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GALACTIC WINDS*

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ABSTRACT

Gas ejected from stars in elliptical galaxies is heated by supernova explosions and produces outward-flowing galactic winds. Except possibly for the dust component, steady-state galactic winds are impossible to observe. However, some galactic winds have thermally unsteady cores which can be observed in optical emission lines. If the gas which is thermally unsteady remains ionized as it goes into free fall at the galactic center, objects more massive than stars tend to form. It is likely that nonthermal radio emission and optical line emission can occur only in those ellipticals with thermally unsteady galactic winds.

I. INTRODUCTION

This paper has been motivated by two questions: (1) Why do most elliptical galaxies show no evidence of interstellar matter? and (2) Why do some elliptical galaxies produce nonthermal radio sources?

The general absence of interstellar gas in ellipticals is a remarkable fact in view of the many different kinds of stars which are known to lose mass. According to the review of Osterbrock (1962), only 15 percent of the moderate or giant ellipticals have emission features characteristic of low-excitation galactic diffuse nebulae. Although this percentage could be increased if a special effort were made to find [O II] λ3727 in ellipticals (Osterbrock 1960), the general shortage of gas in these galaxies is surprising. The total rate of mass loss from red giants, planetary nebulae, binary stars, pulsating stars, and supernovae could easily amount to 0.1–1.0 $\mathfrak{M}_\odot$ year^{-1}. At this rate 10^9–10^{10} $\mathfrak{M}_\odot$ of gas would be ejected in a Hubble time. If a tiny fraction of this gas were ionized, it would be easily observable; if it were neutral, it would almost certainly form into stars (§ VII), yet early-type stars are not observed. Where does all this gas go?

Part of the answer is that the gas ejected from the stars becomes very hot. In NGC 3379, a typical E0 galaxy, the rms stellar velocity is $v_* = 324$ km s^{-1} (Burbidge, Burbidge, and Fish 1961). The stars therefore have an equivalent temperature $T_* = \mu M v_*^2/3k = 2.1 \times 10^6$ ° K, where μM is the mean particle mass. When gas is expelled from stars moving at velocity v_*, it eventually passes through a strong adiabatic (nonradiating) shock and the gas temperature is increased to $T \approx T_*$. If the gas had the same temperature as the stars, it would occupy roughly the same volume of space and would be very difficult to observe. For example, if an elliptical galaxy at $d = 10$ Mpc contained 10^{10} $\mathfrak{M}_\odot$ of gas at $T = 2 \times 10^6$ ° K, the flux at 4 keV would be about 10^{-32} ergs cm^{-2} s^{-1} Hz^{-1}. This flux is a factor of 10^4 times smaller than that observed from M87. The hot gas would therefore be completely unobservable.

Unfortunately, this does not completely answer question (1) above, since the hot gas would cool off on a short time scale. We estimate that a gas cloud of 10^{10} $\mathfrak{M}_\odot$ at $T = 2 \times 10^6$ ° K associated with an elliptical galaxy would have a central density of at least 1 cm^{-3}. The cooling time (eq. [8]) at the galactic center is then $\lesssim 5 \times 10^5$ years! Once

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the gas had cooled, it would be observable either as an H II region or in the form of young stars.

This dilemma can be avoided if the gas has another heat source in addition to the adiabatic shocks that result from the relative motion of gas ejected from stars. We find that the observed rate of supernova explosions in ellipticals provides sufficient energy to heat the gas to temperatures greater than T_*. The gas then flows out of the galaxy and is completely unobservable (§ V). Such steady-state flows with appropriate source terms have also been discussed by Burke (1968) and Johnson and Axford (1971).

Even more astonishing is the fact that many elliptical galaxies become strong nonthermal radio sources. How does a self-gravitating system of 10^{11}–10^{12} weakly interacting mass points lead to ejection of radio clouds with masses $\sim 10^7\, \mathfrak{M}_\odot$ and energies $\sim 10^{55}$–10^{60} ergs? Our proposal in this paper is that nonthermal energy is released as a result of gravitational collapse of the central parts of galactic winds where the flow becomes "thermally unsteady" (§ VII). Steady-state galactic winds are not possible if the central gas density is so high that the gas cools by radiative losses before it can flow away from the galactic center (§ VI). In these cases the cool gas which forms near the galactic center undergoes free fall. We present arguments in § VII which suggest that this collapsing gas is initially ionized and that under these circumstances it is likely that objects more massive than stars will be formed. Such rotating massive objects have often been proposed to account for sources of extragalactic nonthermal energy. Rotation is not considered in detail in this paper—we confine our discussion to giant E0 and E1 galaxies.

II. THE FLOW EQUATIONS

We assume that gas appears in elliptical galaxies solely because of mass loss from stars—accretion of intergalactic gas is therefore ignored. We suppose that gas enters the galaxy at a rate proportional to the local stellar density ρ_*, which for E0 galaxies is a function of radius alone. The Eulerian equation of continuity for the spherically symmetric flow of gas is

$$\frac{\partial \rho}{\partial t} + \frac{1}{r^2}\frac{\partial}{\partial r}(r^2 \rho u) = \alpha \rho_* , \tag{1}$$

where $\alpha = \alpha_{sn} + \alpha_s$ is the total specific rate of mass loss from supernovae (α_{sn}) and all other stars (α_s for binaries, planetary nebulae, variable stars, etc.). Normally, $\alpha_s \gg \alpha_{sn}$, so $\alpha \approx \alpha_s = \mathfrak{M}_G^{-1} d\mathfrak{M}_G/dt$, where $\mathfrak{M}_G$ is the mass of the galaxy. The corresponding equation of motion is

$$\frac{Du}{Dt} = -\frac{1}{\rho}\frac{\partial P}{\partial r} - \frac{G\mathfrak{M}(r)}{r^2} - \alpha \frac{\rho_*}{\rho} u , \tag{2}$$

where D/Dt is the Lagrangian time derivative. The gravitational force depends on $\mathfrak{M}(r) = \mathfrak{M}_*(r) + \mathfrak{M}_g(r)$, the mass of stars and gas interior to r, but normally $\mathfrak{M}_g(r) \ll \mathfrak{M}_*(r)$. Poisson's equation is implicitly satisfied with this definition of $\mathfrak{M}(r)$. The final term in equation (2) represents a momentum sink as the inertia of new gas, injected with zero mean velocity, drains momentum from the total flow. The pressure

$$P = \frac{2k}{M}\rho T \tag{3}$$

corresponds to a fully ionized gas of pure hydrogen. Here M is the proton mass.

The exact manner in which gas escapes from the stars and merges with the local flow is rather complex. The detailed interaction of recently ejected gas with the ambient medium undoubtedly differs if the gas is ejected in bulk, as in a planetary nebula, or over a long period of time, as with a stellar wind or pulsating variable. Nevertheless, its ultimate fate may be quite similar. If a star with typical velocity $v_* \approx 100$–400 km s^{-1}

becomes a planetary nebula, then both the star and the nebula will move together through a tenuous interstellar medium for a very long time ($\gtrsim 10^4$ years if the ambient density is $n_e \lesssim 100$ cm^{-3}). Ultimately, however, the density in the nebula will decrease as it expands away from the star and it will undergo shock interactions with gas from other stars or the ambient medium. In giant elliptical galaxies the random stellar velocities are so high that the shocks in the ejected gas raise the gas temperature to very high values ($\gtrsim 10^6$ ° K) where radiative losses are negligible. We therefore assume that all the gas ejected from stars eventually passes through an adiabatic shock, and its thermal temperature T_s is then equal to the effective kinetic temperature of the stars $T_* = Mv_*^2/6k$. The final thermal temperature of the gas depends almost entirely on the motion of the stars and is practically independent of the gas temperature as it leaves the star. The one important exception to this pattern is the gas ejected in supernova remnants—here the ejection velocity v_{ej} (typically $\approx 20{,}000$ km s^{-1}) is so much greater than v_* that the equivalent temperature after thermalization with shock waves $T_{sn} = Mv_{ej}^2/6k$ is determined by the ejection process itself. Since $T_{sn} \gg T_s = T_*$, the gas can be heated appreciably by supernovae; for this reason we keep the rather small mass loss from supernovae separate from all other forms of stellar mass loss.

With this simple model for stellar mass loss, the conservation of thermal energy is

$$\frac{DE}{Dt} - \frac{P}{\rho^2}\frac{D\rho}{Dt} = \frac{H(T)\rho}{M^2} + \frac{\rho_*}{\rho}\left[\frac{3k}{M}(\alpha_s T_s + \alpha_{sn} T_{sn}) + (\alpha_s + \alpha_{sn})\left(\frac{u^2}{2} - E - \frac{P}{\rho}\right)\right], \quad (4)$$

where $E = 3kT/M$ is the thermal energy per gram. The terms on the left in equation (4) are the usual equations for adiabatic gas flow. Departures from adiabatic flow arise because of radiative heating and cooling and addition of new gas with different effective temperatures. We approximate the rate coefficient for radiative heating with

$$H(T) = 1.5 \times 10^{-20} T^{-1\ 0} - 1.26 \times 10^{-25} T^{0\ 3} - 10^{-21} \exp\{-0.565[\ln(5 \times 10^{-6} T)]^2\} \text{ ergs cm}^3\text{ s}^{-1}. \quad (5)$$

The first term in equation (5) represents an approximation of the photoelectric heating by ultraviolet radiation (probably from horizontal-branch OB stars), which we assume to be present (see § III). The equilibrium temperature for gas which is primarily heated by ultraviolet starlight, given by $H(T_{eq}) = 0$, is slightly less than 6000° K. The second term in equation (5) represents both cooling by forbidden-line emission in the range 10^3 ° $\lesssim T \lesssim 10^4$ ° K and cooling by thermal bremsstrahlung for 5×10^6 ° $\lesssim T \lesssim 10^8$ ° K. The final term approximates the cooling by collisional excitation of oxygen and other ions for 10^4 ° $\lesssim T \lesssim 5 \times 10^6$ ° K (Cox and Tucker 1969; Cox and Daltabuit 1971). We assume that the gas is optically thin to all radiation in the cooling terms and that the level of ionization for $T \gtrsim 10^4$ ° K immediately adjusts to its collisional-equilibrium value determined by the local kinetic temperature. The cooling rate $|H(T)|$ has a maximum at $T_m = 2 \times 10^5$ ° K, and any gas with temperatures near T_m will tend to cool rapidly. The last term in equation (4) accounts for the source of thermal energy as new gas enters the flow from stars and supernovae. In addition to the thermal-energy content of the newly ejected gas, the relative streaming of the ambient gas ultimately converts an additional $\frac{1}{2}\alpha u^2(\rho_*/\rho)$ ergs g^{-1} s^{-1} into thermal energy.

