for most of the matter, both inside and outside of spiral arms.

Clusters themselves are correlated *within* spiral arms. Observationally this is evident in the formation of large H II regions in spiral arms arranged like "beads on a strong" separated by ~ 1 kpc (Mouschovias, Shu, and Woodward 1974). Each "bead" should become a supercavity center. They are separated by more than the supercavity diameter, so each can be regarded as independent of the others.

b) The Value of N

Above we find that stars having $M > 8\ M_\odot$, corresponding to main-sequence spectral type B3, eventually become Type II SNs. The quantity N is the number of such stars per cluster. In Paper I, we used Bruhweiler *et al.*'s (1980) value of 200; however, we believe that this is an overestimate. We use two methods to obtain N. In one, we take derived estimates of the birthrate of stars having $N > 8\ M_\odot$ in the solar vicinity and divide by the birthrate of clusters in the solar vicinity. In the other, we use direct observations of the number of O stars in clusters in the solar vicinity together with the initial mass function derived from clusters in the solar vicinity.

The first method requires two observationally derived birthrates, that for clusters and that for stars. A recent estimate for the birthrate of young clusters and associations, 4.5×10^{-5} kpc^{-2} yr^{-1}, is given by Elmegreen and Clemens (1985). Estimates for the birthrate of young stars—i.e., the initial mass function (IMF)—are less reliable. They are affected more seriously by observational selection, because individual stars can be hidden more easily by extinction than whole clusters. Van Buren (1985) claims to have carefully accounted for these effects in his derivation of the IMF. He obtains a significantly flatter mass distribution than all other authors (see Scalo 1986 for a review). Use of a flatter mass function is also justified by the increase in flatness toward the Galactic interior found by Garmany, Conti, and Chiosi (1982). Van Buren derives a formation rate of 2.5×10^{-5} kpc^{-2} yr^{-1} for $M > 8\ M_\odot$. As noted below in § V*c*, this is in excellent agreement with our estimate of the Type II SN rate in the solar vicinity, obtained on entirely different grounds. The ratio of the cluster and stellar birthrates gives $N = 56$. (As an indication of the uncertainty, a typical, steeper IMF used by Ostriker, Richstone, and Thuan [1974] gives $N = 24$.) Of these 56 stars, roughly half will have left the cluster with typical velocity 30 km s^{-1} as "runaways" (Stone 1979, 1981). Thus we expect the number of stars *observed* in a typical cluster to be only ~ 28. If Stone has overestimated the fraction of runaways, as suggested by

Garmany and Conti (1986), then the number of observed stars should be larger.

The second method uses the observed number of O stars in clusters, together with the IMF for clusters, to estimate N. This was the method used previously by Bruhweiler *et al.* (1980), who began not with the observed number of O stars but with the number of stars B0 and earlier listed by Humphreys (1978). However, Humphreys's B stars are both incomplete and not restricted to stars earlier than B0, which (we believe) caused Bruhweiler *et al.* to overestimate N. The largest uncertainty in this method is the IMF. The IMF for clusters appears to differ from that for field stars (Scalo 1986); this is actually expected, since it is primarily the more massive O stars that tend to be "runaways" (Stone 1979, 1981). This effect should steepen the mass function for stars that lie in clusters. Scalo (1986) derives a slope of -1.43 (cf. -1.03 for van Buren) for clusters. We adopt this value. We adopt 26 $M_\odot$ for the mass of an 09 star (e.g., Ostriker *et al.* 1974; Bowers and Deeming 1984), so that the ratio of the number of stars having $M > 8\ M_\odot$ to those having $M > 26\ M_\odot$ is equal to 6.1 (cf. 9.0 for the steeper mass function of Ostriker *et al.* [1974]). Thus, in the Sco OB 1 association, the 18 O stars listed by Humphreys correspond to $N = 110$. However, the Sco OB 1 association is not typical. Humphreys's Table 5 gives 301 O stars in 71 clusters for an average of 4.24 O stars per cluster; this translates into $N = 26$. This is very close to the number per cluster of 28 found in the above paragraph with the first method.

We emphasize that both methods suffer considerable uncertainty. The slope of the IMF is not well determined. In addition, its parameters vary from cluster to cluster. For example, the Pleiades has B stars but no O stars and because it is surrounded by filamentary dust no SNs have yet exploded; apparently its IMF has a maximum mass that is smaller than the 8 $M_\odot$ required for a SN! The slope of the IMF also varies, at least with position as mentioned above and probably from cluster to cluster as well. We have indicated the quantitative uncertainties above, but there may be other contributions that we have not discussed. These uncertainties concerning the IMF are well known, and the problem of understanding the IMF remains an important and difficult one.

