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Coronal interstellar gas and supernova remnants

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10.1 Introduction
The study of coronal \( T \gtrsim 10^6 \) K interstellar gas is a relatively new branch of astronomy. Before the 1970s, there was little direct evidence for such gas, although theoretical models predicted that it should be found in the interiors of supernova shells. In 1956, Spitzer made the prescient suggestion that the galaxy would likely possess a hot corona much like the solar corona. By the early 1970s, a series of rocket experiments had shown that the Milky Way was glowing in soft X-rays, indicating that coronal gas was pervasive in the interstellar medium; this interpretation was supported by observations by the Copernicus satellite of the interstellar absorption line O VI \( \lambda 1035 \), showing that this tracer of high-temperature gas was extensively distributed throughout the galaxy.

We now have good maps of the brightness and temperature distribution of the soft X-ray emission from the Milky Way. With X-ray telescopes we have seen emission from coronal gas in elliptical galaxies and between the galaxies in clusters. As a result of these observations, the theory of coronal interstellar gas has advanced rapidly. The atomic processes that determine the local temperature, ionization, and spectral emissivity of the gas have been studied in detail. We have also learned much about the energy sources and macroscopic processes that control the global properties of the interstellar gas. It is now clear that the coronal gas in the Milky Way is produced mainly by the blast waves from supernova explosions, although stellar winds and compact X-ray sources may dominate in specific locales. It is also clear that the same physical processes control the properties of coronal gas wherever it may be – between galaxies and in the nuclei of galaxies as well as in the Milky Way.

Interstellar gas can be heated to coronal temperatures by three sources: energetic particles, X-rays, and shock waves. The first two heat sources may be steady; if so, the resulting state of the gas is a function of a single
parameter, the ratio of the gas pressure to the flux of particles or radiation. On the other hand, it is certainly possible or likely that the flux of particles or radiation is transient, and heating by shock waves is intrinsically transient. In those cases, the problem of interpreting or modeling the state of the gas becomes considerably more complicated.

The plan of this chapter is as follows. In Section 10.2 I describe the atomic processes that determine the local state and spectral emissivity of the gas. Then, in Section 10.3, I describe the stationary equilibrium states that result from steady sources of heating or ionization and discuss qualitatively how these results will be modified when the source is transient. In Section 10.4 I describe the physics of the interfaces between hot and cool interstellar gas, where electron thermal conduction may play an important role. Finally, in Section 10.5, I show how these concepts may be used to describe the structure and evolution of expanding interstellar shells caused by the action of supernovae and stellar winds.

### 10.2 Thermally ionized gas: coronal approximation

In the absence of ionizing radiation, the local state of hot interstellar gas is controlled by collisions of ions with thermal electrons. The dominant processes are electron impact excitation, electron impact ionization, radiative recombination, dielectronic recombination, and bremsstrahlung.

Electron impact excitation, in which an ion X\(^{z+}\) (charge state \(z\) of element X) is excited by a passing electron to a state X\(^{z+1}\) and subsequently radiates a photon with energy \(\varepsilon\), is the primary mechanism by which the heat of a coronal gas is converted to radiation in emission lines. The thermal excitation rate coefficient, \(q_e = \langle \sigma_e \nu_e \rangle\), is conventionally written

\[
q_e(T) = 8.63 \times 10^{-9} \ T_e^{-1/2} \Omega_e(T) \omega_{e}^{-1} \exp\left(\frac{-\varepsilon}{kT}\right) \text{ cm}^3 \text{s}^{-1},
\]

(10.1)

where \(T_e = T/(10^6 \text{ K})\), \(\omega_{e}\) is the statistical weight of the ground state and \(\Omega_e\) is the 'collision strength', a dimensionless function of temperature that incorporates all the details of the atomic physics. For resonance transitions \(\Omega_e\) is a weak function of temperature and has a value of order unity.

Electron impact ionization, in which \(e + X^{z} \rightarrow 2e + X^{z+1}\), has a rate coefficient that may be written (for \(kT < I_z\))

\[
q_i(T) = 1.3 \times 10^{-5} \ T_e^{1/2} I_z^{-2} r_{ij} F_{ij} \exp\left(\frac{-I_z/kT}{\text{ cm}^3 \text{s}^{-1}}\right),
\]

(10.2)

where \(I_z\) is the ionization threshold in eV, \(r_{ij}\) is the number of electrons in the outer shell, and \(F_{ij}\) is another dimensionless number \(\approx 1-3\).

Radiative recombination, in which \(e + X^{z+1} \rightarrow X^{z} + e\), has a net (to all
bound states) rate coefficient that may be written

\[ \alpha_{e,z}(T) = 2.06 \times 10^{-4}(z+1)^2T_6^{-1/2}\phi_z(T_6) \text{ cm}^3\text{s}^{-1}, \]  

(10.3)

where the function \( \phi_z(T_6) \) embodies the details of the atomic physics, and \( \phi_z(T_6) \approx 0.63(z+1)^{1/2}T_6^{-1/4} \) for hydrogenic ions. This expression provides a fairly good approximation to the radiative recombination coefficient for any ion as \( z \) becomes large.

In dielectronic recombination, a fast incoming electron simultaneously excites an ion and is trapped into a weakly bound autoionizing state, which then stabilizes by the emission of a resonance line photon, viz.: \( e + X^{z+1} \rightarrow X^z \rightarrow X^0 + \epsilon_{ij} \). The second step, needed to complete the recombination, is rare; usually the first step is followed by the reverse reaction. The rate coefficient for dielectronic recombination is given by the expression:

\[ \alpha_{d,z}(T) = 3.0 \times 10^{-12}T_6^{-3/2}B(z) \times \sum_j f_{ij}A(x)\exp(-\epsilon_{ij}/kT) \text{ cm}^3\text{s}^{-1}, \]  

(10.4)

where \( f_{ij} \) is the oscillator strength of the resonance transition \( i \rightarrow j \), \( E_{ij} = \epsilon_{ij}/[1 + 0.015x^3(z+1)^{-2}] \), \( B(z) = z^{1/2}(z+1)^{5/2}(z^2 + 13.4)^{-1/2} \), \( x = \epsilon_{ij}/(z+1) \), and \( A(x) = x^{1/2}/(1 + 0.105x + 0.015x^2) \). One can see, by comparing (10.3) and (10.4), that dielectronic recombination may dominate radiative recombination in high-temperature plasmas by factors \( \approx 10^2 \), especially for those ions with partially filled L or M shells, which have low-lying excited states and therefore can form autoionizing states with \( \epsilon_{ij} \approx kT \).

In the ‘coronal model’ approximation, the elements are ionized only by thermal electrons. The master equation describing the rates of change of the fractions, \( f_z \), of element X in each ion stage \( z \) is:

\[ \frac{df_z}{dt} = n_e \sum_{z'} M_{z,z'} f_{z'}, \]  

(10.5)

where the matrix \( M_{z,z'} \) is tri-diagonal and has elements

\[ M_{z,z-1} = q_{z-1} \]

\[ M_{z,z} = -\left(q_z + \alpha_{e,z-1} + \alpha_{d,z-1}\right) \]  

(10.6)

and

\[ M_{z,z+1} = \alpha_{e,z} + \alpha_{d,z} \]

If the ionization is stationary, \( df_z/dt = 0 \), equation (10.5) can be recast into a set of equations connecting pairs of adjacent ionization states:

\[ q_z f_z = (\alpha_{e,z} + \alpha_{d,z}) f_{z+1}. \]  

(10.7)
Note that the electron density, \( n_e \), does not appear in equation (10.7) because it is a common factor of all ionization and recombination processes. Therefore, in stationary coronal ionization balance the ion fractions, \( f_z \), are functions only of temperature. (This statement is not strictly true at high electron densities, for which \( \alpha_{d,z} \) can be suppressed.)