We have ignored the dynamic effects of cosmic rays and magnetic fields. Cosmic rays produced in supernova explosions permeate the galaxy and lead to pressure gradients which couple with the ionized interstellar gas primarily through magnetic fields. For gas flows in elliptical galaxies, however, the cosmic-ray pressure is likely to be small if most of the energy released by the supernovae is in the kinetic energy of the ejected shells. Similarly, if most of the interstellar magnetic field present in ellipticals has been ejected from the stars along with the gas, then magnetic forces can be ignored since gas pressure generally exceeds the magnetic field pressure in stellar atmospheres. Finally, we note that dust may be expelled from red giants—some of the consequences of this are discussed in § V.

a) Boundary and Initial Values

Initially, we assume that the elliptical galaxy has no interstellar gas, but at time $t = 0$ gas begins to flow from the stars at a rate $\alpha\rho_*$ g cm^{-3} s^{-1}. After a small time, the gas density is proportional to the stellar density, i.e., $\rho(r, \Delta t) = \alpha\rho_*(r)\Delta t$, where Δt is the first time step in the numerical integration. The gas is assumed to be injected at zero mean velocity, $u(r, 0) = 0$.

Partly for mathematical convenience, we assume that the initial gas temperature distribution is determined by the static solution of equation (4). After a simple rearrangement of terms, this condition leads to

$$T(r, \Delta t) - \tfrac{3}{5}T_0 = \frac{H[T(r, \Delta t)]\alpha\rho_*(r)(\Delta t)^2}{5kM}, \tag{6}$$

where

$$T_0 = (\alpha_s T_s + \alpha_{sn} T_{sn})/\alpha\,. \tag{7}$$

We have chosen a value of Δt which is sufficiently small that the right-hand side of equation (6) can be ignored and the initial gas temperature is constant ($\frac{3}{5}T_0$) throughout the galaxy.

As usual, the boundary condition at the origin requires that $u(0, t) = 0$. At the outer boundary of the gas flow we ignore the small pressure of the intergalactic medium by setting the gas pressure equal to zero. We show below that the external intergalactic medium has almost no effect on the flow of gas within the galaxy.

III. DISCUSSION OF PARAMETERS

a) Model Galaxy

Among the various functions and constants which appear in the equations of § II which must be specified, those which describe the properties of the galaxy are most accurately determined from the observations. As a model for the distribution of stellar mass, we have chosen NGC 3379, a well-observed and relatively nearby giant E0 galaxy. Miller and Prendergast (1962) have determined the stellar mass density $\rho_*(r)$ of NGC 3379 for $r \leq R_G = 14$ kpc (see Fig. 1). These densities are based on an assumed distance of $d = 10$ Mpc, which we also adopt here. At the galactic center, the density of stellar mass is $\rho_{*c} = 440\ \mathfrak{M}_\odot$ pc^{-3} or 1.87×10^4 protons cm^{-3}, but for $r = 14$ kpc the density has dropped to $2.2 \times 10^{-6}\ \rho_{*c}$. We assume also that $\rho_*(r) = 0$ for $r > R_G$. In § V we show that truncation of the stellar mass in this manner has no effect on the galactic-wind solutions for $r < R_G$. The total mass of the galaxy within R_G is $\mathfrak{M}_G = 0.91 \times 10^{11}\ \mathfrak{M}_\odot$.

b) Supernovae

Supernovae are the only directly observed form of stellar mass loss in giant elliptical galaxies. The frequency of supernova outbursts in ellipticals has recently been reexamined by Katgert and Oort (1967), who found an average frequency ν_{sn} of one supernova every 40 years. Earlier determinations of ν_{sn} were lower because of insufficient correction for incompleteness effects. In this paper we assume one supernova every 50 years per 10^{11} stars. Apparently, all supernovae with known light curves in E or S0 galaxies are of Type I (Bertola and Sussi 1965), normally associated with Population II stars. The properties of galactic Type I supernovae have been reviewed by Minkowski (1967), who suggests that a kinetic energy $E_{ej} \approx 10^{51}$ ergs is associated with the expanding shell. A typical ejection velocity $v_{ej} \approx 20{,}000$ km s^{-1} corresponds to a shell of mass

$\mathfrak{M}_{ej} = 0.25\, \mathfrak{M}_\odot$. With these assumed values, the specific rate of mass ejection by supernovae is $\alpha_{sn} = \mathfrak{M}_{ej}\nu_{sn}/\mathfrak{M}_G = 1.6 \times 10^{-21}$ s^{-1}. This value of α_{sn} is uncertain by a factor of at least 2, but probably not greater than 10. The equivalent temperature associated with the kinetic energy of the supernova shells is $T_{sn} = Mv_{ej}^2/6k = 8 \times 10^9$ ° K, where the ejected gas is assumed to be completely ionized. T_{sn} is uncertain to about an order of magnitude.

The velocity of ejection of supernova shells v_{ej} far exceeds the escape velocity ($\sim$500 km s^{-1}) from the centers of giant ellipticals. Nevertheless, none of these shells directly escapes, even if no additional interstellar gas is present. During the time necessary for a supernova shell to cross the galaxy $R_G/v_{ej} \approx 7 \times 10^5$ years, at least 10^4 other supernovae have occurred in the galaxy. In the absence of an ambient medium, the shells from all these explosions would interact and ultimately thermalize. Since we assume that $\alpha_s \gg \alpha_{sn}$, the supernova shells are further impeded by the presence of larger amounts of gas ejected from other stars.

c) *Stellar Mass Loss*

Mass ejection from stars can occur by many different mechanisms; among the most common possibilities are planetary nebulae, ejection associated with stellar pulsation or binary star gaseous streams, mass loss during contraction to the main sequence, supernovae, novae, and stellar winds of various types (see Deutsch 1969 for a recent review). It is difficult, however, to make a good quantitative estimate of the total rate of mass loss from all these sources. In our Galaxy the rate of star deaths indicates that stars may return about 0.1–10 $\mathfrak{M}_\odot$ to the interstellar medium each year. Observations of the outflow of gas from the nucleus of our Galaxy (Burbidge 1970 and references therein) suggest mass-loss rates of $\sim$1 $\mathfrak{M}_\odot$ year^{-1} from stars in the nuclear region alone. It is likely, however, that this outflowing gas may have its origin outside the galactic center.

In giant elliptical galaxies the total rate of stellar mass loss $d\mathfrak{M}_G/dt$ is extremely difficult to estimate. Presumably, star formation in ellipticals occurred over a shorter period of time than in spirals and the resulting stellar mix is therefore quite dissimilar. Sandage (1957) estimated the total rate of mass loss in the globular cluster M3 from a detailed study of the luminosity function. If these results are adjusted to the higher mass-to-light ratios appropriate to ellipticals ($\mathfrak{M}/L \sim 50$), the galactic mass-loss rate is $d\mathfrak{M}_G/dt \approx 0.1$–$1.0\, \mathfrak{M}_\odot$ year^{-1}, with considerable uncertainty. The fraction f_{RG} of stars in ellipticals which are red giants can also be used to estimate the rate of mass loss. The optical spectrum of many ellipticals is similar to that of the nucleus of M31 (Oke and Sandage 1968). The stellar makeup of the central regions in M31 has been recently studied by Spinrad and Taylor (1971), who find that G III subgiants and K III giants are present with $f_{RG} \approx 1.4 \times 10^{-3}$. If these stars spend about 10^9 years as giants and lose half their mass by gas ejection, then the corresponding rate of mass loss (scaled to a galaxy of 10^{11} stars) is $d\mathfrak{M}_G/dt \approx 0.1\, \mathfrak{M}_\odot$ year^{-1}. This estimate of the rate of mass loss is very uncertain and would be increased further if M giants are also present in M31 in reasonably large numbers. In any case, it would appear from these estimated values of $d\mathfrak{M}_G/dt$ that the specific rate of stellar mass loss $\alpha_s = 3.5 \times 10^{-19}\, d\mathfrak{M}_G/dt$ s^{-1} (where $d\mathfrak{M}_G/dt$ is in $\mathfrak{M}_\odot$ year^{-1}) will considerably exceed α_{sn}.

