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GAS-RICH DWARFS AND ACCRETION PHENOMENA IN EARLY-TYPE GALAXIES

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ABSTRACT

Recent constraints on intergalactic H I clouds suggest that allowable accretion rates by several luminous early-type galaxies are too low to account for their observed H I content. We have therefore developed an alternative model, wherein gas-rich dwarf galaxies are accreted into galactic halos. This process is significant in groups of galaxies only when a sufficiently high density of gas-rich dwarfs (~ 30 Mpc^{-3}) is present. The dwarf galaxy gas content plays a crucial role in enabling the galaxy to be trapped in the halo by interaction with a galactic wind or extensive gaseous corona. Gas stripping occurs, resulting in the formation of dwarf spheroidal systems that populate the outer halos of massive galaxies and in the injection of a system of clouds into the halo. If the clouds are initially confined by the pressure of the ambient halo medium, dissipation and continuing infall enable the clouds to accrete into the central regions of galaxies before becoming gravitationally unstable and presumably forming stars. Consequences of this scenario include the production of a radial abundance gradient and infall of adequate amounts of neutral gas to account for the observations of H I in early-type galaxies. This gas accretion rate is also sufficient to feed active nuclei and radio sources. An important cosmological implication of our model is that, because the characteristic time scale of a gas-rich dwarf galaxy to be accreted and lose its gas is comparable to a Hubble time, there may have been a far more extensive primordial distribution of such systems at earlier epochs. This implies that accretion rates were greatly enhanced at relatively recent epochs ($z \gtrsim 0.5$) and could account both for the rapid cosmological evolution inferred for radio galaxies and quasars, and for the observed frequency of occurrence of quasar absorption-line systems.

Subject headings: galaxies: formation — galaxies: intergalactic medium — galaxies: structure

I. INTRODUCTION

It has recently seemed attractive to appeal to the existence of intergalactic H I clouds of mass greater than, or order of, $10^8\ \mathfrak{M}_\odot$ in order to account for observations of gas in early-type galaxies. Three prominent examples may be cited. In the case of NGC 4278, the inclination of the extended gas disk indicates that the gas must probably have originated from a source whose angular momentum distribution is different from that of the galaxy. Accretion of an intergalactic gas cloud has provided a possible explanation (Bottinelli and Gouguenheim 1977; Gallagher *et al.* 1977; and Knapp, Kerr, and Williams 1978). A similar event may also have occurred in the case of NGC 1052 (Knapp, Faber, and Gallagher 1978; Fosbury *et al.* 1978; Reif, Mebold, and Goss 1978). Another interesting case is the Spindle galaxy NGC 2685, observations of which (Shane 1977) have been interpreted to show evidence for the recent accretion of intergalactic H I. Oort (1967) has long maintained that the extensive high-velocity H I provides evidence for the infall of intergalactic matter into our Galaxy. Gunn (1977) has proposed that accretion of intergalactic neutral gas clouds can produce extended H I disks in elliptical galaxies, the interaction of which with nuclear winds will lead to substantial gas infall into the central nucleus. Direct evidence for the existence of intergalactic H I clouds in the Sculptor group has been given by Mathewson, Cleary, and Murray (1975).

Unfortunately, there is very little actual evidence for the widespread occurrence of intergalactic H I clouds of mass greater than $10^7\ \mathfrak{M}_\odot$. The results of Mathewson *et al.* have been disputed by Haynes and Roberts (1979), who in a more extensive survey find that the high-velocity H I toward Sculptor may be part of the Magellanic Stream and therefore be relatively local. More seriously, Lo and Sargent (1979) have extensively searched for intergalactic gas clouds in nearby groups. For example, *none* were found of mass greater than $\sim 4 \times 10^7\ \mathfrak{M}_\odot$ in the M81 group, and the ensuing limits effectively demonstrate that the average rate of accretion of intergalactic H I in clouds of more than $\sim 10^8\ \mathfrak{M}_\odot$ is less than $\sim 0.1\ \mathfrak{M}_\odot$ yr^{-1} per luminous early-type galaxy. A more serious constraint applies if more massive clouds are required: for example, less than $0.01\ \mathfrak{M}_\odot$ yr^{-1} per galaxy seems to be available in clouds of mass greater than $\sim 5 \times 10^8\,\mathfrak{M}_\odot$. Accretion of rates of this order are probably too low to account for the observed H I in some ellipticals (Knapp, Kerr, and Williams 1978).