A potential source of confusion is the common use of the term ‘coronal equilibrium’ to describe the stationary ionization balance represented by (10.7). Equation (10.7) does not describe a state of thermodynamic equilibrium because it does not represent the result of a detailed balance of time-reverse processes (e.g., photoionization vs. radiative recombination, electron impact ionization vs. 3-body recombination). Indeed, according to the Saha equation, the ion fractions in thermodynamic equilibrium depend on electron density as well as temperature (at fixed \( T, f_{z+1}/f_z \propto n_e^{-1} \)).

Fig. 10.1 is a typical result of a coronal ionization balance calculation. Note that the temperature, \( T_e \), at which the ionization state shifts from state \( z \) to \( z+1 \) is typically about \( kT_e \approx 0.2-0.3 I_\nu \). The effects of dielectronic recombination are apparent in Fig. 10.1. For example, the He-like Fe XXV and Ne-like Fe XVII are very durable and span large temperature ranges because they have low dielectronic recombination rates (\( \alpha_{d}/\alpha_e < 1 \)) and high ionization potentials, whereas ions with partially filled L or M shells span relatively narrow temperature ranges because their dielectronic recombination rates are high (\( \alpha_{d}/\alpha_e > 10 \)).

Fig. 10.1. Coronal model fractional abundances, \( f_z(T) \), of ionization stages of iron. (Courtesy of Dr. J. M. Shull).
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Given the ionization fractions, $f_i(T)$, one can calculate the emission spectrum by summing the spectral emissivities due to the various radiative processes. The line spectrum is dominated by radiative decays following electron impact excitation; the volume emissivity of a given line is given by:

$$\eta_i(T) = n_e n_H A X f_i(T) q_e(T) \text{ erg cm}^{-3} \text{ s}^{-1},$$

(10.8)

where $A_X$ is the abundance of element X relative to hydrogen (atomic number density $n_H$) and $q_e(T)$ is given by (10.1). The emission line spectrum of the gas includes the sum of all these collisionally excited lines, plus a relatively small contribution of recombination lines. The continuum emissivity is dominated by bremsstrahlung and recombination, and has a spectral shape given approximately by $d\eta_c/d\varepsilon \propto \exp(-\varepsilon/kT)$.

The total power radiated by a unit volume of coronal gas is calculated by integrating the line plus continuum emissivity over the spectrum; it may be written $P = n_e n_H \Lambda(T)$. Fig. 10.2 shows the radiative cooling function $\Lambda(T)$ and indicates the contributions of different elements to the emission line cooling in particular temperature ranges. For $T > 10^5 \text{ K}$ it may be

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Fig. 10.2. Radiative cooling function $\Lambda(T)$, for coronal model. Contributions to cooling from individual elements are indicated. (From Gaetz & Salpeter (1983); courtesy of Astrophysical Journal.)
approximated by
\[ \Lambda(T) = 1.0 \times 10^{-22} T_6^{-0.7} \times 2.3 \times 10^{-24} T_6^{0.5} \, \text{erg cm}^3 \text{s}^{-1}, \] (10.9)
where the former term accounts for line emission and is accurate within a factor 2 and the latter term accounts for bremsstrahlung. For \( T > 5 \times 10^6 \, \text{K} \) the gas radiates most of its power in the X-ray (\( \varepsilon > 0.5 \, \text{keV} \)) band. The emission line spectrum is dominated by Lyman-\( \alpha \) transitions of hydrogenic ions and the \( 1s^2 \rightarrow 1s2p \) transitions of He-like ions of abundant elements such as O, Si, S, and Fe. The spectrum also has a rich cluster of emission lines near 1 keV due to L-shell transitions of Fe ions. When the temperature drops below \( \approx 5 \times 10^6 \, \text{K} \), the spectrum changes radically. The emission line spectrum is dominated by many strong EUV (\( 13.6 < \varepsilon < 100 \, \text{eV} \)) lines from ions of O, Si, S, with L-shell electrons and from Fe ions with M-shell electrons. (The strongest emission lines with \( \varepsilon < 13.6 \, \text{eV} \) are Lyman-\( \alpha \), C IV \( \lambda 1550 \) and O VI \( \lambda 1035 \)). In the temperature range \( 10^4 < T < 3 \times 10^6 \, \text{K} \) the power in this EUV line emission exceeds the continuum power by factors \( \approx 10^2 \) and is much greater than the power radiated at higher or lower temperatures.

The peak in \( \Lambda(T) \) forms a ‘thermal barrier’ that causes interstellar gas to avoid this range of temperatures. One way to understand this effect is to consider the thermal relaxation of a gas that has been heated suddenly to some initial temperature \( T_0 \). Then, if the gas cools at constant pressure, its temperature will relax according to the equation
\[ \frac{d}{dt} \left( \frac{5}{2} kT \right) = -n_c \Lambda(T). \] (10.10)
It follows from (10.10) that the thermal relaxation time scale is roughly \( t_r \approx 5kT/[2n_c \Lambda(T)] \). For typical interstellar pressures, \( nT = 10^4 \, \text{cm}^{-3} \, \text{K} \), we may estimate from (10.10) that \( t_r \approx 10^7 T_6^{0.7} \) yr. We see from this example that the gas temperature will linger near \( T \) for a time \( \approx t_r \) and then will decrease at an accelerating rate to \( T \lesssim 10^4 \, \text{K} \), at which point the radiative cooling drops rapidly and the temperature may be stabilized by the heating of ultraviolet starlight. We also see that radiative cooling is unimportant for interstellar gas with \( T \gtrsim 10^7 \, \text{K} \). This result, which depends only on atomic physics and is independent of the details of the heating mechanism, is the main reason why interstellar gas is usually found with \( T \lesssim 10^4 \, \text{K} \) or with \( T \gtrsim 10^6 \, \text{K} \), but rarely with intermediate temperatures.

The above example is not self-consistent, because we used the radiative cooling function \( \Lambda(T) \) that was derived with the assumption of a stationary ionization balance in order to estimate timescales for relaxation of a
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nonstationary gas. In fact, fairly substantial errors arise from this approximation, because changes in ionization will lag changes in temperature. Consequently, gas which has been cooling down will be 'over-ionized' compared with a stationary gas at a given temperature, while gas which has recently been heated impulsively will be 'under-ionized'. Since lower ionization stages of a given element usually have more low-lying bound states that may be excited by thermal electrons, the radiative cooling at a given temperature is enhanced compared with a stationary gas if the gas is under-ionized, and suppressed if it is over-ionized. The enhancement or suppression factors may be substantial, say, \( \approx 3 \). As will be discussed later, these nonstationary effects can be important for the interpretation of the spectra of supernova remnants. However, they do not obviate the qualitative conclusions of the previous paragraph.