The equivalent thermalized temperature T_s of the bulk of the gas ejected from stars will be almost entirely determined by the relative motions of the stars. If the gas ejected from various stars collides through adiabatic shocks, then $T_s = T_* \approx 2.1 \times 10^6$ ° K for NGC 3379. (We assume that the mean stellar velocity v_* is independent of radius.) Although it would seem unlikely, if there are some radiative losses behind the shocks, then $T_s < T_*$. Fortunately, however, the exact value of T_s does not greatly influence the results discussed below, provided the supernova blast waves involve nearly adiabatic shocks. This follows from equation (4) and the fact that $T_s \ll T_{sn}$.

d) *Photoionization*

In many globular clusters there are a number of early-type stars on or above the horizontal branch which provide some ionizing radiation ($h\nu > 13.6$ eV). If the stellar population in giant ellipticals approximately resembles that in M3, a globular cluster with a rather high number of OB horizontal-branch stars, then there would be more than enough ionizing radiation to account for the observed [O II] emission-line flux; this is true even in NGC 1052 and NGC 4278, which have strong emission features (Minkowski and Osterbrock 1959; Osterbrock 1962). Indeed, in his simulation of the stellar population in several ellipticals, Wood (1966) suggested that OB horizontal-branch stars could account for the observed optical ultraviolet flux. The numbers of these hot stars required by Wood would also ionize enough gas to account for all the observed emission. More direct observational evidence for the presence of hot stars has recently become available from preliminary results obtained with OAO II. Code (1969) and Savage (private communication) point out that the continuum flux (per Å) from several elliptical galaxies is greater at 1700 Å than at 3000 Å. In view of these observations, and because there is no strong evidence (except perhaps that of Robinson and Koehler 1965) for the presence of neutral hydrogen in elliptical galaxies, we have decided to ignore the possibility of neutral hydrogen in the computed gas flows.

IV. COMPUTING PROCEDURE

Lagrangian forms of equations (1), (2), and (4) were solved by the standard numerical method described by Richtmeyer and Morton (1967). The source terms in equations (2) and (4) can be differenced in a straightforward manner. For example, we have solved equation (1) in the form

$$\rho_{j-1/2}{}^{n+1} = \frac{(R_j{}^n)^3 - (R_{j-1}{}^n)^3}{(R_j{}^{n+1})^3 - (R_{j-1}{}^{n+1})^3}\, \rho_{j-1/2}{}^n + \tfrac{1}{2}\alpha(\rho_{*j-1/2}{}^n + \rho_{*j-1/2}{}^{n+1})\Delta t^{n+1/2}\,,$$

where j and n denote space and time intervals. The Courant condition was always satisfied. Occasionally it was necessary to rezone the space grid to accelerate the calculations or to improve the accuracy—the rezoning procedure conserved mass as well as internal energy. At the outer edge of the gas flow, zones were dropped from the calculation. The procedure we have used conserves mass to within several percent (i.e., the total mass of gas at any time must be $\alpha\mathfrak{M}_G t$) and thermal energy to within about 10 percent.

V. STEADY-STATE HOT GALACTIC WIND

a) *Computed Model*

The solutions to equations (1), (2), and (4), with boundary conditions as discussed in § II, are of two general types: (i) steady-state outward-flowing hot ($T \gtrsim 10^6$ °K) galactic winds, and (ii) galactic winds with "thermally unsteady" cores of low temperature ($T \lesssim 10^4$ °K). In this section we discuss solutions of the first type only, reserving § VII for thermally unsteady flows. As an example of this first type of solution, we describe the gas flow which results from the following choice of parameters: $\alpha_s = 1.2 \times 10^{-18}$ s^{-1} (corresponding to a mass-loss rate $d\mathfrak{M}_G/dt = 3.4\ \mathfrak{M}_\odot$ year^{-1}), $T_s = T_* = 2.1 \times 10^6$ °K, $\alpha_{sn} = 1.6 \times 10^{-21}$ s^{-1}, and $T_{sn} = 8 \times 10^9$ °K.

The gas density $n(r, t) = \rho(r, t)/M$, velocity $u(r, t)$, and temperature $T(r, t)$ for this solution are shown in Figures 1–3, respectively. Figure 1 shows how the density gradually increases throughout the galaxy as the stars expel gas into the flow. Initially, $\rho(r, t)$ has the same distribution as the stellar distribution $\rho_*(r)$ (also shown in Fig. 1), but, with time, a unique steady-state distribution is reached. Steady-state conditions set in first at the galactic center, but after $\tau_{ss} \approx 5 \times 10^7$ years, the entire flow within $R_G = 14$ kpc

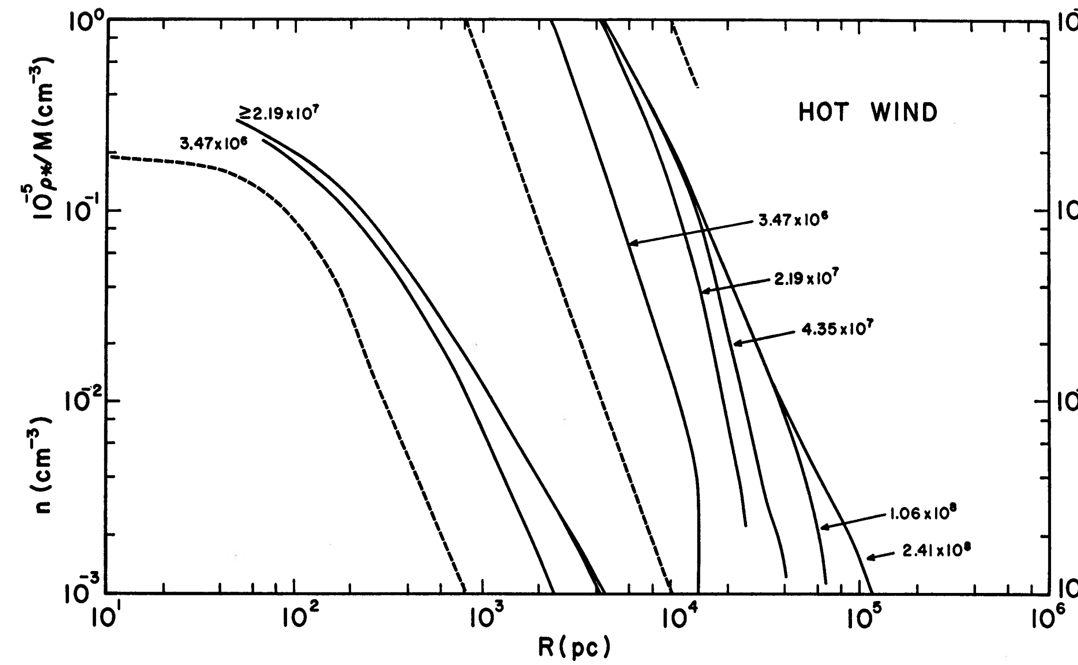

FIG. 1.—Variation of gas density $n(r, t)$ with time since the beginning of the calculation. Each curve is labeled with the time in years. The hot-wind solution is for $d\mathfrak{M}_G/dt = 3.4\ \mathfrak{M}_\odot$ year^{-1}, $T_s = T_*$, and $\alpha_{sn}T_{sn} = 1.28 \times 10^{-11}$ ° K s^{-1}. The smoothed-out stellar density is shown as a dashed curve. Curves entering the top of the diagram refer to the scale on the right.

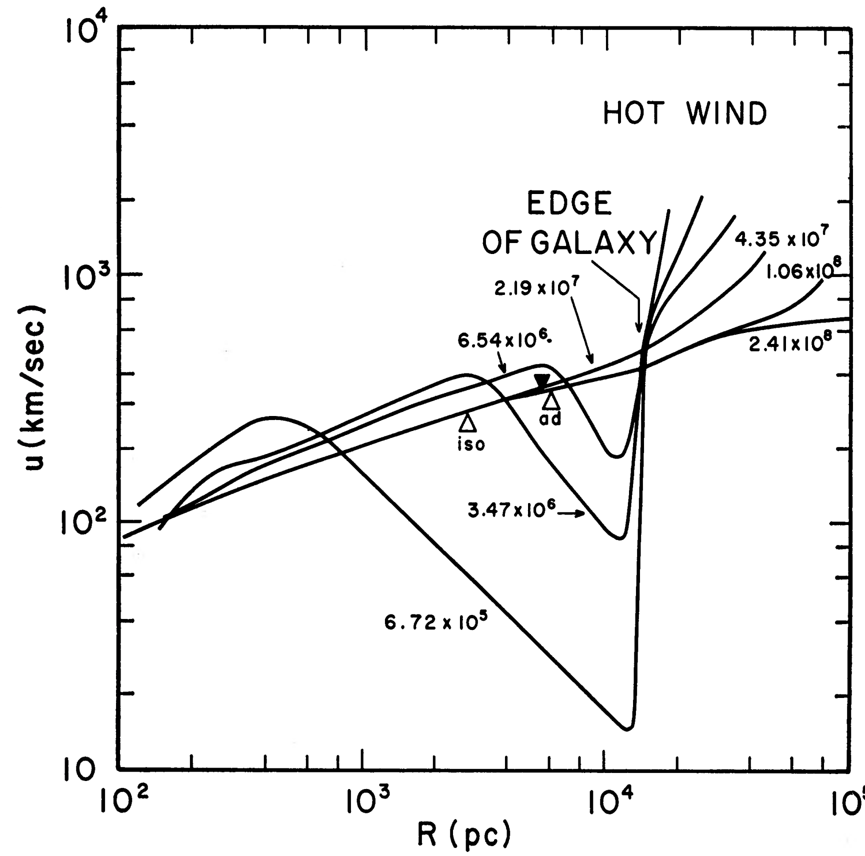

FIG. 2.—Variation of gas velocity $u(r, t)$ with time (in years) for hot-wind solution. Open triangles mark the adiabatic and isothermal sonic points; filled triangle shows where the flow velocity equals the local escape velocity.