This appears to leave us in a serious quandary. Either a substantial amount of gas is stored at a great distance from the nucleus, or else gas evidently has to be supplied to the gas-rich elliptical galaxies at a rate $\sim$(0.1–1) $\mathfrak{M}_\odot$ yr^{-1}. The difference in symmetry axes of the gas disk and the stellar distribution in NGC 4278 may be difficult to reconcile with storage in a disk at great distances because of the relatively short precession time (Gunn 1977). Moreover, observational constraints are only marginally consistent with the latter possibility of intergalactic (as opposed to circumgalactic) gas accretion. In fact, in the case of a galaxy such as NGC 4636 with a hydrogen mass of $8 \times 10^8\,\mathfrak{M}_\odot$, the limit on accretion is at least an order of magnitude too low for the gas to have originated from a single encounter with an intergalactic H I cloud.

There is a possible resolution to this problem that we shall explore in the present paper. Several nearby low-surface-brightness gas-rich dwarf galaxies ($M_{pg} \approx -10$) with ratios of total hydrogen mass to luminosity $\mathfrak{M}_{\rm H\,I}/L \approx$ 3–4 have recently been discovered in the Local Group (Longmore *et al.* 1978; Cesarsky *et al.* 1977), in the Sculptor group (Laustsen *et al.* 1977; Cesarsky, Falgarone, and Lequeux 1977), and in the Phoenix region (Canterna and Flower 1977). A number of such systems have also been found by Lo and Sargent (1979), who have estimated that there may be a substantial number of faint hydrogen-rich dwarf galaxies in groups, the corresponding space density with M_{pg} between -9 and -11 being $\sim$30 Mpc^{-3}. This estimate is based on an extrapolation of the luminosity function of galaxies and on a preliminary examination of IIIa-J plates taken of two of the surveyed groups. Lo and Sargent (1979) have extended the Fisher-Tully relation (Fisher and Tully 1975) for $\mathfrak{M}_{\rm H\,I}/L$ to $M_{pg} = -10$ (the ratio of hydrogen mass to luminosity increasing to 1–10 in this magnitude range). This result extrapolates the indication of an increase in $\mathfrak{M}_{\rm H\,I}/L$ toward somewhat lower absolute magnitudes than was found by Fisher and Tully (1975) and suggests that the faintest dwarfs with $M_{pg} = -10$ may predominantly consist of H I. Such an inference is based on at least two independent lines of argument: (*a*) $\mathfrak{M}/L$ for the stellar component is likely to lie in the range 1–3, and (*b*) the indicative dynamical masses determined from H I line widths (Fisher and Tully 1975) yield an upper limit on the total mass of the system.

We shall henceforth assume that there is a population of faint ($M_{pg} \gtrsim -10$) hydrogen-rich galaxies in groups with the parameters indicated by these admittedly preliminary observations. In addition, there is likely to be a more numerous population of dwarf spheroidal galaxies. Van den Bergh (1968) estimates from the Palomar sky survey that there may be as many as 500 dwarf spheroidal systems Mpc^{-3} if the Local Group is uniformly filled with systems similar to Draco and Sculptor. Such low-surface-brightness systems are not easily detectable at a characteristic distance greater than several hundred kpc, where individual stars can no longer be easily resolved. Gas-rich dwarfs are more easily detectable than dwarf spheroidals of similar M_{pg}, and a frequency of such systems in the Local Group as high as that inferred for the M81 group would not have been detected. Nor would 21 cm surveys of high-velocity gas (Hulsbosch 1978) have easily detected these gas-rich dwarfs; conversely, we might speculate that deep IIIa-J plates of the regions where the high-velocity gas column density peaks should reveal the presence of faint stellar counterparts. The existence of a population of faint hydrogen-rich dwarf galaxies in the Local Group would appear to provide an important test of the Lo-Sargent inference of their existence in other similar nearby groups. Our motivation in assuming that such a population exists in groups is that intergalactic H I clouds probably fail to fuel gas-rich ellipticals at the required rate. Observers who have postulated the frequent occurrence of massive intergalactic H I clouds to explain their observations must seek an alternative source of gas. Thus we shall investigate whether accretion of dwarf hydrogen-rich galaxies can provide a significant mass input into giant ellipticals. A related concept, due to Einasto and his colleagues (Einasto 1978), is that of hypergalaxies, which are depicted as consisting of giant spiral galaxies surrounded by dwarf elliptical companions. We shall argue that hypergalaxies may be an end product of the interaction that we shall describe.