The result $N = 26$ or 28 should be an underestimate because it does not include the "runaway" stars, most of which are O stars with short lifetimes and many of which will deposit energy within the supercavity, defined by R_f in equation (6) above. According to Stone, the velocity dispersion of the runaways is ~ 30 km s^{-1} or greater. We estimate that about one-half of these will release their energy within distance R_f. Thus we adopt $N = 40$. Stone probably overestimated the fraction of runaways, because his estimate was based on proper motions, which have large errors (Garmany and Conti 1986). If so, our adopted value of N is somewhat too low.

It is conceivable that we are seriously underestimating the value of N, either because OB associations may themselves be clustered or because only the dense core of associations can be easily recognized with the Galaxy. This latter possibility is suggested by van den Bergh (1965), who finds that the median diameter of OB associations in M31 is 400 pc—5 times that in our Galaxy. Similarly, he obtains an association birthrate in M31 of only 10^{-5} yr^{-1}, far smaller than the Galactic rate. Assuming a Type II SN rate of 0.023 yr^{-1} for M31, as we have assumed for the Galaxy, this implies $N = 2300$! With the large diameter of these associations, not all of these SNs could contribute energy to the supercavity, and the details of our derivation of Q would be incorrect.

Is there any evidence for a variation of N with Galactocentric radius? We think not, on the basis of both Galactic and extragalactic data. The only way to estimate N for clusters too distant to observe optically is through the Hα or radio emission of the associated H II regions. These luminosities are proportional to the number of ionizing photons emitted by the exciting stars, which is proportional to N if the H II regions are ionization-bounded. In NGC 628, which to our knowledge is the only external galaxy for which data exist, Kennicutt and Hodge (1980) find that the H II region luminosity function is independent of radius. In our Galaxy, Smith, Biermann, and Mezger (1978) have compiled radio luminosities for a statistically complete sample of H II regions within the Galaxy. A plot of the radio luminosity versus Galactocentric radius reveals no obvious dependence of the luminosity on radius. One can go further and predict the number of O stars per cluster in these regions. This involves using the mass function for clusters, together with the dependence of ionizing photon luminosity on stellar mass given by Panagia (1973), to estimate the number of O stars for a given radio luminosity, and then deriving the average value of N for the sample of Smith *et al.* This procedure is very sensitively dependent on the assumed value of the high mass cutoff because the photon luminosity increases drastically with stellar mass and the number of stars decreases with stellar mass. Assuming an upper mass cutoff of 90 $M_\odot$, we obtain

$$\text{Number of O stars} = \frac{\text{Number of ionizing photons s}^{-1}}{6.6 \times 10^{49}} . \tag{18}$$

Applying this to the sample yielded 40 O stars, in reasonable agreement with the 26 obtained in the above paragraph. Thus we conclude that our runaway-corrected result $N = 40$ applies throughout the Galaxy.

These 40 stars yield energy in the form of stellar winds and SNs. We assume that all SNs yield 10^{51} ergs and all O stars yield an additional 10^{51} ergs in wind energy. This is an approximation, because the wind luminosity varies with optical luminosity (Abbott 1982). With this approximation we obtain an average of 1.17×10^{51} ergs per star. We further follow McCray and Kafatos (1986) and assume that the combined power output of winds and SNs is uniform during the period $\tau_{\rm SN}$; in fact, the power probably is larger during the latter stages when the more numerous less massive stars explode as SNs. With these assumptions and with $\tau_{\rm SN} = 30$ Myr,

$$L_{38} = 1.24 \times 10^{-2} N . \tag{19}$$

Thus, our adopted value for L_{38} is 0.5.

We note parenthetically that Smith *et al.*'s sample contains five H II regions with anomalously bright radio luminosities, with N ranging from 120 to 330. They are randomly distributed in Galactocentric radius between 5.1 and 10.3 kpc.

c) *The Type II SN Rate as a Function of Galactocentric Radius*

In this section we derive the relative distribution of Type II SNs within the Galaxy from the distribution of related populations: pulsars, which are formed from SNs, and molecular clouds, which form the stars that eventually becomes SNs.