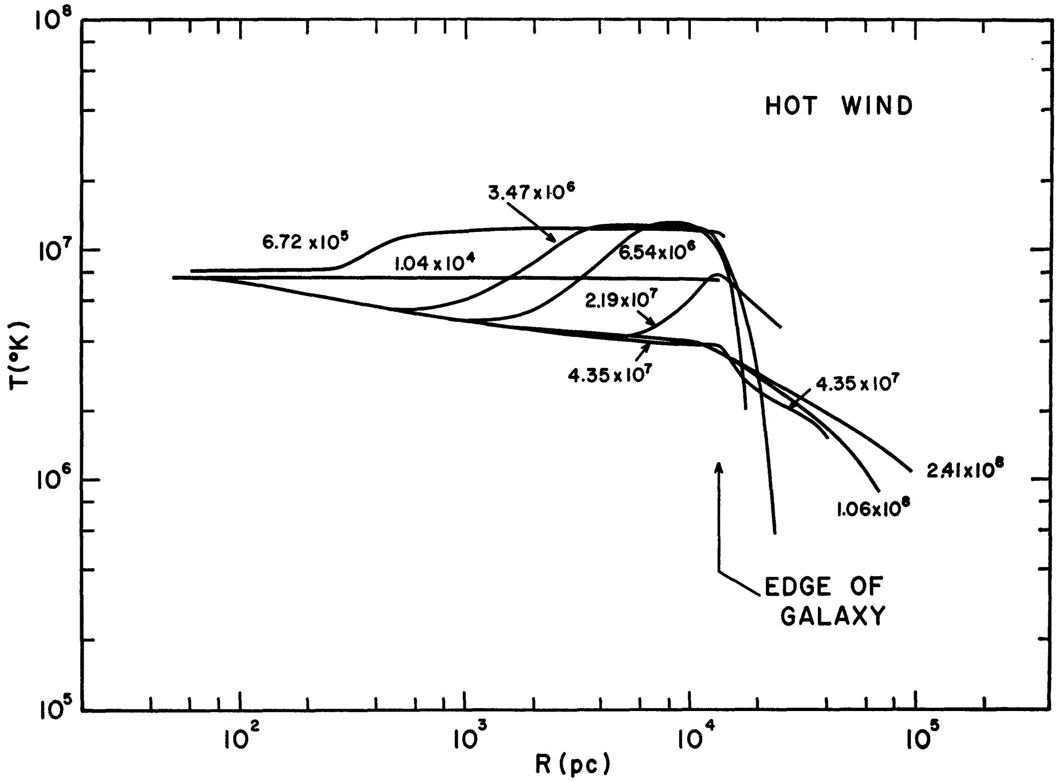

FIG. 3.—Variation of gas temperature $T(r, t)$ with time (in years) for hot-wind solution

remains unchanged. The final value of the gas density at the center is $n_c = 0.4 \pm 0.1$, where the uncertainty arises from the fact that the density is evaluated only at the middle of the first spatial zone. The transient velocity flow shown in Figure 2 begins at the center where the pressure gradient is largest. A compression wave moves outward through the galaxy, growing in strength (but never quite shocking) until it merges with freely expanding gas beyond R_G. The gas becomes supersonic with respect to the local isothermal (adiabatic) sound velocity at 2.5 kpc (6.23 kpc). Since the flow is supersonic in the outer parts of the galaxy, no information can move upstream; therefore, truncation of the galactic star density at $R_G = 14$ kpc has no effect on the solution within R_G. The flow velocity first exceeds the local escape velocity at $R_e = 5.6$ kpc, also indicated in Figure 2. Finally, Figure 3 shows the development of the gas temperature with time. After the first time step ($\Delta t = 10^4$ years) the temperature is constant, $T(r, \Delta t) = \frac{3}{5}T_0 = 7.6 \times 10^6$ ° K (§ II*a*). Following an initial transient phase, where the gas is heated by the compression produced as new gas is added into the flow, the temperature eventually obtains a monotonically decreasing profile through the range $(7.5–3.0) \times 10^6$ ° K for $r < R_G$. (This temperature gradient is so small that thermal conduction can safely be ignored in eq. [4].) The excess temperature above $T_* = 2.1 \times 10^6$ ° K is responsible for the general outflow of gas. The temperature decreases beyond R_G since the gas is no longer subject to supernova heating when $\rho_* = 0$, according to the assumption involved in equation (4).

Several global properties of the hot-wind solution are of interest. The "flushing time," i.e., the time for a given fluid element to move with the gas from the center to R_G, is $\tau_{fl} \approx 4 \times 10^7$ years. The time to reach steady-state flow out to R_G, τ_{ss}, is virtually the same as τ_{fl}. The total amount of gas interior to R_G when steady-state conditions are reached is 7.7×10^7 $\mathfrak{M}_\odot$. The gravitational binding energy of the gas in the gravitational field of the stars is $E_{grav} = -1.58 \times 10^{56}$ ergs.

b) Observability

Another important consideration is the observability of hot-wind solutions such as the one discussed above. Since interstellar gas has not been detected in the majority of elliptical galaxies, we consider a successful hot wind to be one that is completely undetectable by any direct observation. The model shown in Figures 1–3 seems to satisfy this definition of success by a wide margin. It is certain that no [O II] λ3727 radiation can result from a gas with $T \gtrsim 10^6$ ° K, since oxygen would exist only in very high stages of ionization. Assuming that the flux distribution of starlight is similar to that given by Oke and Sandage (1968) for giant ellipticals, we find that the ratio of observed Hβ flux (Gould 1971) to stellar continuum flux in a 10 Å band pass (unreddened and centered on Hβ) is $\sim 8 \times 10^{-7}$. Alternatively, if all the iron in the hot wind is in the Fe^{+14} stage of ionization, then the flux in the coronal line [Fe XV] λ7060 can be estimated and integrated over the flow. For this line the ratio of line flux to 10 Å of stellar continuum is $\sim 7 \times 10^{-6}$. If only the central 500 pc of the galaxy is observed (corresponding to a diaphragm of about 10″ with the adopted distance $d = 10$ Mpc), these flux ratios can be increased only by a factor of about 10. Since the signal-to-noise ratios for emission lines would be even smaller for lower values of the mass loss rate (we regard $d\mathfrak{M}_G/dt = 3.4$ $\mathfrak{M}_\odot$ year^{-1} as a rather high value), we conclude that optical observation of hot galactic winds is impossible. This explains why most elliptical galaxies which are not too highly flattened show no optical evidence of diffuse gas, even if its component stars are losing mass at significant rates.

Observation of galactic winds outside the optical window is also extremely difficult. The X-ray bremsstrahlung at 4 keV integrated over the galaxy is 1.2×10^{-34} ergs cm^{-2} s^{-1} Hz^{-1}, which is very much less than the 4-keV flux of $\sim 9 \times 10^{-28}$ observed from M87 by Bradt *et al.* (1967) and Bowyer *et al.* (1970). The X-ray power of individual galactic winds in 1–10 keV is $\sim 10^{36}$ ergs s^{-1}, about 10^3–10^4 times less than the X-ray power of our own Galaxy, and far below the level which would be needed to account for the diffuse X-ray background (Gould and Burbidge 1963; Silk 1968). Finally, radio bremsstrahlung from the galactic wind is $\sim 10^{-6}$ f.u. at 10 cm, i.e., completely unobservable.

It is therefore not surprising that the majority of elliptical galaxies show no evidence of interstellar gas. Even the combined luminosity of $\sim 10^3$ galactic winds in the Coma cluster would be quite unobservable. However, the gas lost from each component galaxy in the Coma cluster during a Hubble time would not be able to support itself thermally at $\sim 10^6$ ° K against the gravitational field of the cluster. In this event, the gas expelled from each galaxy would collect at the center of the cluster and form into an optically visible object of galactic dimensions. More likely, however, after the gas has been ejected from the galaxies it is heated further to a temperature ($\sim 10^8$ ° K) equivalent to the thermal velocity of the component galaxies ($\sim 10^3$ km s^{-1}). This heating would be accomplished by adiabatic bow shocks associated with each galactic wind in the cluster. In this event, the hot intercluster medium would occupy roughly the same volume as the cluster galaxies and would be difficult to detect, even if the rate of mass loss per galaxy exceeds 1 $\mathfrak{M}_\odot$ year^{-1} (Felten *et al.* 1966; Woolf 1967). Additional sources of diffuse radiation will therefore be necessary to account for the nonthermal radio halo in the Coma cluster (Willson 1970) and the recently reported X-ray detection of Coma in the 0.7–4-keV band (Fritz *et al.* 1971).

Planetary nebulae, which may be responsible for a significant fraction of the stellar mass loss in ellipticals, would also be unobservable. The total luminosity of 10^4 planetaries each with mass 0.5 $\mathfrak{M}_\odot$ with $N_e = 5 \times 10^2$ cm^{-3} and $T = 10^4$ ° K is $L_\beta = 3 \times 10^{38}$ ergs s^{-1}. The luminosity of galactic starlight at Hβ is $L_{*\lambda} = 8.3 \times 10^{39}$ ergs s^{-1} Å^{-1}.