10.3 Photoionized gas: nebular approximation

The ionization state and radiative cooling of interstellar gas is very different if the gas is illuminated by X-rays from a compact galactic X-ray source or an active galactic nucleus. In that case photoionization dominates ionization by thermal electrons for most ions. The X-rays may also dominate the heating of the gas. As a result, the ionization and temperature of the gas may reach a stationary state that is determined by the gas pressure and the flux and spectrum of radiation. Models for the ionization and spectral emissivity of gas around a compact X-ray source are called 'X-ray nebular models', by analogy with models for planetary nebulae.

Suppose that a parcel of gas with given pressure \( P = 2.3 n_{\text{H}} kT \) is located a distance \( r \) from a compact isotropic X-ray source with luminosity \( L \) and spectrum \( dL/d\varepsilon = L_\varepsilon \). Assume further that the primary spectrum is not significantly attenuated by photoabsorption. Then the photoionization rate for electrons in shell \( s \) of an ion, \( X^+ \), is given by:

\[
J_{\text{s}} = \frac{L}{4\pi r^2} \int_{I_{\text{s}}}^{\infty} \frac{d\varepsilon}{\varepsilon} g(\varepsilon) \sigma(\varepsilon) \quad \text{s}^{-1},
\]  

(10.11)

where \( I_{\text{s}} \) is the photoionization threshold. The photoionization cross section, \( \sigma(\varepsilon) \), is a rapidly decreasing function of \( \varepsilon \); for highly stripped ions or inner shells it may be approximated fairly well by

\[
\sigma(\varepsilon) = 7.9 \times 10^{-18} (13.6 \text{ eV}/I_{\text{s}}) r_9 F_s (I_{\text{s}}/\varepsilon)^3 \quad \text{cm}^2,
\]  

(10.12)

where \( r_9 \) is the number of electrons in shell \( s \) and the numerical factor \( F_s \approx 1 \) for hydrogenic ions, \( F_s \approx 0.8 \) for filled K shells, and \( F_s \approx 0.5 \) for L-shell electrons.
Photoionization of inner shell electrons is often followed by the emission of one or more Auger electrons. For example, if an ion has two or more L-shell electrons, photoionization of a K-shell electron will leave the resulting ion in an autoionizing state that can relax by ejecting one or more L-shell electrons as another L-shell electron drops to the K-shell. Likewise, the L-shell vacancy may result in further Auger electrons if the ion has two or more M-shell electrons. Thus, there is a set of probabilities, \( p_{z,z'}(s) \), that photoionization from shell \( s \) of ion \( X^z \) will result in ion \( X^{z'} \). Given these probabilities, we may write an expression for the rate to produce ions \( X^{z'} \) by photoionization of ion \( X^z \):

\[
J_{z,z'} = \sum_s J_{z,s} p_{z,z'}(s) = \frac{L}{4\pi r^2} j_{z,z'},
\]

where we have removed the explicit dependence on \( L/(4\pi r^2) \) in defining the quantities \( j_{z,z'} \), which are numbers depending on the spectral shape function \( g(\xi) \).

As in the coronal case, we may write the master equation (10.5) describing the rates of change of the ion fractions, \( f_z \), and look for stationary solutions such that \( df_z/dt = 0 \). However, the matrix \( M_{z,z'} \) is no longer tri-diagonal because the Auger emission populates ionization stages \( z \) from stages with \( z' < z - 1 \). The matrix elements including photoionization and Auger emission may be written:

\[
M_{z,z'} = j_{z,z'}, \quad z' < z - 1
\]

\[
M_{z,z-1} = q_{z-1} + \xi j_{z,z-1},
\]

\[
M_{z,z} = -(q_z + \alpha_{r,z-1} + \alpha_{d,z-1} + \xi j_{z,z'}),
\]

and

\[
M_{z,z+1} = (\alpha_{r,z} + \alpha_{d,z}),
\]

where we have defined an ‘ionization parameter’, \( \xi = L/(4\pi r^2 n_e) \). The stationary ionization fractions are found by solving the matrix equation \( \sum_{z'} M_{z,z'} f_{z'} = 0 \), with the normalization \( \sum_{z'} f_{z'} = 1 \). We see that the solutions depend on two parameters: \( f_z = f_z(T, \xi) \). Of course, if \( \xi = 0 \) we recover the coronal model, and with increasing \( \xi \) the ionization at a given temperature shifts to higher stages.

As with the coronal model, the spectral emissivity of the photoionized gas is calculated by summing the line and continuum emissivities due to all radiative transitions as in (10.8). As before, the total power radiated per unit volume can be written \( P = n_e n_H \Lambda(T, \xi) \). Generally, the radiative cooling function \( \Lambda(T, \xi) \) is a decreasing function of \( \xi \) for fixed \( T \) because as elements
become more highly ionized they tend to have fewer low-lying resonance transitions. Indeed, as \( \xi \) becomes very large the peak in \( \Lambda(T, \xi) \) for \( 10^4 < T < 10^7 \) K due to line emission vanishes, leaving only bremsstrahlung cooling.

In the nebular model it is assumed that the gas temperature obtains a stationary state in which this radiative cooling is balanced by heating due to photons from the central source. Thus, \( T \) is no longer an arbitrary parameter of the model but is determined as a function of \( \xi \). The photoelectric heating per ion is given by an expression similar to (10.11):

\[
\Gamma_{x,z} = \frac{L}{4\pi r^2} \sum_s \sum_{e,z} \int_{e_{\text{th}}}^{\infty} \frac{1}{e} g(e) \sigma_{x,s}(e) E_{e,s}(e),
\]

where \( E_{e,s}(e) \) is the net kinetic energy of photoelectrons and Auger electrons resulting from an ionization from shell \( s \). The net heating per hydrogen atom is calculated by summing over all species:

\[
\Gamma_{p}(T, \xi) = \sum_x A_x \sum_z f_z(T, \xi) \Gamma_{x,z}.
\]

In addition, the electrons are heated directly by Compton scattering at a net rate (per hydrogen atom):

\[
\Gamma_c = \frac{L}{4\pi r^2} \frac{n_e}{n} \sigma_T \int \text{d}e g(e) \left[ \frac{e}{m_e c^2} - \frac{4kT}{m_e c^2} \right]
\]

where the second term in brackets represents cooling of hot electrons by scattering off soft photons. This Compton heating term vanishes when the \( kT \approx \epsilon_{\text{max}} \), where \( \epsilon_{\text{max}} \) is a characteristic maximum X-ray photon energy. Since \( L/(4\pi r^2) \) is a common factor of all heating terms, the net heating per unit volume can be written in the form

\[
\Gamma_{\text{TOT}} = \frac{L}{4\pi r^2} \gamma(T, \xi).
\]

Now, by equating this radiative heating to the radiative cooling, we find:

\[
\xi \gamma(T, \xi) = \Lambda(T, \xi).
\]

Equation (10.19) is a complicated nonlinear equation that depends implicitly as well as explicitly on \( \xi \) and \( T \) through the ionization fractions \( f_z(T, \xi) \). However, it can be solved on a computer, giving the solution \( \xi = L/(4\pi r^2) = \xi(T) \).