Tammann (1982) gives the total Galactic Type II SN rate as 0.023 yr^{-1}, corresponding to an interval of 44 yr. This should agree with the pulsar birthrate, because neutron stars are thought to be formed primarily by Type II SNs. Recent estimates of the Galactic pulsar birthrate have decreased for several reasons and now range from 0.008 to 0.033 yr^{-1} (Lyne, Manchester, and Taylor 1985). The two rates are commensurate, although Tammann's SN rate lies near the high end of the pulsar range. This rough agreement supports our use of the Galactic distribution of pulsars to derive the Galactic distribution of Type II SNs.

It is conceivable that the Type II SN rate is very much smaller. First, and perhaps trivially, circumstantial evidence favors a very much smaller rate. No SN has been known in our Galaxy since Cas A, which exploded some 300 yr ago. And in M31, only one SN has been observed in astronomical history. Second, we can estimate the number of O stars currently residing in a galaxy by measuring the total Hα or radio free-free emission. For NGC 628, Kennicutt and Hodge (1980) find $\sim 2 \times 10^{53}$ ionizing photons s^{-1}; and for our Galaxy, Mezger (1978) finds $\sim 3 \times 10^{53}$ s^{-1}. The near equality of these numbers lends some confidence to their validity, which is helpful because they are difficult to determine. From equation (18) above, this corresponds to 4500 O stars. For a main-sequence lifetime of 7 Myr, the O-star birthrate is 650 Myr^{-1}; the corresponding birthrate for stars greater than 8 $M_{\odot}$ is 0.004 yr^{-1}, i.e., one SN every 250 years. This is 5.8 times smaller than Tammann's rate. Furthermore, according to Mezger only 26% of the O stars reside in clusters; the majority ionize extended, diffuse regions in the ISM. It is difficult to assess the severity of the discrepancy: some O stars lie in dense regions and produce unobservably small H II regions, and in addition both the observational data concerning the Hα or radio emission and especially the interpretation are subject to quantitative error.

The distribution of pulsars with Galactocentric radius is given by Lyne *et al.* (1985). This distribution peaks at ~ 6 kpc radius; inside this is falls precipitiously and outside it drops gradually with radius, falling to essentially zero by ~ 14 kpc. The derived distribution is obtained from the observed distribution after generous correction for selection effects. These corrections make the derived distribution more uncertain with increasing distance from the Sun, and in particular within ~ 6 kpc Galactocentric radius.

The Type II SN distribution should also mimic the molecular distribution, because the star formation rate increases linearly with H_2 density (Young and Scoville 1982*b*). The molecular distribution with Galactocentric radius is given by Sanders, Solomon, and Scoville (1984), and it is indeed comparable to the pulsar distribution inside the solar circle and outside Galactocentric radius 2 kpc. Inside 2 kpc, there is a sharp peak in the molecular abundance; this peak is not related to star formation in the diffuse gaseous disk, and we exclude this central molecular mass in the following discussion.

Outside the solar circle, there is a curious difference between the molecular and pulsar distributions: the molecular abundance decreases with radius faster than the pulsars. This may well be a selection effect resulting from less complete observational coverage outside the solar circle, particularly for the molecular observations. CO survey observations are restricted in both Galactic longitude and latitude; the restriction in latitude is particularly serious because of the Galactic warp and the increase in thickness of the H I layer in the Galactic exterior (see, e.g., Kulkarni *et al.* 1982; Henderson, Jackson, and Kerr 1982). The data of Sanders *et al.* (1984) cover $l = -4°$ to 170°, $|b| < 2°$, which is too small to cover the Galactic exterior. A more recent survey by Dame and Thaddeus (1985) does adequately cover the exterior but of course could not have been included in the earlier analysis of Sanders *et al.* (1984).

Adopting the Sanders *et al.* (1984) molecular distribution for the shape of the Type II SN distribution, and the Type II SN frequency of 0.023 per year from Tammann (1982), the SN distribution becomes the same as the molecular distribution given in Figure 12 and Table 3 of Sanders *et al.* (1984), with an appropriate scaling factor. This factor is:

$$\nu_{\mathrm{SN\,II}} = 8.2 \times 10^{-6} \text{ Type II SN } M_{\odot}^{-1} \text{ Myr}^{-1} , \quad (20)$$

where the *mass refers to* H_2 *alone*. As mentioned above, the H_2 density may be underestimated outside the solar circle; if so, the numerical factor on the right-hand side of equation (20) is too large, but not by much.