With a bandwidth of 10 Å, the signal-to-noise ratio for observing Hβ in emission from all the planetaries is 3.6×10^{-3}. This could be increased by only a factor of $\sim$10 if the central 500 pc were observed. Such a low level of emission could never be disentangled from stellar absorption features. Also, the general level of excitation is higher in planetaries than that observed in the centers of ellipticals. Osterbrock (1962) has suggested that the total nonthermal emission from all the supernova remnants within a giant elliptical could be observed at radio frequencies. Although this may be unlikely, it cannot be strictly ruled out at present.

One possible means of observing the galactic flow would be an indirect or direct measurement of the dust which is probably ejected from red giants along with the gas. The dust would move along with the general flow of gas. The reddening observed near the central regions of many ellipticals by Tifft (1969, 1970) may be due in part to the presence of interstellar dust. If we assume, for example, that all of the carbon in the hot wind is in the form of small graphite particles having radii $a \approx 0.05$–$0.10\ \mu$, then the optical depth at 5000 Å along a radius is $\tau_d \approx 0.1$ for the hot wind shown in Figures 1–3. This value of τ_d is perhaps consistent with the small amount of reddening observed in NGC 3379 by Miller and Prendergast (1962). In view of the many uncertainties involved in the estimate of τ_d, it is at least possible that still larger values could be found in other ellipticals. It may also be possible to observe infrared emission from these grains directly. As far as the authors know, however, no infrared measurements have been made of normal ellipticals (i.e., with no optical or radio evidence of interstellar gas). We note that the sputtering destruction time for uncharged graphite grains is about a tenth of the time ($\sim 10^6$ years) for the gas flow to cross the central regions of the galaxy. However, if the grains have a relatively small positive charge, they could survive (Mathews 1969).[1]

c) Further Discussion

Implicit in the gas-flow equations (1), (2), and (4) is the assumption of "microscopic evenness," i.e., we ignore the spatial cavitation which results from supernova blast waves and we assume that the equivalent thermal energy of the supernovae is deposited smoothly in the gas. In fact, however, the supernova energy is released at specific points in the galaxy. We have therefore attempted to estimate the mean interval of time $t_{\rm bl}$ between passages of a supernova blast wave at a given point in the galaxy. For consistency, $t_{\rm bl}$ must be small compared to the radiative cooling time $t_{\rm cool}$. From the theory of adiabatic blast waves we find that $t_{\rm bl} = 2 \times 10^6 n_e^{3/11} n_s^{-5/11}$ years, where n_e and n_s are the densities (in protons cm^{-3}) of gas and stellar matter, respectively, and the properties of supernovae are those given in § III*b*. The cooling time is easily found by comparing the first and third terms in equation (4):

$$t_{\rm cool} = \frac{3kT}{n_e|H(T)|} . \qquad (8)$$

When $t_{\rm cool}$ is evaluated at $T = 5 \times 10^6\ ^\circ$ K, we find that $t_{\rm cool} \approx 500 t_{\rm bl}$ at the galactic center in the hot wind. Since $t_{\rm cool}/t_{\rm bl}$ increases rapidly with distance from the center of the galaxy, we conclude that the gas does not have time to cool below T_s before it is struck by another strong adiabatic shock. The assumption of uniform spatial deposition of supernova energy does not appear to be seriously in error.

If elliptical galaxies are losing gas by hot winds, a low-density intergalactic medium is produced. The density of this medium is easy to estimate. Assume that the density of galactic matter in the Universe is 7×10^{-31} g cm^{-3} and that 15 percent of this mass is in ellipticals. Then each elliptical galaxy (with mean mass $10^{11}\ \mathfrak{M}_\odot$) losing 1 $\mathfrak{M}_\odot$ each year during a Hubble time results in an intergalactic gas density of only $n_{\rm ig} \approx 10^{-8}$ cm^{-3}.

[1] The sputtering yield slope should be $s = 2.5 \times 10^{-3}$ eV^{-1}, rather than the value given in Mathews (1969). This correction has been pointed out to us by Mr. Per Aannestad.

This is substantially less than the upper limit ($n_{ig} \lesssim 10^{-6}$ cm^{-3}) of an ionized intergalactic medium set by Lα observations in distant quasi-stellar objects. We note, however, that if the intergalactic gas density were as high as 10^{-6} cm^{-3}, the hot galactic wind would not be seriously affected by the motion of the galaxy through the intergalactic gas. The dynamic pressure produced by a galaxy moving at velocity v_G is $P_{dy} = \rho v_G^2$. With $v_G = 500$ km s^{-1} and $\rho = 10^{-6} M$ g cm^{-3}, we find that $P_{dy} = 4 \times 10^{-15}$ dyn cm^{-2}, which is much smaller than the gas pressure in the outer parts of the galactic wind; $P(R_G) = 1.5 \times 10^{-13}$ dyn cm^{-2}. The bow shock where the wind collides with the intergalactic medium would be at least 50 kpc away from the center of the galaxy. The outflow of gas from elliptical galaxies therefore makes accretion of low-density intergalactic gas quite difficult.

Finally, we note that the gravitational binding energy of the hot wind $|E_{grav}| = 1.6 \times 10^{56}$ ergs is very much less than the typical energy $\sim 10^{60}$ ergs involved in the ejection of strong radio sources. If a galaxy with a hot wind were to become a radio source, then it would seem likely that the whole galaxy would be outgassed. In this event our initial condition of zero gas density (§ II*a*) would be appropriate, and the wind would again reestablish steady-state flow in $\sim 5 \times 10^7$ years. However, for a number of reasons which we discuss below, it seems more likely that a galaxy which has a hot wind will never become a radio source. The exotic situation in the nuclei of some ellipticals which results in the release of high kinetic and thermal energies may develop only when the gas flow is thermally unsteady. We discuss thermally unsteady galactic winds in the following sections.

VI. GENERAL PROPERTIES OF STEADY AND UNSTEADY WINDS

Steady-state galactic winds as discussed in the previous section do not occur for all combinations of the parameters α_s, $\alpha_{sn}T_{sn}$, T_s, and $\rho_*(r)$. This is apparent from equation (4), where the two source terms depend on gas density and temperature in quite different ways. Because the radiative-cooling term depends on binary collisions of atomic particles, the volumetric cooling rate (ergs cm^{-3} s^{-1}) varies as ρ^2. On the other hand, the major source of heating depends on the term $\alpha_{sn}\rho_* 3kT_{sn}/M$ (ergs cm^{-3} s^{-1}), which is independent of gas density. For a given galaxy (i.e., T_s, $\alpha_{sn}T_{sn}$, and ρ_* fixed) as α_s is increased, the central value of the gas density will increase until ultimately radiative losses exceed the heating by supernovae; the gas then cools rapidly and a thermally unsteady flow develops as the gas collapses under the gravity field of the stars. Alternatively, if T_s, α_s, and ρ_* are fixed, but T_{sn} (or α_{sn}) is lowered, then the heating rate will ultimately fall below the cooling rate, again resulting in a thermally unsteady situation at the galactic center.

The properties of steady-state hot-wind solutions (§ V) at the galactic center as well as the condition for thermally unsteady flow can be roughly determined from equations (1), (2), and (4) without detailed calculations. Let R_c be the scale height of the gravitational field at the galactic center. Then for steady-state flow, equation (1) gives

$$\rho u/R_c \approx \alpha\rho_* , \qquad (9)$$

where u and ρ are the characteristic gas velocity and density within R_c. The central temperature T_c for the hot-wind solutions can then be found from equation (4) by replacing D/Dt with u/R_c:

$$\frac{3T_c u}{R_c} - \frac{2T_c u}{R_c} = \frac{H\rho}{Mk} + \frac{\alpha\rho_*}{\rho}(3T_0 - 5T_c) , \qquad (10)$$

where terms of $O(u^2)$ are neglected. From equations (9) and (10) we find

$$T_c = \frac{H\rho^2}{6Mk\alpha\rho_*} + \tfrac{1}{2}T_0 , \qquad (11)$$

where T_0 is given by equation (7). If parameters are such that the steady-state galactic wind is far from being thermally unsteady, then the first term on the right of equation (11) can be neglected, and the central temperature is

$$T_c = \tfrac{1}{2}T_0 = \tfrac{1}{2}(T_s + \alpha_{sn}T_{sn}/\alpha_s) . \tag{12}$$

For the values of α and $\alpha_{sn}T_{sn}$ used in § V, $T_0 = 1.07 \times 10^7$ ° K, so $T_c = 5.3 \times 10^6$ ° K, in excellent agreement with the value shown in Figure 3.

The central gas density is

$$n_c = \frac{\rho c}{M} = \frac{\alpha \rho_* R_c}{M c_s} , \tag{13}$$

where we have assumed that the characteristic gas velocity is the sound speed $c_s = (10kT_c/3\ M)^{1/2}$. In NGC 3379 the stellar density drops to $\sim$0.1 of its central value at $R_c = 250$ pc. We find from equation (13) that $n_c = 0.45$ cm^{-3}, in good agreement with the central density shown in Figure 1.