It is convenient to define a new function,

\[
\Xi(T) = \frac{L}{4\pi r^2 c} \frac{\xi(T)}{2.3 c k T},
\]
which is the dimensionless ratio of radiation pressure to gas pressure. This function determines the possible stationary thermal states of gas that is heated by a cosmic X-ray source. It depends on cosmic abundances and on the shape, \( g(\varepsilon) \), of the X-ray source spectrum; a typical example, with standard cosmic abundances and \( g(\varepsilon) \propto \varepsilon^{-0.7} \), is plotted in Fig. 10.3. For points to the right of the curve, heating exceeds cooling and the temperature must rise; to the left, the temperature must fall. Thus, specifying gas pressure, \( P \), source luminosity, \( L \), and radial distance, \( r \), determines the value of \( \Xi \), and given that value we can find the corresponding value of \( T \) from Fig. 10.3. Far from the source, \( \Xi \) has a low value, say, \( \Xi_1 \), and the stationary temperature will be \( T_\Xi \approx 10^4 \) K, where the strong cooling due to emission of EUV resonance lines balances the X-ray photoelectric heating. Close to the source, \( \Xi \) has a high value, say, \( \Xi_3 \), and most elements are fully ionized so that there is very little line emission and the stationary temperature will obtain a high value, \( T_h \), for which the net Compton heating (10.17) is almost zero. In this limit the equilibrium temperature is determined solely by the spectral shape function \( g(\varepsilon) \).

A very interesting phenomenon occurs when \( \Xi = \Xi_2 \), where \( \Xi_2 \) falls in the

Fig. 10.3. Gas temperature, \( T \), as a function of ionization parameter, \( \Xi \), in the nebular model. Source spectrum is power law, \( dL/d\varepsilon \propto \varepsilon^{-0.7} \), with high energy cut-off at \( \varepsilon = 500 \) keV. (Courtesy of Dr. T. Kallman.)
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intermediate range $\Xi_{\text{min}} < \Xi < \Xi_{\text{max}}$. In this case there are three temperatures, $T_c$, $T_h$, and $T_h$, for which heating equals cooling. However, the intermediate state, $T = T_c$, is thermally unstable: a perturbation to higher (or lower) temperature at fixed $\Xi$ will grow because it puts the state of the gas to the right (or left) of the curve where heating is greater than (or less than) cooling. By the same argument, both the cool ($T = T_c$) and hot ($T = T_h$) states are stable. Thus, in the region where $\Xi_{\text{min}} < \Xi < \Xi_{\text{max}}$, relatively dense 'clouds' with $T = T_c \approx 10^4$ K can coexist in pressure equilibrium with low-density coronal gas with $T = T_h \approx 10^7 - 10^8$ K.

Although the stationary nebular model discussed here is very different from the time-dependent cooling model discussed in Section 10.2, the qualitative conclusion is the same: interstellar gas is rarely found in the intermediate temperature range $10^4 < T < 3 \times 10^6$ K where the atomic cooling function $\Lambda(T)$ peaks.

The detailed shape of the function $\Xi(T)$ is sensitive to the shape of the source spectrum shape $g(\epsilon)$. Generally, $\Xi(T)$ doubles back on itself if the source spectrum is fairly flat up to the X-ray band, i.e., $n < 2$ for $g(\epsilon) \approx \epsilon^{-n}$. The range from $\Xi_{\text{min}}$ to $\Xi_{\text{max}}$ increases with decreasing $n$ and if the source spectrum is deficient in low-energy photons. However, if the source spectrum is steep, $n > 2.5$, $\Xi(T)$ does not double back on itself and there is no range of $\Xi$ for which two phases may coexist.

In this two-phase model the mass fractions of gas in the cool and hot phases are indeterminate. However, one can imagine a scenario that would have the gas pressure at any distance from the source rising to that value for which $\Xi = \Xi_{\text{max}}$. Suppose that there is initially a distribution of clouds with $T = T_c$, but at low pressure so $\Xi > \Xi_{\text{max}}$. Then the X-rays will heat the gas at the surface of the clouds to $T_h$, causing it to expand and fill the volume between the clouds. This process will continue until the pressure of the hot intercloud medium is great enough that $\Xi = \Xi_{\text{max}}$, at which point the clouds become stable.

This multiphase behavior plays an important role in the dynamics of gas near any cosmic X-ray source. For example, the model predicts that hot coronal gas will surround the dense gas associated with an accretion disk and stream near a compact galactic binary X-ray source. Another prediction is that the dense clouds responsible for the broad emission lines in active galactic nuclei and quasars ('AGN') should be confined by the pressure of a hot coronal gas. If this model for the broad-line clouds is correct, the parameter $\Xi$ should be restricted to the fairly narrow range, $\Xi_{\text{min}} < \Xi < \Xi_{\text{max}}$. The value of $\Xi$ for the broad-line clouds can be inferred by fitting theoretical models for the emission line spectra to the observed spectra, and it indeed
seems to be in the right range. The model also provides a possible explanation for the absence of broad emission lines in active galactic nuclei, such as BL Lacertae objects, which have steep (spectral index \( n > 2 \)) X-ray spectra; for such sources the curve \( \Xi(T) \) does not double back on itself, so there is no region in which dense clouds can coexist with coronal gas.

In the description of the nebular model up to this point, we have neglected the effects of photon absorption and have assumed that the source spectrum is attenuated only by the \( 1/r^2 \) geometrical dilution. This approximation is adequate for determining whether broad-line clouds may exist, provided that the Compton optical depth of the hot gas is small. However, the approximation is not usually adequate for calculating the emission line spectrum of broad-line clouds. Since the cooler emission line clouds have much greater density at constant pressure, \( n_e T_e = n_h T_h \), they are likely to have photoelectric optical depths, \( \tau(e) \gg 1 \). To take account of the photoelectric absorption of the source spectrum, one must replace \( g(e) \) by \( g(e) \exp \left[ -\tau(e) \right] \) in (10.11) and (10.15) and one must add the diffuse emission of ionizing radiation to \( g(e) \). Then it becomes necessary to solve a radiative transfer problem, in which the system of equations (10.11)–(10.19) must be solved at every depth point in the cloud.

Several authors have constructed such radiative transfer models for the emission line spectra of AGN. Since the line emitting regions must be relatively cool, \( T = T_e \approx 10^4 \) K, and dense, \( n_e \approx 10^9 - 10^{11} \) cm\(^{-3} \), we require \( \Xi < \Xi_{\text{max}} \) at the illuminated surface of the cloud. The line emitting region is thin, \( d \approx 10^{12} \) cm, compared with its radial distance, \( \approx 10^{18} - 10^{19} \) cm from the source, so \( \Xi \) is almost constant throughout the cloud. Generally, the surface of the cloud is an H II zone in which most of the ionizing uv luminosity is converted into uv emission lines, the strongest of which are C IV \( \lambda 1549 \) and O VI \( \lambda 1035 \). Beneath the H II region is a deeper, partially ionized H I zone, in which the luminosity of the more penetrating X-rays is converted primarily to hydrogen (\( L_\alpha, H_\alpha, H_\beta \), etc.), Mg II \( \lambda 2798 \), and Fe II line emission. It is remarkable, and possibly fortuitous, that theoretical emission line spectra agree fairly well with the observed spectra even for models consisting of a single slab at a given distance from the source.

### 10.4 Thermal conduction

In Section 10.3 we described how relatively cool (\( T \approx 10^4 \) K) gas might coexist in pressure equilibrium with hot (\( T \approx 10^7 - 10^8 \) K) gas, and we pointed out that gas at intermediate temperatures would be thermally unstable and therefore improbable. However, this statement cannot be strictly true at the interfaces between the cool and hot gas, where the
temperature must change continuously. Otherwise the heat flux due to conduction, \( F = K(T) \nabla T \), would be infinite. As we shall see, thermal conductivity at these interfaces produces observable amounts of intermediate temperature gas and can play an important role in the thermodynamics of the interstellar medium.