There is considerable uncertainty in equation (20) because the H_2 mass is not observed directly. The "observed" H_2 mass depends on the adopted ratio of CO to H_2, and some authors quote different (especially lower) H_2 masses (see Sanders *et al.* 1984, and quoted references; Blitz and Shu 1980; Gordon and Burton 1976). The SN rate is fixed by other data. Thus, if different values for the total Galactic H_2 mass or the SN rate are adopted, the numerical factor in equation (20) must be adjusted accordingly.

The radial distribution of the Type II SN rate per unit area in the disk is shown in Figure 1. This radial dependence peaks sharply at 6 kpc. The Type II SN rate in the solar vicinity is $\sim 3 \times 10^{-5}$ kpc^{-2} yr^{-1}. This should agree with the SN birthrate, i.e., the birthrate for stars having $M > 8\ M_{\odot}$. Van Buren's (1985) birthrate for such stars, which we adopted above in § V*b* to estimate N, predicts 2.5×10^{-5} kpc^{-2} yr^{-1}—in excellent agreement. On the other hand, Kennicutt's (1984) rate differs substantially if his results for a "shallow" IMF are used; this is an indication of the uncertainty in both the SN rates and the IMF in our Galaxy.

d) Type II SNs versus the H I *Disk*

We now use the characteristic values of parameters for the Galaxy to determine the porosity and mass transfer rates to the halo. From Table 1 and equation (8),

$$Q_{2\mathrm{D}} = 8.3 \sigma s N^{-1} . \quad (21)$$

Figure 1 shows that for $N = 40$ the ratio σN^{-1} peaks to $\sim 3.75s$ at $R_{\mathrm{Gal}} = 6$ kpc, dropping to less than $0.75s$ at the solar circle. These yield $Q_{2\mathrm{D}}$-values for (inside, outside) spiral arms of (65, 22) at 6 kpc, and (13, 4.4) at the solar circle. These translate into volume filling factors for hot gas of (0.98, 0.96) and (0.93, 0.81), respectively.

Note that the values of Q predicted for superwinds are considerably larger than values that would be predicted if the SN were randomly distributed. With $h = 185$ pc, $S_{-13} = 0.027\ \sigma$. At $R_{\mathrm{Gal}} = 6$ kpc, S_{-13} would be ~ 3.8 if Type II SNs were uncorrelated. Equation (1) would give $Q_{3\mathrm{D}} = 3.9$—17 times smaller than for correlated SNs!

The value used for N is crucial. In Paper I, we used $N = 200$ and obtained much smaller values of $Q_{2\mathrm{D}}$. If N were smaller than the value of 40 used here, $Q_{2\mathrm{D}}$ would be even larger.

Outside the solar circle, σ decreases and h increases, causing Q to decrease. As R_{Gal} increases, h increases to the point where breakthrough cannot occur. At this point, large individual quasi-spherical cavities are formed instead of tightly packed cylindrical ones. These "supershells" have been observed by Heiles (1979), particularly in the outer Galaxy; an approximate theory has been given by Bruhweiler *et al.* (1980). In contrast, in the inner Galaxy the cavities tend to be open at the top (the "worms" of Heiles 1984), corresponding to breakthrough.

It is interesting to consider the number of cylindrical cavities in existence at any one time. Roughly, this is just the cluster formation rate multiplied by τ_p, or $\sim$12,000. Solomon, Sanders, and Rivolo (1985) find that in the region bounded by Galactic longitude 20°–50° there are $\sim$2200 CO emission centers. The total number of centers in the Galaxy should be perhaps 4 times this number, or $\sim$9000. The approximate equality implies that, on average, each CO emission center produces a star cluster.

The large values of Q_{2D} found above are in violent disagreement with all observational data, both Galactic and extragalactic, known to the author. In the Galaxy outside $R_{Gal} = 8$ kpc, Heiles (1980) estimated that no more than 20% of the interstellar volume is occupied by large supercavities, corresponding to $Q_{3D} = 0.25$. While his estimate is subject to considerable quantitative uncertainty because of interpretive difficulties, it can hardly be incorrect by more than one order of magnitude. Similarly, the relatively smooth angular distribution of H I in the solar vicinity (Jahoda *et al.* 1985; Lockman *et al.* 1986) is inconsistent with large values of Q. Finally, the ISM into which the Cygnus Loop is expanding is not the hot, rarefied gas inside a supercavity (Hester and Cox 1985). All this shows that the large, rarefied cavity in which the Sun is embedded is atypical.