It is necessary to provide a mechanism by which trapping of dwarf systems can occur; otherwise, little, if any, accretion will result. One can readily demonstrate that dynamical friction is inadequate (by a factor proportional to the mass ratio of the accreted and accreting galaxies). We shall therefore explore the role of galactic winds and extensive gaseous halos or coronae in enhancing the accretion rate.

One interesting clue to the possible effects of accretion is the tentative correlation between hydrogen content and nuclear activity in several ellipticals (Ekers 1978*a*). It seems likely that infall may lead to nuclear activity. In turn, one might expect gas accretion and ensuing star formation to lead to the development of galactic winds and to the formation of extended gaseous coronae. Indeed, strong theoretical reasons have been given by Faber and Gallagher (1976) that winds provide the dominant gas removal mechanism in early-type galaxies.

Here we shall analyze the combined effects of cloud accretion and galactic winds and coronae. Many interesting phenomena result from their interaction. In particular, we shall study the accretion by early-type galaxies of gas-rich dwarfs. Accretion will be found to fuel the wind, thereby regulating the accretion flow and yielding a time-dependent model for star formation, enrichment, and nuclear activity. In § II we discuss the permissible parameter range for intergalactic gas clouds and galaxy groups, and in § III we discuss the frequency of gas-rich dwarfs and give an argument that may account for their large ratios of gas mass to luminosity. The interaction of gas-rich dwarfs with galactic winds and gaseous coronae is described in § IV. The evolution of the gas clouds formed by gas stripping is discussed in § V. The orbits of infalling clouds are evaluated in § VI, and some effects associated with star formation are considered in § VII. In § VIII the cosmological evolution of the gas infall rate is discussed and applied to radio sources and quasar evolution, quasar absorption-line systems, and the formation of stellar halos. Section IX summarizes our results.

II. PARAMETERS OF INTERGALACTIC CLOUDS IN GROUPS OF GALAXIES

We first consider the various theoretical constraints on intergalactic clouds in order to examine the significance of the upper limits on their spatial frequency obtained by Lo and Sargent (1979).

If substantial amounts of gas are present in galaxy groups, as is suggested, for example, by Ekers (1978*b*), it is likely that the gas will be clumpy. At the density of 10^{-3} cm^{-3} required by the thermal confinement model for these radio-trail galaxies, the cooling time for a primordial mixture of H and He is only $\sim 4 \times 10^9$ yr if $T \approx 10^6$ K. The temperature cannot greatly exceed this value if the intergalactic gas is to be confined within the groups. If the intergalactic gas is enriched, as is suggested by cluster observations, the cooling time is reduced by up to two orders of magnitude. Once the gas can cool, it is likely to form cooler clouds (Silk 1970). Studies of quasar absorption lines have revealed the presence of considerable numbers of low-excitation components along the line of sight. It seems likely that some of these are associated with intergalactic H I or H II clouds in galaxy groups (Bahcall 1978). The identification of gas-rich, low-mass galaxies could conceivably provide evidence that galaxy formation from intergalactic clouds may have occurred within significantly less than a Hubble time (Sargent and Searle 1970). Finally, the evidence we have previously cited for ongoing accretion may require the presence of intergalactic clouds.

Strong constraints can be imposed on the parameters of intergalactic clouds in galaxy groups. Cowie and McKee (1976) have discussed these restrictions in the context of a widespread hot intergalactic plasma of density $\sim 10^{-6}$ cm^{-3} and temperature $T \gtrsim 10^8$ K. Here we specifically restrict our attention to galaxy groups because the H I rich dwarf galaxies for which intergalactic clouds are plausible progenitors have hitherto been found only outside of rich clusters, and generally in groups. We therefore consider the allowable parameter range for clouds contained in a local intergalactic medium confined to groups.

The cloud masses are restricted to a relatively narrow range. On the one hand, the Jeans criterion must be satisfied in order to avoid premature collapse. On the other hand, a lower limit on cloud mass is obtained in order to avoid evaporation by the hot confining medium. For nominal parameters, we shall adopt an intergalactic medium with $T_{i,6} \equiv T_i/10^6$ K, density $n_{i,4} \equiv n_i/10^{-4}$ cm^{-3}, and pressure $n_i T_i$ K cm^{-3}. This implies a total mass in hot gas of $\sim 10^{13}\, \mathfrak{M}_\odot$ and a soft X-ray luminosity of $\sim 10^{43}$ ergs s^{-1}. These parameters are consistent with current observational constraints. For the clouds, we adopt a temperature $T_{c,4} \equiv T_c/10^4$ K, radius $R_1 \equiv R/10$ kpc, and density $n_{c,3} \equiv n_c/10^{-3}$ cm^{-3}. If the clouds are thermally supported and confined by the pressure of the intergalactic medium, the Jeans mass is

$$\mathfrak{M}_J = 10^8 T_{c,4}{}^2 n_{i,4}{}^{-1/2} T_{i,6}{}^{-1/2}\, \mathfrak{M}_\odot\,. \tag{1}$$