According to the classical kinetic theory, the thermal conductivity of a fully ionized plasma is given by \( K(T) = CT^{5/2} \), where the constant \( C \approx 6 \times 10^{-7} \text{ erg cm}^{-1} \text{ s}^{-1} \text{ K}^{-7/2} \). We may write the second law of thermodynamics for a compressible gas, including thermal conduction, as follows:

\[
\frac{D}{Dt} \left( \frac{3 P}{2 \rho} \right) - \frac{P}{\rho^2} \frac{D}{Dt} \frac{\rho}{\rho} = \nabla \cdot [K(T) \nabla T] - \frac{\eta \mu}{\rho} \Lambda(T) \quad (10.21)
\]

where the gas pressure \( P = \rho k T / \mu, \mu = 0.6 m_{H} \), and \( D/Dt \) represents the time derivative in Lagrangian coordinates. The terms on the right-hand side represent heat loss due to conduction and radiation, respectively. The dynamics and thermodynamics of interstellar gas may be described by (10.21) and Euler's equation, a very complicated set of equations in general.

Here we shall describe only a highly simplified model problem, which nevertheless illustrates the main effects of thermal conduction. The model is as follows. Assume that there is a spherical cloud of cold \( (T = T_{c} \approx 0) \) gas with radius \( R \) embedded in hot gas with temperature \( T_{h} \) at very large distance. Assume further that there is a stationary flow of mass from the cloud, \( \dot{M} = 4\pi r^{2} \rho v \). (For truly stationary flow, there must be a hypothetical mass source within the cloud.) Assume that the gas flow is everywhere subsonic, in which case Euler's equation may be replaced by the condition \( P = \text{const} \). Finally, assume that the radiative cooling term may be neglected. (The validity of these assumptions will be discussed below.) Then, replacing \( D/Dt \) by \( v d/dr \), we obtain:

\[
\frac{5}{2} \frac{\dot{M}}{4\pi \mu} \frac{k dT}{dr} = \frac{d}{dr} \left[ \frac{r^{2} C T^{5/2} dT}{dr} \right]. \quad (10.22)
\]

The solution of (10.22), subject to the boundary condition \( T(R) = 0 \), is

\[
T(r) = T_{h} (1 - R/r)^{2/5}. \quad (20.23)
\]

The existence of such a stationary solution requires that

\[
\dot{M} = \frac{16\pi \mu C}{25} RT_{h}^{5/2} \approx 4 \times 10^{-7} M_{\odot} \text{ yr}^{-1} R_{pc} T_{h}^{5/2}, \quad (10.24)
\]

where \( T_{h} = T_{h}/(10^{6} \text{ K}) \). Thus, we see that the heat conducted from hot gas into cool clouds causes them to 'evaporate', with mass loss rate proportional
to the cloud radius. From (10.24) we may estimate the time in which a cloud of atomic density \( n_e \) and radius \( R_{pc} = R/(1 \text{ pc}) \) will be completely evaporated:

\[
t \approx 3 \times 10^5 \text{ yr} \ n_e R_{pc}^2 T_6^{-5/2}.
\]  (10.25)

Equation (10.23) tells us the distribution of intermediate temperature \( (10^5 < T < 10^7 \text{ K}) \) gas in the interfaces between hot interstellar gas and cool clouds. Such gas can be detected through narrow ultraviolet absorption lines of ions such as C IV, N V, and O VI that predominate at these temperatures (cf. Section 10.2). However, such ions are not unambiguous tracers of intermediate temperature gas because they can also be produced in cooler gas by photoionization. The most reliable tracer of such gas is probably the O VI \( \lambda 1035 \) line, because (except near a cosmic X-ray source) there is probably very little interstellar radiation at photon energies \( \varepsilon > 114 \text{ eV} \), the ionization potential of O V. Assuming that the O VI fractional abundance, \( f_{OVI}(T) \), is given by the stationary coronal approximation, its column density in the conduction front may be derived from (10.23) and the condition of pressure equilibrium, \( n(r)T(r) = n_h T_h \). The result is

\[
N_{OVI} = \frac{5}{2} X_O n_h T_h^{-3/2} R \int_0^{T_h} \frac{f_{OVI}(T) T^{1/2} dT}{[1-(T/T_h)^{5/2}]}.
\]  (10.26)

where we have written \( n_{OVI} = n(r) X_O f_{OVI}(T) \). We can evaluate the integral in (10.26) \( [f_{OVI}(T) \text{ peaks at } T \approx 3 \times 10^5 \text{ K}] \) to obtain \( N_{OVI} \approx 10^{14} n_h T_6^{-3/2} R_{pc} \) \( \text{cm}^{-2} \), where we have assumed a cosmic oxygen abundance \( X_O = 6.8 \times 10^{-4} \). Observed column densities of interstellar O VI are typically in the range \( N_{OVI} \approx 10^{13} \text{ to } 10^{14} \text{ cm}^{-2} \). These values can be understood if the line of sight passes through several clouds with, say, \( R \approx 1 \text{ pc}, \ T_6 \approx 1, \) and \( n_h \approx 10^{-2} \text{ cm}^{-3} \).

This order-of-magnitude agreement is appealing and suggests that the observed O VI column density does come mainly from conduction fronts. However, this interpretation must be regarded as tentative for a variety of reasons. First, the predicted column densities are sensitive to uncertain parameters such as \( n_h, T_6, R_{pc}, \) and number of interfaces along the line of sight. Second, the observed O VI column density may contain a significant but very uncertain contribution due to photoionization by soft X-rays. Third, there is a great deal of uncertainty in the theory itself.

One may rightly question the assumptions that led to (10.22)–(10.26). First, consider the assumption that the flow is everywhere subsonic, \( v/c_s = v/(kT/\mu)^{1/2} < 1 \). One finds for the stationary flow above that the ratio \( v/c_s \) reaches a maximum for \( r = 1.1 R \), and that the condition \( v/c_s < 1 \) at that point can be written approximately \( R_{pc} \approx 10^{-2} T_6 n_h^{-1} \), where \( n_h \) is the density of
the hot intercloud gas at large distance. Thus, for typical parameters of the hot gas, \( T_6 \approx 1 \) and \( n_h \approx 10^{-2} \text{ cm}^{-3} \), this condition is violated for clouds with radius \( R < 1 \text{ pc} \). This is the regime of ‘saturated conductivity’, in which the heat flux is limited to a few times \( \rho_h c_h^3 \), where \( c_h = (kT_h/\mu)^{1/2} \). Mass evaporates from the cloud surface supersonically with a flux of a few times \( \rho_h c_h \), and the rocket force due to this supersonic expansion raises the cloud pressure substantially above that of the hot intercloud gas.

The other assumption that one must check is that of neglecting the radiative cooling term in (10.21). This may be done by using the solution (10.23), in which radiative cooling is neglected, to evaluate the ratio of the radiative cooling term to the advected enthalpy term in (10.21). Using (10.9), we find that this ratio is small for \( r \approx R \) only if \( R_{\text{pc}} < 0.1 \ T_6^{2/3} n_h^{-1} \). If this condition is violated, radiative losses compete effectively for the heat that is conducted toward the clouds and suppress the evaporative mass flux. Indeed, the mass flux may even reverse, causing gas to condense onto large clouds while it is evaporating off the small clouds. Note that the domain for which (10.22)–(10.26) are valid is only a factor \( \approx 10 \) in cloud radius.