In external galaxies, Q_{2D} can be estimated from high-resolution maps of the 21 cm line emission. In his detailed mapping and analysis of M31, Brinks (1984) observed large H I holes in the disk and found $Q_{2D} \approx 0.01$. While M31 may well have a lower SN rate than the Galaxy, it can hardly be three orders of magnitude lower. The H I in M81 and M101 has also been mapped with high angular resolution (Rots 1975; Allen and Goss 1979). These data have not been statistically analyzed as they have been in M31, but from simple inspection it is clear that, as in M31, H I holes occupy only a small fraction of the disk area. An additional rough statement about Q can be made because, in all face-on spirals for which adequate data exist, H I spiral structure is correlated with the optical spiral arms. These galaxies include M81, M101, UGC 2885 (Roelfsema and Allen 1985), NGC 628 (Shostak and van der Kruit 1984), and NGC 1058 (van der Kruit and Shostak 1984). If Q were as high as estimated above, the H I could exist only in small clouds. Given the observed 21 cm line intensities, the clouds would have significant 21 cm line opacity, so that the observed line intensity would reflect the number of clouds rather than the average column density. It is hard to believe that prominent spiral arms would appear in 21 cm line maps under these conditions.

Statistical data on giant H II regions in the Sc spiral galaxy NGC 628 (Kennicutt and Hodge 1980) offer another approach for estimating Q. These giant H II regions are produced by the same OB associations and clusters that produce supercavities. The area of the disk occupied by a giant H II region is larger than the area occupied by the corresponding supercavity, because the Strömgren radius is larger than R_f. The giant H II regions last a shorter time than the supercavities. We can multiply the fractional area of the disk occupied by giant H II regions by the appropriate factors to derive the fractional area occupied by supercavities. The quantitative result depends on the minimum diameter of the giant H II regions included in the analysis, because the number of H II regions increases rapidly with decreasing size. We adopt the minimum diameter that would be produced by a cluster with $N = 12$, which is the number of SNs required for breakthrough, and estimate that 6.1% of the disk area is occupied by H II regions near $R_{Gal} = 6$ kpc, where the surface density of H II regions is largest (Kennicutt and Hodge 1976). This corresponds to $\sim$35% of the area being occupied by supercavities, i.e., $Q_{2D} \approx 0.5$.

In summary, all data on both the Galaxy and external galaxies are inconsistent with the large values of Q predicted above.

VI. SNs VERSUS THE HALO

The mass input rate to the halo is derived by multiplying the mass evaporated in each supercavity by the cluster formation rate. To account for runaways, we take the cluster formation rate, Σ, to be $0.75/N$ times the Type II SN rate; owing to our probable overestimate of the fraction of runaways, Σ is probably somewhat larger than this value. Applying equations (9) and (11), we obtain

$$\dot{M}_{\rm halo} = 72\Sigma N^{-2/7} h_{100}{}^{2/7}(26 N^{-1/3} n_0{}^{1/3} h_{100}{}^{5/3} + \tau_{{\rm SN},6})\, M_\odot\ {\rm yr}^{-1}\ . \quad (22)$$

For Galactic parameters, $\dot{M}_{\rm halo} = 22\ M_\odot\ {\rm yr}^{-1}$. This is a very large value: during the lifetime of the Galaxy, the entire mass of interstellar gas would have been recycled into the halo many times. Modern observational estimates of the inflow of H I in high-velocity clouds are much lower, ranging from 0.2 (Mirabel and Morras 1984) to 2.1 $M_\odot\ {\rm yr}^{-1}$ (Kaelble, de Boer, and Grewing 1985). (These estimates depend on the assumed distance to the high-velocity clouds, and are correspondingly uncertain). The inflow may, in fact, have nothing to do with recycled halo gas; Mirabel (1981) points out that the high-velocity gas is distributed asymmetrically with respect to the Galactic plane and believes the origin of the gas is tidal interaction between the Galaxy and the Magellanic clouds. Thus the recycling rate may be lower still, and the value predicted by equation (22) (which is, in fact, an upper limit; see § II*b*) may be much too large.