The condition for a thermally unsteady galactic wind can be found by comparing the left-hand side of equation (10) with the radiative-cooling term. Cooling and unsteady collapse occur whenever the cooling time of the gas (eq. [8]) is less than the time t_{flow} necessary for gas to flow out to R_c, i.e.,

$$t_{cool} = \frac{3kT_c}{n_c|H(T_c)|} < \frac{3R_c}{c_s} = t_{flow} , \tag{14}$$

where the factor 3 follows from equation (10). All the factors which occur in condition (14) can be expressed in terms of T_c. Since T_c depends on T_s, α_s, and $\alpha_{sn}T_{sn}$ (eq. [12]), condition (14) can be presented in a plot of α_s and $\alpha_{sn}T_{sn}$. Figure 4 shows such a plot for $T_s = 2.1 \times 10^6$ ° K, appropriate to NGC 3379. The solid line represents the condition $t_{cool} = t_{flow}$. Points representing the parameters chosen in §§ V and VII fall on either side of the instability line, as expected. Higher rates of mass loss or less energetic (or less frequent) supernovae favor the instability. In the region of Figure 4 where steady hot winds occur, we have drawn in lines of constant T_c determined from equation (12).

VII. THERMALLY UNSTEADY GALACTIC WIND

a) Computed Thermally Unsteady Galactic Wind

In order to demonstrate the nature of thermally unsteady galactic winds, we now discuss the solution which results when the calculations of § V*a* are repeated with $\alpha_{sn}T_{sn}$ lowered by a factor of 10. Specifically, we have taken $T_{sn} = 8 \times 10^8$ ° K and left α_{sn} unchanged—all other parameters are exactly the same as in § V*a*. Figures 5–7 show the initial variation of $n(r, t) = \rho(r, t)/M$, $u(r, t)$, and $T(r, t)$, respectively. The gas density (Fig. 5) begins to increase as in Figure 1, but for $t \gtrsim 2 \times 10^6$ years radiative losses dominate at the center of the galaxy, and the gas within $\sim$100 pc cools rapidly (in $\sim$10^5 years) to $T_{eq} \approx 6000$° K, as determined by photoionization (§§ II and III*c*). The gas, once cooled to 6000° K, undergoes free fall toward the galactic center. Since $\rho_* \gg \rho$, the self-gravity of the gas can be ignored for the initial stages of collapse. The central value of the gas density increases rapidly with time as the free fall develops.

The gas velocity shown in Figure 6 begins with an outward-moving compression wave, similar to that in Figure 2. When strong radiative cooling sets in, however, the gas velocity reverses and becomes negative in the whole central region of the galaxy. As soon as gas at the very center undergoes radiative cooling, the pressure forces become ineffective and the gas goes into free fall with hot gas farther out rushing inward to take its place. As the cool gas ($T \approx 6000$° K) approaches the center, an outward-facing shock forms. Following Richtmeyer and Morton (1967), we have used artificial viscosity Q as a numerical artifice which couples equations (2) and (4) across the shock. For this reason,

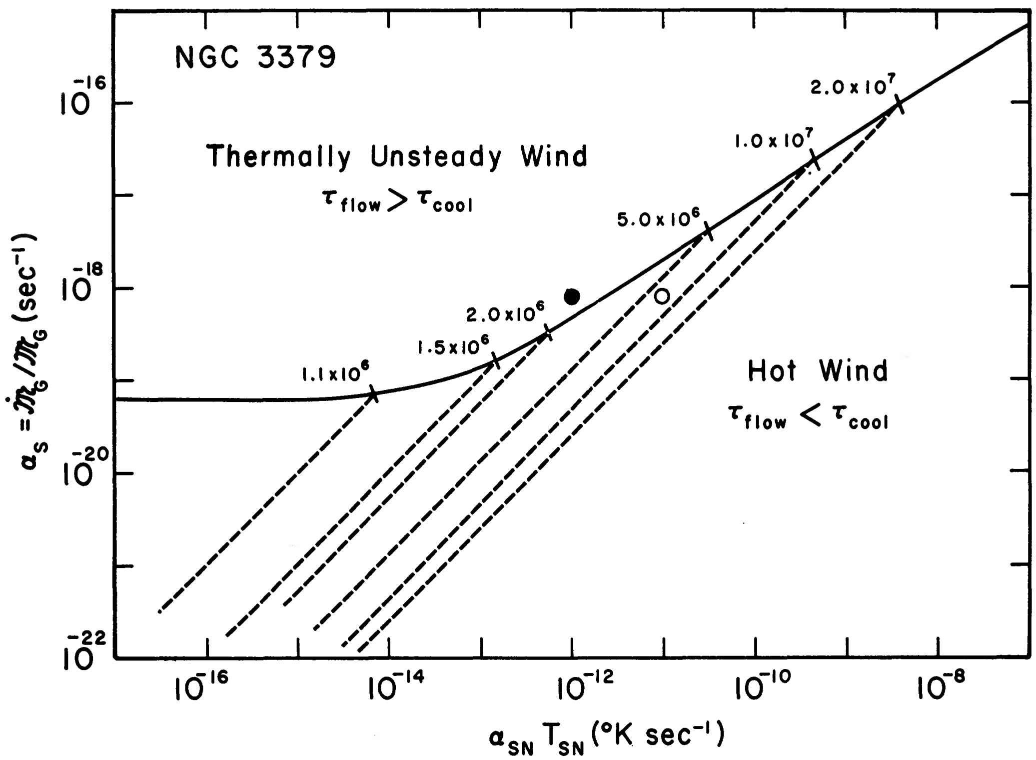

FIG. 4.—Regions of steady and unsteady winds in NGC 3379 for various α_s and $\alpha_{sn}T_{sn}$. Open and filled circles correspond to parameters chosen for detailed solutions in §§ V and VII, respectively. Numerical values of the central temperature T_c label each dashed line in the region of steady flow.

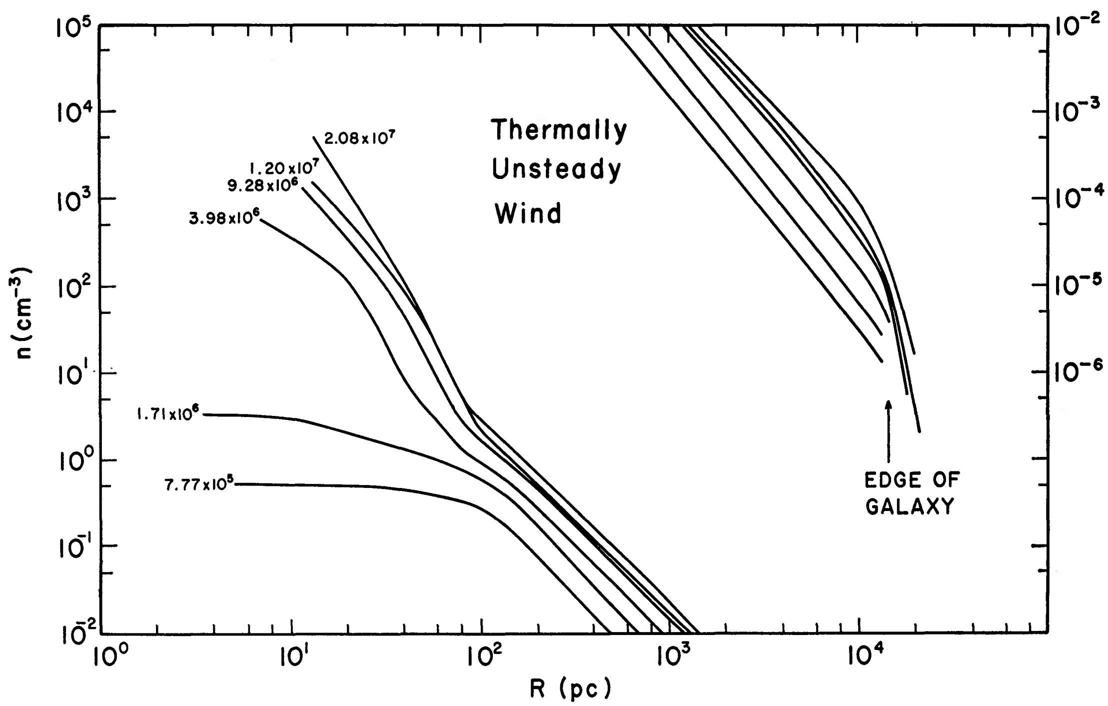

FIG. 5.—Variation of gas density $n(r, t)$ with time (in years) for thermally unsteady solution. Parameters are identical with the hot-wind solution except $\alpha_{sn}T_{sn} = 1.28 \times 10^{-12}$ ° K s^{-1}.