When one attempts to include the effects of radiative losses and saturation, the theory of conduction fronts becomes much more complicated and is not yet fully understood. A further major uncertainty arises from the possible effects of magnetic fields. Since magnetic fields can greatly suppress thermal conductivity across field lines, the structure of conduction fronts may depend strongly on the topology of interstellar magnetic fields at the interfaces between the hot and cool gas.

The conclusions that we may draw from this section are rather unsatisfying. The simple classical theory indicates that thermal conductivity may play a major role in the thermodynamics of coronal interstellar gas, cooling the hot medium and evaporating the smaller clouds. Furthermore, the observed column densities of O VI are roughly consistent with the theory. However, we know that the classical theory is valid only over a very limited range of parameters, and that the theory itself becomes very complicated and uncertain outside this range or if magnetic fields are present.

10.5 Supernova remnants, interstellar bubbles, and superbubbles

10.5.1 Supernova remnants

In a supernova explosion, a star ejects a mass \( M_\odot \approx 1 - 10 \ M_\odot \) into interstellar space with terminal velocity \( V_\odot \approx 10^4 \text{ km s}^{-1} \) and kinetic energy \( E_\odot \approx 10^{51} \text{ ergs} \). The ejecta expand at nearly constant velocity until they encounter a comparable mass of ambient interstellar gas (density \( \rho_0 = n_0 m_0 \),
where $m = 1.4 m_\odot$). This occurs at a time $t_0 \approx 200 \ yr \ (M_\odot/M_\odot)^{1/3} n_0^{-1/3}$ and radius $R_0 \approx 2 \ pc \ (M_\odot/M_\odot)^{1/3} n_0^{-1/3}$. From the surface of impact a shock propagates outward into the interstellar gas and another shock propagates inward into the expanding ejecta. After a few times $t_0$ the ejecta have slowed down and have transferred most of their kinetic energy into the energy of the outgoing shock. This blast wave continues to expand for $> 10^4 \ yr$, sweeping the interstellar gas into a thin shell and leaving a hot rarefied cavity of coronal gas that may linger for millions of years before it is once again filled with 'normal' ($T < 10^4 \ K$) interstellar gas.

Suppose, as a first approximation, that the ambient interstellar gas (temperature $T_0$) has uniform density $\rho_0$. In that case a detailed theory has been developed for the hydrodynamics of the supernova blast wave. Here we present a simplified approximate version of that theory, which illustrates the main physical effects and is surprisingly accurate. We assume that the mass of the system resides in a thin expanding shell of radius $R(t)$, and that it is dominated by the mass of swept-up interstellar gas, $M(t) = \frac{4}{3} \pi R(t)^3 \rho_0$. We assume that the shell is expanding hypersonically, $\dot{R} \gg (kT/\mu)^{1/2}$. Then, according to the jump conditions for a monatomic gas, the shocked gas in the shell will have density $\rho_1 = 4 \rho_0$ and temperature $T_1 = \frac{3}{4} \mu k R^2/k$. The thickness of the shell, $\Delta R \approx R/12$, may be estimated from the condition $4\pi R^2 \Delta R \rho_1 = \frac{4}{3} \pi R^3 \rho_0$. The dynamics of the shell follows from Newton's Second Law applied to a conical section of solid angle $\Delta \Omega$: $(d/dt)(\Delta M \dot{R}) = \Delta F$, where $\Delta M = \frac{1}{3} R^3 \rho_0 \Delta \Omega$ and the driving force $F = 3 R^2 \Delta \Omega$. Thus, we have the general equation for an expanding spherical system:

$$\frac{d}{dt} \left[ \frac{1}{3} R^3 \rho_0 \dot{R} \right] = R^2 \dot{P}.$$ \hspace{1cm} (10.27)

To solve (10.27) we need an expression for the interior pressure, $P$. In the early part of the expansion, the total energy of the blast wave is conserved. Assume that this energy is mostly thermal energy and that the interior pressure is uniform: then $P = \frac{2}{3} E_0/(\frac{4}{3} \pi R^3)$. Inserting this result into (10.27), we find the solution,

$$R(t) = \left[ \frac{25 E_0 t^2}{4 \pi \rho_0} \right]^{1/5} \approx 13 \ pc \ \left( \frac{E_{51}}{n_0} \right)^{1/5} t_4^{2/5};$$ \hspace{1cm} (10.28)

where $E_{51} = E_0/(10^{51} \ ergs)$ and $t_4 = t/(10^4 \ yr)$. Equation (10.28) agrees with the exact result within 1 percent. From it we may also derive the expansion velocity, $\dot{R}(t)$, and the temperature of the shocked gas,

$$T = 3.3 \times 10^6 \ K \left( \frac{E_{51}}{n_0} \right)^{2/5} t_4^{-6/5}.$$ \hspace{1cm} (10.29)
Note that the shell remains hot enough \((T > 10^6 \text{ K})\) to emit observable soft X-rays for \(\approx 3 \times 10^4 \text{ yr}\). The X-ray spectrum is dominated by continuum for \(\varepsilon > 1.5 \text{ keV}\) and by strong emission lines of C, N, O, Ne, Mg, and Fe for \(\varepsilon < 1.5 \text{ keV}\).

The instantaneous power radiated by the shell is given by

\[
L(t) = 4\pi R^2 \Delta R n_e^2 \Lambda(T) = \frac{16\pi}{3} R^3 n_e^2 \Lambda(T).
\]  

(10.30)

If we use (10.28), (10.29), and the approximate formula (10.9) for \(\Lambda(T)\), we find that \(L(t)\) rises slowly at first, \(L(t) \approx t^{0.6}\) for \(T > 10^7 \text{ K}\), then rapidly, \(L(t) \approx t^{2.04}\) for \(T < 10^7 \text{ K}\), when strong line emission sets in. Actually, \(L(t)\) rises even more rapidly than this. Because the shocked gas is 'under-ionized' (cf. Section 10.2), the cooling due to line emission is enhanced by factors \(\approx 3\). The X-ray emission spectra of supernova shells clearly show enhanced line emission due to these non-equilibrium effects.

The energy conserving phase ends at a time \(t_1\) when \(\int L(t) \, dt \approx 0.3 \, E_0\). From (10.28)–(10.30) and (10.9) we obtain

\[
t_1 \approx 3 \times 10^4 \text{ yr} \, E_{51}^{0.22} n_0^{-0.55},
\]

and

\[
R_1 = 20 \text{ pc} \, E_{51}^{0.39} n_0^{-0.42},
\]

(10.32)

where we have increased \(\Lambda(T)\) by a factor of 3 to account for the enhanced line emission. After this time the shell collapses, becoming cool, dense, and very thin, and ceases to emit observable X-rays. However, the shell continues to expand, pushed by very hot low-density interior gas, which does not radiate significantly. The pressure of the interior gas decreases according to the adiabatic law, \(PV^{5/3} = \text{const.}\), or \(P = P_1 (R/R_1)^{-5}\), where \(P_1 = \frac{3}{2} E_0/\left(4\pi R_1^2\right)\). Inserting this expression into (10.27), we obtain the expansion law for late times:

\[
R(t) = \left[\frac{147 E_0 R_1^2 t_4^2}{4\pi \rho_0}\right]^{1/7} \approx 19 \text{ pc} \, E_{51}^{0.23} n_0^{-0.26} t_4^{2/7},
\]

(10.33)

where we have used (10.32) to derive the second expression. The supernova shell will continue to expand for \(t > 10^6 \text{ yr}\), reaching a radius \(> 60 \text{ pc}\). Finally, after several million years, relatively cool and dense gas will flow back into the cavity.