In the snowplow phase, when most of the cold gas evaporates, the gas is injected with a temperature of $\sim 1.9 \times 10^6$ K. This temperature is simply a measure of the total Type II supercavity-producing SN energy per injected particle. However, there is additional energy input to the halo from Type I SNs located outside the H I disk and, in addition, a small contribution from runaway Type II SNs. For our choices of parameters, 62% of the Type I SNs explode outside the H I layer. Retaining only the very uncertain $h_{\rm SN\,I}$ as a parameter, and including the extra heat input from runaway Type II SNs by assuming that they have the same scale height as Type I SNs, total power input to the halo is

$$\dot{E}_{\rm halo} = 1.04 \times 10^{42}\left(1 - \frac{128}{h_{\rm SN\,I}}\right) {\rm ergs\ s^{-1}}\ (h_{\rm SN\,I} > 185\ {\rm pc})\ . \quad (23)$$

We can express this as a temperature by assuming that 72% of this energy is distributed uniformly among the injected particles (as it is in a Sedov-Taylor blast wave):

$$T_{\rm halo} = 5.7 \times 10^7\left(1 - \frac{128}{h_{\rm SN\,I}}\right)\dot{M}_{\rm halo}{}^{-1}\ {\rm K}\ (h_{\rm SN\,I} > 185\ {\rm pc})\ . \quad (24)$$

Here $\dot{M}_{\rm halo}$ is in units of $M_\odot$ yr^{-1}. For $h_{\rm SN\,I} = 325$ pc, $\dot{M}_{\rm halo} = 22\ M_\odot$ yr^{-1} and $T = 1.6 \times 10^6$ K. This is slightly less energy per particle than the 1.9×10^6 K produced during evaporation in the snowplow phase. The actual temperature should be the sum, 3.5×10^6 K.

The halo gas heating can be drastic. If the total mass input is really $\sim 2.1\ M_\odot$ yr^{-1} or less, as found from high-velocity clouds, then $T \gtrsim 1.7 \times 10^7$ K. This is greater than the critical temperature for a wind in the galactic center region, 3.4×10^6 K, computed by Chevalier and Oegerle (1979) for no massive halo.

However, our Galaxy has a flat and even slightly rising rotation curve out to at least $R_{\rm Gal} = 20$ kpc (Kulkarni *et al.* 1982; Schneider and Terzian 1983), as do other spiral galaxies, which implies that there is a massive halo. A number of model mass distributions for galaxies have been developed by Bahcall, Schmidt, and Soneira (1982, 1983); here we use the model in Table 2 of the former work because it roughly matches the observationally determined rising rotation curve outside the solar circle. Inserting these into equation (5) of Chevalier and Oegerle (1979) for the critical temperature for a wind in the absence of radiative cooling, we obtain

$$T_{\rm cr} \approx 4.6 \times 10^6\left[1 + 0.28 \ln\left(\frac{10}{R_{\rm Gal}}\right)\right] \text{K} \qquad (2 < R_{\rm Gal} \lesssim 20 \text{ kpc})\,. \quad (25)$$

Most of the SNs occur outside $R_{\rm Gal} = 2$ kpc; for $R_{\rm Gal} > 2$ kpc, $T_{\rm cr} \lesssim 7 \times 10^6$ K. This is also below the temperature obtained for $\dot{M}_{\rm halo} = 2.1\ M_\odot$ yr^{-1}. For this and smaller values of $\dot{M}_{\rm halo}$, it seems likely that at least some of the halo gas will exit as a wind.

Diffuse X-ray observations place severe constraints on the Galactic halo. Nousek *et al.* (1982) find that the "halo" gas is distributed in a thick disk, with emission measure (EM) toward the Galactic poles $\sim 4 \times 10^{-3}$ cm^{-6} pc and $T \approx 2.5 \times 10^6$ K. If $h_{d,5}$ is the height of the thick disk in units of 5 kpc, then the total halo mass is

$$M_{\rm halo} = 9.8 \times 10^7 h_{d,5}{}^{1/2}\ M_\odot\,. \quad (26)$$

At $T > 2.5 \times 10^6$ K, the cooling time is

$$\tau_{\rm cool} > 9.2 \times 10^8 h_{d,5}{}^{1/2} \text{ yr}\,. \quad (27)$$

Note that this estimate for $\tau_{\rm cool}$ includes clumping effects, because the diffuse X-ray measurements respond to the rms electron density, not the average density. The gas can fall back to the gaseous disk only after it cools. Thus, the rate of infall to the Galactic plane is

$$\dot{M} = \frac{M_{\rm halo}}{\tau_{\rm cool}} < 0.11\ M_\odot \text{ yr}^{-1}\,. \quad (28)$$

This is far smaller than the mass input to the halo estimated above from equation (22). More seriously, for such a small $\dot{M}_{\rm halo}$ the injected gas should be very hot, so that $T \gg T_{\rm cr}$ and, in addition, $\tau_{\rm cool}$ is very much longer than the lower limit of equation (27). The Galaxy would necessarily have a wind.