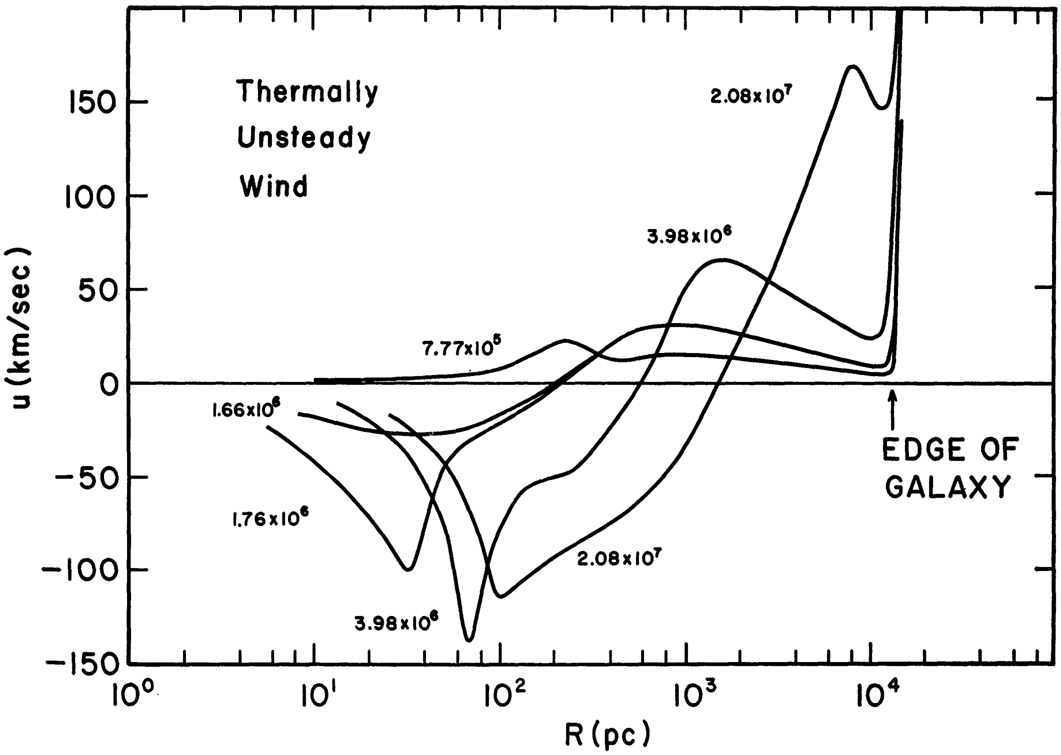

FIG. 6.—Variation of gas velocity $u(r, t)$ with time (in years) for thermally unsteady solution

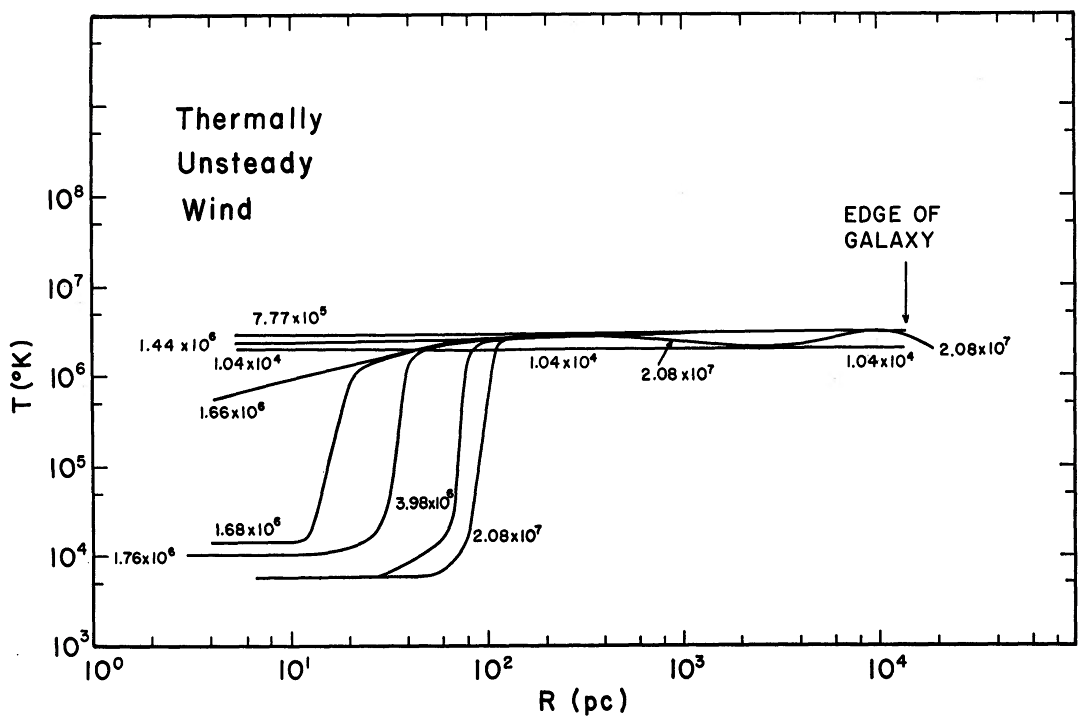

FIG. 7.—Variation of gas temperature with time (in years) for thermally unsteady solution

and because of our rather coarse spatial zoning, the solution is artificially smeared out in Figure 6 within the radius of maximum negative velocity. In fact, in a more refined calculation we would expect the maximum negative velocity (i.e., just before the gas enters the shock) to be somewhat larger than indicated in Figure 6. Because the hot gas ($T \approx 5 \times 10^6$ ° K) farther out ($r \gtrsim 100$ pc) is actually pushing the thermally unstable gas ($T \approx 6000°$ K) toward the origin, we would expect the maximum velocity of infall for a given element of gas to exceed the free-fall velocity determined at the point at which the radiative losses occurred.

The thermally unsteady region is best illustrated in Figure 7, which shows the gas temperature as a function of time. After 2×10^7 years the region of thermally unsteady flow extends out to $r \sim 100$ pc, and about $1.0 \times 10^7\ \mathfrak{M}_\odot$ of gas has a temperature less than 10^5 ° K. The mean density of this cool gas is $\sim$100 cm^{-3}, which is comparable with the observed gas densities in ellipticals (Osterbrock 1962). The fluxes in Hβ and [O II] λ3727 are easily observable at $d = 10$ Mpc with signal-to-noise ratios of about unity relative to 10 Å of stellar continuum. The exact value of the computed optical line fluxes is somewhat uncertain because of the smearing produced by artifical viscosity. After 2×10^7 years, the stellar mass and density interior to $r = 100$ pc are $\mathfrak{M}_* = 1.4 \times 10^9\ \mathfrak{M}_\odot$ and $n_* = 10^4$ protons cm^{-3}, both considerably larger than the mass and density of the gas.

b) Swallowing

When the calculations are carried beyond $t = 1.59 \times 10^7$ years, a curious quasi-cyclic process develops at the galactic center. The reason for this (with certain reservations explained below) is that not all of the gas that ultimately cools and flows toward the galactic center actually does so in the initial collapse. As the first cloud of cool gas falls toward the galactic center, the hot gas farther out rushes in toward the center, and the density at the innermost shells of this hot ($\sim 3 \times 10^6$ ° K) gas increases. When the density becomes high enough, a new shell of hot gas cools and falls toward the center. It appears from our calculations that spherical shells of thermally unsteady gas are deposited in nonlinear cycles with a mean period of 3.4×10^5 years. Each shell contains about $3.1 \times 10^5\ \mathfrak{M}_\odot$ of new cooled gas. We have followed about 15 of these cycles before $t = 2.08 \times 10^7$ years, where the calculation was terminated. The galactic nucleus is therefore unable to ingest all of the thermally unsteady gas during one smooth process, but rather swallows the gas over a period of many cycles.

Although the cyclic behavior of the unsteady flow is reasonable and has a simple explanation, we cannot be entirely certain of the reality of this effect because of the distortion of the solutions at the center due to artificial viscosity. During most of the swallowing phases, the region in which artificial viscosity exceeded the gas pressure roughly coincided with the extent of the region occupied by the cool gas ($T \approx 6000°$ K). More refined calculations (with a finer spatial zoning) will be necessary to determine the reality of the swallowing cycles. Perhaps the details of the cycles also depend on the particular distribution of stellar mass near the galactic center. It is nevertheless certain that a much greater quantity of cool ($T \approx 6000°$ K) gas will be produced than shown in Figures 5–7, and that the duration of the unsteady flow will extend considerably beyond 2×10^7 years.

c) Ultimate Fate of the Thermally Unsteady Gas

Any speculation concerning the final state of thermally unsteady galactic cores depends critically on whether the gas that falls toward the center remains ionized or becomes neutral. We therefore begin this section with a review of observational and theoretical arguments which suggest that the gas may be largely ionized.

For reasons outlined in § III*c* we have assumed that enough ionizing photons are present at the centers of giant ellipticals to maintain the temperature at a minimum value of 6000° K. If no ionizing photons were present, and if the gas became neutral following radiative cooling, then stars would certainly form in the neutral gas. Some of these stars would necessarily be massive and hot enough to ionize an observable amount of neutral gas. The Hβ flux resulting from ionization by only 10–20 main-sequence O stars in a dust-free H I cloud would be observable at 10 Mpc. Even more stringent is the observational restriction that very few if any main-sequence A- or F-type stars are present in ellipticals (Wood 1966; Spinrad and Taylor 1971). Since early-type stars would naturally form in an H I cloud, we conclude that such clouds do not exist at the centers of ellipticals. It also appears extremely unlikely that large amounts of neutral gas in the nuclei of ellipticals could be sufficiently agitated by stellar motions to discourage star formation.