The simplified theory presented above agrees very well with more-detailed theoretical models. However, the theory might have limited applicability to actual supernova explosions because a key assumption of the model, that of uniform density ambient gas, may be invalid. In fact, the interstellar medium
is known to be intrinsically nonuniform, containing a chaotic distribution of relatively dense gas clouds and filaments embedded in a low-density substrate, much of which may be coronal ($T \approx 10^6$ K) gas. In such a complex region the development of a supernova blast wave will be very different from that described above. The blast wave will propagate rapidly through the low-density gas, passing and enclosing the dense clouds. The mass of the expanding system may be dominated by gas evaporated from the clouds rather than by the hot gas that is swept up. If so, one can show that the radius of the system increases as $R \propto t^{2/5}$ rather than $R \propto t^{1/5}$ as in (10.28).

Moreover, the ambient interstellar gas may be modified substantially by the pre-supernova star. Most supernovae in the galactic disk are thought to come from fairly massive stars, say, $> 6 M_\odot$, corresponding to main-sequence spectral type B4. If the pre-supernova star (or any associated star) is somewhat more massive, say, $> 20 M_\odot$, it would probably be luminous enough to ionize any interstellar clouds within a radius $> 30$ pc. If so, the clouds would expand and fill the region with gas of fairly uniform density. On the other hand, massive stars are also likely to have strong stellar winds, and such winds will create cavities of low-density gas around the pre-supernova star. Furthermore, since massive stars tend to be found in clusters and associations, there is a good chance that a given supernova will occur in a region of interstellar gas that has been disturbed by other supernovae during the past few million years. These possibilities are discussed below.

### 10.5.2 Interstellar bubbles

Early-type stars with bolometric luminosity $L_* > 10^5 L_\odot$ are known to lose mass in strong stellar winds, with mass-loss rates $10^{-8} < \dot{M} < 10^{-5} M_\odot$ yr$^{-1}$ and terminal velocities $1000 < V_w < 4000$ km s$^{-1}$. The mechanical luminosity of the wind, $L_w = M V_w^2 / 2$, is a strongly increasing function of bolometric luminosity, $L_w \approx 10^{35}$ erg s$^{-1} \left[ L_* / (10^5 L_\odot) \right]^{1.7}$. Over its lifetime, such a star can eject a substantial fraction of its mass with such velocities and a net mechanical energy $\approx 10^{50}$ ergs. The result is a huge cavity in the ambient interstellar medium that resembles a supernova remnant. Here we sketch the theory of these expanding ‘interstellar bubbles’.

As with supernova remnants, interstellar bubbles at first conserve most of their energy in a hot expanding shell, reach a radius at which radiative losses cause the shell to collapse, and thereafter expand more slowly. However, the energy-conserving phase of the bubble lasts longer, typically a few times $10^6$ yr instead of a few times $10^4$ yr. Thus, typical interstellar bubbles are likely to be somewhat larger than supernova remnants, even though they are formed with less total energy.
The expansion law for interstellar bubbles can be derived fairly accurately from the thin shell approximation. For the energy conserving phase one takes \( P = \frac{2}{3} L_w t / (\frac{1}{3} \pi R^3) \) for the right-hand side of (10.27) and obtains

\[
R(t) = \left[ \frac{25}{14 \pi} \frac{L_w t^2}{\rho_0} \right]^{1/5} \approx 31 \text{ pc} \left( \frac{L_{36}}{n_0} \right)^{1/5} t_6^{3/5},
\]

(10.34)

where \( L_{36} = L_w / (10^{36} \text{ erg s}^{-1}) \) and \( t_6 = t / (10^6 \text{ yr}) \).

The interior structure of an interstellar bubble differs significantly from that of a supernova blast wave. The thin shell of swept-up interstellar gas is driven by the pressure of the hot shocked stellar wind, which occupies most of the interior volume. (The wind itself is stopped in a shock at a radius \( r_1 < R \).) There is a conduction front at the interface of the shell and the hot interior gas, where some of the swept-up interstellar gas in the shell is evaporated and mixed with the hot shocked wind.

Very early in the evolution of the bubble, \( t \approx 10^4 \text{ yr} \) \( n_0^{0.14} L_{36}^{0.3} \), the outer shell collapses due to radiative losses, becoming dense and thin. However, the bubble continues to expand according to (10.34), because the hot interior retains a fixed fraction, \( \frac{4}{11} \), of the total wind energy and continues to drive the shell. (The numerical coefficient in (10.34) drops by a factor \( \left( \frac{t_6}{t_7} \right)^{1/2} \) to 27 pc.) Eventually, the radiative losses of the hot interior become comparable with \( L_w \), and the shocked wind region loses pressure and collapses. The radius, \( r_1 \), where the wind shocks move out so that \( r_1 \approx R(t) \). This occurs at a time

\[
t_2 \approx 9 \times 10^6 \text{ yr} \ L_{36}^{0.3} n_0^{-0.7}
\]

(10.35)

and radius

\[
R_2 \approx 115 \text{ pc} \ L_{36}^{0.4} n_0^{-0.6}.
\]

(10.36)

(In order to derive (10.35), one must solve for the evolution of the density and temperature of the interior using the evaporation theory of (10.24), then calculate the radiative loss rate.)

Subsequently, the bubble continues to expand, driven now by the ram pressure of the wind itself, \( P = \rho_w V_w^2 = L_w / (2 \pi R^2 V_w) \). Putting this into the right-hand side of (10.27), we obtain the law for the late expansion of the bubble:

\[
R(t) \left[ \frac{\frac{3}{2} L_w t^2}{\pi \rho_0 V_w} \right]^{1/4} \approx 15 \text{ pc} \left( \frac{L_{36}}{n_0 V_{1000}} \right)^{1/4} t_6^{3/2},
\]

(10.37)

where \( V_{1000} = V_w / (1000 \text{ km s}^{-1}) \).

Interstellar bubbles are difficult to observe in emission because of their large size and low surface brightness. However, a few large faint ring-shaped
emission nebulae have been seen around early-type stars. Another potentially observable characteristic of the system is the narrow blue-shifted O VI $\lambda 1035$ absorption line due to the conduction front inside the bubble. The predicted column density, $N_{\text{OVI}} \approx 5 \times 10^{13}$ cm$^{-2}$, is in the range of observed column densities of O VI toward early-type stars. However, it may be difficult to distinguish the O VI in the interstellar bubbles from that in other conduction fronts that may lie along the line of sight to the star.

10.5.3 Superbubbles

There is now a substantial body of evidence for expanding shells of interstellar gas with radii exceeding 100 pc and sometimes even 1 kpc. These 'superbubbles' have been observed in various ways. Expanding HI shells have been seen in 21-cm emission line surveys of the Milky Way. Optical emission from giant shells of H II have been seen in interference filter photographs of the Milky Way and other nearby galaxies, particularly the Magellanic Clouds. Large-scale coherent structures of high-velocity gas have been observed through surveys of interstellar absorption lines in the ultraviolet spectra of hot stars. There is also one example, the 'Cygnus superbubble', of a region of soft X-ray emission in the Milky Way with radius $\approx 400$ pc. For some of these systems the kinetic or thermal energy can be inferred from observations; often the net energy is tens or hundreds of times $10^{51}$ erg, the typical energy of a supernova remnant.