X-ray observations of external galaxies yield similar conclusions, but with less confidence. X-ray observations of two edge-on spirals show that less than 10^{-3} of the expected SN energy is radiated as thermal X-rays (Bregman and Glassgold 1982), unless the gas temperature lies outside the range to which the *Einstein* IPC is sensitive ($\sim 6 \times 10^5$ K to 4×10^6 K). And in the face-on galaxy M101, McCammon and Sanders (1984) find that less than 0.1 of the SN energy is being radiated by a hot gas with $T > 5 \times 10^5$ K. More sensitivity and spectral range is required to establish better limits, but the trend is clear.

The Galactic situation seems to be tightly constrained. If $\dot{M}_{\rm halo}$ is large, it either violates the diffuse X-ray observations or produces a Galactic wind with an unacceptably large mass loss rate—i.e., such that the total mass lost over the Galactic lifetime exceeds the present mass of interstellar gas. The only apparent resolution is for $\dot{M}_{\rm halo}$ to be small and to be lost as a Galactic wind.

VII. SUMMARY AND DISCUSSION

We have used an approximate theory, together with observational data, to estimate Q, the fraction of volume or area occupied by supernova cavities. We have compared the results with observational determinations of Q. For Type I SNs, which explode singly and independently, the calculated values of $Q_{\rm 3D}$ are modest and do not violate observational data in the solar neighborhood. Type I SN rates should increase toward the Galactic interior, and we suggest that the resulting increase in $Q_{\rm 3D}$ is responsible for the H I "holes" observed in the centers of most spiral galaxies.

For Type II SNs, which are correlated in space and time, the explosion cavities are large enough to break through the gaseous disk. The calculated values of $Q_{\rm 2D}$ for the Galaxy are much larger than unity. This is in violent disagreement with Galactic and extragalactic observational data. Furthermore, the cavities inject mass into the halo upon breakthrough. Estimates of the amount of mass violate diffuse X-ray observations. Type I SNs, located outside the gaseous disk, should heat this mass directly; for amounts of halo mass that do not violate the diffuse X-ray observations, the gaseous halo is heated well beyond the critical temperature for a Galactic wind.

Reconciling the large discrepancy between the theoretical estimates of $Q_{\rm 2D}$ and observations requires major, fundamental modifications in either the theory or the observational input data. One theoretical possibility is a defect in the derivation of Q: Cioffi (1985) has shown that shells push neutral matter into previously cleared cavities; this effect is not included in the derivations of Q. The simplest observational reconciliation would involve decreasing the Type II SN rate, σ, and increasing the value of N, the number of Type II SNs per OB association or cluster, so that the ratio σ/N is smaller than assumed by a factor of ~ 30. This could be accomplished, for example, by adopting a 120 yr interval between Type II SNs—near the upper end allowed by the pulsar formation rate—and increasing N from 40 to ~ 300. Some increase in the effective value of N is perhaps indicated by a comparison of extragalactic and Galactic observations.

This would leave a more fundamental quandary: where does the SN energy go? In the analyses used in this paper, it must go either into mechanical motion of the ISM or into mass injection into and heating of the halo. Observations imply that much less energy goes to these sinks than is expected.

It seems that SN explosion energy that is released into these sinks must be smaller than expected. A simple and straightforward way to accomplish this would be by a decrease in SN explosion energy itself. A second, more sophisticated way involves evaporation. In our analysis of supercavities, we have totally neglected the condensation of the hot interior supercavity gas onto molecular clouds and its evaporation of cold gas from diffuse clouds. These processes decrease the pressure that drives the supershell expansion and also decrease the temperature of the gas in the supercavity interior. If they are sufficiently important, they could greatly diminish the energy available for driving the shell expansion and cause the interior supercavity gas to cool very rapidly, greatly reducing the effects on the gaseous halo. A similar process, involving the large surface area of the H I gaseous disk, might soak up the energy of the randomly distributed Type I SNs located outside the Galactic disk, removing their direct energy input into the halo gas which is responsible for the high temperatures predicted by equation (24).

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