To illustrate this last point, we consider a hypothetical H I cloud at the center of NGC 3379 and estimate the heating produced by gravitational agitation and bow shocks, which result from stellar motions. Since the stars pass through the cloud with highly supersonic velocities, we can ignore the gas temperature altogether in an estimate of the amount of heating by gravitation stirring. According to Dodd and McCrea (1952) the drag force on a star of mass $\mathfrak{M}_*$ and velocity v_* moving through a zero-temperature gas cloud of density ρ and size R_{cl} is

$$F_d = -2\pi\rho G^2\mathfrak{M}_*^2 v_*^{-2} \ln [1 + (R_{cl}v_*^2/G\mathfrak{M})^2] .$$

The work done by a single star crossing the cloud is $F_d R_{cl}$, and the total power delivered to the cloud by all the stars is $F_d R_{cl} n_* v_* \pi R_{cl}^2$ ergs s^{-1}, where n_* is the number density of stars. The volumetric heating rate is therefore $h_* = \frac{3}{4}F_d n_* v_*$ ergs cm^{-3} s^{-1}. With values appropriate to the center of NGC 3379, we find that $h_* \approx 5 \times 10^{-32} n$ ergs cm^{-3} s^{-1}, where $n = \rho/M$. This heating rate is much lower than the radiative-cooling rate for all cases of interest. Consider, for example, collisional excitation of the 63-μ line of O I by collisions with hydrogen atoms (see Penston 1970 for details and references). As long as the hydrogen density is less than $2.7 \times 10^5 T^{-1/3}$ cm^{-3}, collisional de-excitation can be ignored and the cooling rate due to 63-μ emission alone is much greater than h_* for $T \gtrsim 10^2$ ° K. The cooling rate may be modified if the H I cloud is optically thick to 63-μ radiation, but it would seem difficult to avoid the conclusion that heating by gravitational stirring as a result of the passage of stars through the nucleus is insufficient to inhibit star formation. Heating of a neutral cloud at the center of a giant elliptical can also be produced by bow shocks associated with the supersonic motion of stars (and their stellar winds), but again, this is likely to be small. In order to estimate this effect, we assume that each star creates a shock of cross section πl^2. The mass of gas shocked by one star traversing the H I cloud is $\rho\pi l^2 R_{cl}$, where R_{cl} is the cloud radius. If the mean stellar velocity is v_*, then $3n_*v_*/4R_{cl}$ stars pass through unit volume each second. If we assume that the shock is adiabatic (maximizing the heating rate), then $9\, v_*^2/32$ ergs are delivered to each gram of gas as it crosses a bow shock. When these factors are combined, we find that the bow-shock heating rate is $h_{sh} \lesssim 3 \times 10^{-55} n l^2$. For this to balance the O I 63-μ radiative cooling at $n = 10^2$ and $T = 10^2$ requires $l \gtrsim 10^{15}$ cm. It is very unlikely that stellar winds could provide enough dynamic pressure (ρu^2) to support a bow shock of this dimension. We conclude from these arguments that star formation could not be seriously suppressed in a cloud of neutral gas if it were placed at the center of a giant elliptical. Therefore, all or most of the gas at the centers of ellipticals must be ionized even if the density is quite high ($\sim 10^2$–10^3 cm^{-3}). This is also consistent with the observations discussed in § III*c*.

As a consequence of the ionization of gas at the centers of ellipticals, it is easy to show that gravitational collapse will favor formation of objects considerably more massive than ordinary stars. The Jeans radius for a self-gravitating gas is

$$R_J = \tfrac{1}{4}\lambda_J = \left(\frac{\pi kT}{16\mu MG\rho}\right)^{1/2} . \tag{15}$$

During the isothermal collapse of a cloud under its own gravity, the Jeans mass decreases as $\mathfrak{M}_J \sim \rho R_J^3 \sim \rho^{-1/2}$ so that smaller masses become unstable and fragmentation occurs. For gas which collects at the centers of giant ellipticals, however, the gravitational field is at first entirely due to the stars, i.e., $\rho_* \gg \rho$. The Jeans radius generalized to this situation is then

$$R_J \approx \left[\frac{\pi k T}{16 \mu M G(\rho + \rho_*)}\right]^{1/2} . \tag{16}$$

The essential point which may determine the ultimate fate of gas in the nuclei of ellipticals is that the Jeans radius depends only on the gas temperature and ρ_* as long as $\rho_* \gg \rho$. For NGC 3379 the observed stellar density is roughly constant within 100 pc, so the Jeans radius and the Jeans mass are also constant during the first phases of isothermal collapse. This means that fragmentation will not occur until the gas density becomes comparable to ρ_*. When $\rho > \rho_*$, however, fragmentation in the normal sense may not have time to occur if the infall is very fast (supersonic).

The gas masses which are originally Jeans unstable in galactic nuclei retain their coherence with the ensuing collapse. If $T \approx 10^4$ ° K and $\rho_* = 440\ \mathfrak{M}_\odot$ pc^{-3}, the Jeans radius according to equation (16) is $R_J \approx 4$ pc and the Jeans mass is $(\mathfrak{M}_J/\mathfrak{M}_\odot) \approx 7\ N_e$. Therefore, if $N_e \approx 10^2$ in a cloud of $10^7\ \mathfrak{M}_\odot$, 1.4×10^4 objects of $\sim 700\ \mathfrak{M}_\odot$ become unstable. Such massive objects may release an enormous supply of gravitational energy which may ultimately account for the nonthermal outbursts observed at the centers of giant ellipticals. Rotation must obviously play a major role in the development of these massive objects, and it will also help stablize even larger masses. In fact, rotating massive objects have often been suggested as a source of extragalactic nonthermal energy.

VIII. DISCUSSION AND SUMMARY

The absence of observable gas in most elliptical galaxies can be explained quite naturally by the heating produced in supernova blast waves. The observed rate of supernova outbursts is sufficient to heat the gas to temperatures in excess of T_*; a galactic wind is then produced which carries the gas into the intergalactic medium. Although we assume in this paper that gas ejected from the stars is heated to T_* by adiabatic shocks, this is not a critical assumption. A simple analysis (as in § VI) shows that galactic winds still result if $T_s \ll T_*$. Possibly these winds can be observed at infrared wavelengths if there is an admixture of dust (§ V). Tifft (1969, 1970) has in fact shown that the nuclear reddening observed in some ellipticals is consistent with internal interstellar reddening by dust. This conclusion may have to be slightly modified if the stars near the centers of these galaxies are also intrinsically redder, a conclusion suggested by the CN band observations of McClure (1969) and others. The OAO-II observations also indicate significant absorption by dust inside ellipticals.

Hot winds tend not to occur in galaxies which are supported primarily by rotation rather than by random motion of stars. However, in view of the remarks above, the smallness of the peculiar stellar velocities is unlikely to be the complete explanation for the absence of galactic winds in spirals and rapidly rotating ellipticals. Perhaps of more importance are the higher gravitational force and greater gas density in spirals. A flattened galaxy will provide a deeper potential well for the gas (at some fixed distance from the center) than a spherical system of the same mass. In rotating galaxies, radiative cooling to $T \ll T_s$ will tend to occur throughout the galactic plane. Also, the large amount of gas in spirals which has not yet formed into stars discourages the formation of hot galactic winds. The higher gas density reduces the velocity of supernova blast waves until radiative losses dominate. When this occurs, the kinetic energy of the supernova shells is less effective in heating the gas. The gas layer in our Galaxy is thinner than the stellar disk population. It is therefore possible that supernovae at large distances from the galactic plane could heat the tenuous local gas sufficiently to cause an outflow from the galaxy.

There is a strong possibility that nonthermal nuclear activity or optical emission lines can only occur in those ellipticals which have thermally unsteady winds. Although we have not computed final steady-state solutions of equations (1), (2), and (4) for galactic winds that undergo unsteady collapse, it is possible that gravitational instabilities will occur and smooth flow will break down before steady-state flow is reached. If the gas remains ionized during the early phases of this instability, as seems quite likely, then the smallest gravitationally stable clouds have masses considerably in excess of normal stellar masses. Perhaps the gravitational collapse of such massive objects, modified by rotation, will lead to nonthermal emission of various kinds. Although we believe that cosmic rays and magnetic fields probably play a minor role in leading to thermally unsteady flows in galactic nuclei, as the cool H II gas collapses in the stellar gravity field, it is likely that the influence of these nonthermal components on the observed radiation will increase. Thermally unsteady galactic winds are more likely to occur in massive galaxies which have rather high central star densities. Galaxies (or galactic nuclei) with masses less than about $10^8\ \mathfrak{M}_\odot$ will not be sufficiently massive to contain gas even at 10^4 ° K.

Although we believe that only those ellipticals with thermally unsteady winds can produce observable optical emission lines, the exact connection with the observations also depends on other factors. The free-fall time of the thermally unsteady gas is only $t_{ff} \approx (4\pi\rho_* G)^{-1/2} = 3 \times 10^5$ years, very much shorter than the Hubble time. The duration of freely falling thermally unsteady gas in a galaxy may be extended by the repeated "swallowings" discussed in § VII*b*, but it is likely that a radio source will be produced in the nucleus before many nonlinear cycles have taken place. In this event, the gas in the galactic nucleus which surrounds a new source of nonthermal activity may be agitated to rather high velocities. Since the gas clouds observed in galactic nuclei often have velocities in excess of the escape velocity, other energy sources in addition to gravity are required.

Recently Heeschen (1970) has shown that there is a correlation between the nonthermal radio spectra of elliptical galaxies and the radio structure of the sources. He finds that galaxies with normal power-law spectra ($F_\nu \sim \nu^{-0.9}$) have radio sizes of $\gtrsim 60$ arc seconds. Conversely, those sources with complex (optically thick) turnovers at centimeter wavelengths usually have the bulk of their energy emitted within 1″.5. Disney and Cromwell (1971) have shown that optical emission is common in those galaxies with complex radiofrequency spectra, but those galaxies with normal radio spectra apparently have little or no optical emission. Presumably, all these galaxies have thermally unsteady winds, and the creation of a new center of nonthermal activity (with apparent dimensions $\lesssim$ 1″.5) strongly agitates the ambient gas. However, the gravitational binding energy of the thermally unsteady gas is only 10^{51}–10^{54} ergs, so if a radio source were ejected or expanded rapidly at the center, the whole central region of the galaxy would be outgassed and intense shocks preceding the expanding radio source would heat the gas to very high temperatures where it would become unobservable. This may explain why optical emission is absent in the larger sources studied by Heeschen.

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