A natural explanation for these superbubbles is that they are formed by the combined action of tens or hundreds of supernovae from a cluster of massive stars. Theory and observations indicate that stars with initial masses $>6 M_\odot$ will terminate their evolution as type II supernovae. Type II supernovae are associated with the massive stars and interstellar gas of the galactic disk (Population I); they account for roughly half the supernovae that occur in our Galaxy. (The other half, type I supernovae, are associated with the galactic bulge and halo (Population II) and do not significantly affect the dynamics of the disk gas.) Since massive stars are normally born in clusters and have relatively short lifetimes, type II supernovae should occur in clusters of tens or hundreds, spaced out over $\approx 5 \times 10^7$ yr, the lifetime of a $6 M_\odot$ star. Since the interval between supernova explosions is less than the time for interstellar gas to fill the cavity created by previous explosions, the repeated explosions will hammer out a huge cavity in interstellar space.

The theory for the evolution of such superbubbles is very similar to that for a stellar wind bubble. For a first approximation, one may merely replace the quantity $L_v$ in (10.34) by the quantity $r_{\text{SN}} E_0$, where $r_{\text{SN}}$ is the rate of supernova explosions in the cluster and $E_0$ is the energy of each supernova. A
simple idealized model for an OB cluster suggests that \( r_{\text{SN}} \) should be almost constant for \( t < 5 \times 10^7 \) yr. If the star cluster is formed with a typical initial mass function, \( dN_*/d(\log M_*) \approx M_*^{-\gamma} \), and the lifetime of the star can be written \( t_* \approx M_*^{-\alpha} \), we find that \( r_{\text{SN}} \approx t_*^{1/\gamma} \), where \( \gamma = \alpha/\beta - 1 \). For stars with \( 6 M_* < M_* < 20 M_* \), we have \( \alpha \approx 1.6 \) and \( \beta \approx 1.6 \), so \( \gamma \approx 0 \). Therefore, we may write \( r_{\text{SN}} \approx N_*/(5 \times 10^7 \text{ yr}) \), or \( r_{\text{SN}} E_0 \approx 0.7 \times 10^{35} \text{ erg s}^{-1} N_*/E_{51} \), where \( N_*/(5 \times 10^7 \text{ yr}) \) is the total number of stars with mass \( > 6 M_* \) initially in the cluster. Substituting this expression in (10.34), we obtain the expansion law for the superbubble in the energy conserving phase:

\[
R(t) \approx 116 \text{ pc} \left[ \frac{N_*/E_{51}}{n_0} \right]^{1/5} t^{3/5},
\]

where \( t_\gamma = t/(10^7 \text{ yr}) \).

If the assumptions of this model held true indefinitely, a cluster would continue to grow according to (10.38) until \( t_\gamma \approx 5 \), when the last supernova occurs. Thus, for a cluster such as Cyg OB2, with \( N_*/200 \), \( E_{51} = 1 \), \( n_0 = 1 \text{ cm}^{-3} \), (10.38) predicts a final \( t_\gamma = 5 \) radius \( R > 800 \text{ pc} \). (Cyg OB2 is a very young \( t < 3 \times 10^6 \text{ yr} \) cluster containing several very luminous stars with strong stellar winds. At this early time in its evolution, the Cygnus superbubble may be driven primarily by the combined action of the stellar winds rather than by supernovae.) However, there are several reasons why the assumptions of (10.38) should break down before the bubble grows to such a large radius.

The first is that radiative cooling may set in at some earlier time. As a first estimate, we may use \( L_{36} = 0.7 N_*/E_{51} \), \( N_*/200 \), \( E_{51} = 1 \), \( n_0 = 1 \text{ cm}^{-3} \), (10.35) and (10.36) to obtain \( t_2 = 5 \times 10^7 \text{ yr} \) and \( R_2 = 800 \text{ pc} \) for the time and radius at which cooling becomes important. However, these values may be overestimates if, as seems likely, the superbubble overtakes and entrains interstellar clouds which subsequently evaporate and enhance the interior cooling. But even if cooling becomes important very early, the superbubble will grow to a minimum radius determined by the momentum imparted by the supernova ejecta, given by (10.37). Taking \( V_e \approx 10^4 \text{ km s}^{-1} \), the velocity of the ejecta, we obtain for this limit \( R \approx 90 \text{ pc} t_\gamma^{1/2} \). Thus, although there is considerable uncertainty in the time and radius at which cooling becomes important, it seems clear that a star cluster with a few hundred OB stars will create an expanding superbubble that will grow to a radius \( R > 90 \text{ pc} \) in less than \( 10^7 \text{ yr} \).

The massive hot stars in a young cluster can probably ionize and heat interstellar gas to large (\( > 100 \text{ pc} \)) distances. This ionization and heating will cause filamentary gas to expand, homogenizing the interstellar density. The
ionization will also cause the expanding shell to fluoresce as a filamentary H II region. However, after \( t \approx 10^7 \) yr, all the stars with large ionizing flux (spectral type earlier than B0) will have evolved off the main sequence. Subsequent supernovae will continue to hammer the expanding shell, causing it to grow until \( t \approx 5 \times 10^7 \) yr. If there are no other ionizing stars in the region, the shell will be invisible optically, although it may show up as an H I shell in 21-cm emission line surveys.

The most important fact that we have neglected in our discussion up to now is that the interstellar gas density decreases with vertical distance, \( z \), above the galactic plane with a scale height \( \Delta z \approx 100 \) pc. When the superbubble grows to this dimension, the shell will break up and the interior energy will squirt into the galactic corona. Thus, the maximum radius of the superbubble in the Milky Way is limited to the scale height of the gas in the galactic disk (which increases with galactocentric distance).

This is probably the reason why the most spectacular superbubbles have been found in the Magellanic Clouds, where several giant H II shells have been observed with radii ranging from 300–600 pc. In contrast to the Milky Way, the Magellanic Clouds have low gravity and large gas scale height, so that a superbubble can contain its energy and grow to large dimension before it squirts into intergalactic space.

Additional evidence for superbubbles in the Magellanic Clouds comes from observations of their supernova remnants. In contrast to the Milky Way supernova remnants, they appear to expand freely, with \( R \propto t \) rather than \( R \propto t^{2/5} \), until they reach large radii, \( R \approx 50 \) pc. This phenomenon could be explained if most of the Magellanic Cloud supernovae occur inside superbubbles which have been almost evacuated by the action of previous supernovae.

Finally, we might speculate on the possibility that star formation may be induced in the expanding shell of a superbubble. If so, we would expect star formation to be most efficient in a system where the gas scale height is large. Indeed, the Magellanic Clouds are unusually rich in young massive stars. Perhaps they provide a modest example of the ‘starburst’ phenomenon that is seen on a more spectacular scale in some galaxies.

References

Section 10.1

Section 10.2

The formulae for atomic cross sections and rate coefficients given in this section are approximate. Extensive references to more accurate results can be found in the references below.


Section 10.3


Section 10.4


Section 10.5