In order for the clouds not to evaporate within a Hubble time, their mass must exceed

$$\mathfrak{M}_{EV} = 5 \times 10^6 T_{i,6}{}^{13/4} T_{c,4}{}^{1/2} n_{i,4}{}^{-1/2}\, \mathfrak{M}_\odot \tag{2}$$

and their radius must exceed $R_{EV} = 1.7 T_{i,6}{}^{3/4} T_{c,4}{}^{1/2} n_{i,4}{}^{-1/2}$ kpc. The clouds are constrained to lie within the range defined by these two critical masses. In particular, their column densities must satisfy

$$5 \times 10^{17} T_{i,6}{}^{3/4} n_{i,4}{}^{-1/2} T_{c,4}{}^{-1/2}\ \mathrm{cm}^{-2} < n_c R_c < 3 \times 10^{20} n_{i,4}{}^{1/2} T_{i,6}{}^{1/2}\ \mathrm{cm}^{-2}\,. \tag{3}$$

Let us now consider the survival of these clouds. If no heavy elements are present, the cooling time is $\sim 10^{18} T_{i,6}{}^{1/2} n_{i,4}{}^{-1}$ s. Evidently, the clouds cannot cool within a Hubble time unless $n_i T_i > 100$ K cm^{-3}. However, cloud-cloud collisions may provide an important energy sink. Because of the long cooling time, cloud-cloud collisions at relative velocities characteristic of galaxy groups will result in cloud disruption. If we require the mean cloud-cloud collision time to exceed a Hubble time, the number density of clouds must satisfy

$$N_c \lesssim 10^4 R_1{}^{-2} v_{cc,2}{}^{-1}\ \mathrm{Mpc}^{-3}\,, \tag{4}$$

where $v_{cc,2} = v_{cc,2}/10^2$ km s^{-1} is the average relative velocity for cloud-cloud encounters. For clouds at the evaporation limit to survive collisions, their number density cannot exceed $3 \times 10^5 T_{i,6}{}^{-3/2} T_{c,4}{}^{-1} v_{cc,2}{}^{-1} n_{i,4}$ Mpc^{-3}.

If we specify the average gas density $\langle n \rangle$ of intergalactic gas in the group, including both clouds and intercloud material, we can derive the gas pressure:

$$\frac{p}{k} = \frac{\langle n \rangle T_i}{1 + f T_i/T_c}\,, \tag{5}$$

where f is the fractional volume occupied by the clouds. If there are comparable amounts of matter in clouds and in the intercloud material, we have $f T_i/T_c \approx 1$, yielding $p/k \approx 100 T_{i,6}(\langle n \rangle/10^{-4}\ \mathrm{cm}^{-3})$ cm^{-3} K.

There can be exceedingly few intergalactic clouds with masses in excess of $5 \times 10^8\, \mathfrak{M}_\odot$ present in groups. If the pressure is of order $nT = 100$ K cm^{-3}, as inferred from the observations of radio trails in groups, we expect that (eq. [1]) the more massive clouds will be unstable and presumably form dwarf galaxies. This suggests that the minimum masses of dwarf galaxies may be $\sim(10^7-10^8)\,\mathfrak{M}_\odot$, with the smaller clouds stabilized by the intragroup gas. Since the lower mass limit against evaporation is $\sim 10^7\,\mathfrak{M}_\odot$, there appears to be only a narrow mass range where intergalactic H I clouds can exist. This suggests that the existence of a substantial mass fraction at the present epoch in groups in the form of intergalactic H I clouds must be considered implausible. Our argument has assumed that the clouds are randomly distributed within the group. It is more likely that the clouds are preferentially formed in or near the potential wells of the largest galaxies and their subsequent evolution will be similar to that described for individually accreted clouds or gas-rich dwarfs.

III. GAS-RICH DWARFS

Statistical arguments suggest that the blue compact dwarf galaxies ($M_B \approx -15$) first described by Sargent and Searle (1970) are relatively old. An estimate of the minimum mean lifetime τ_d can be derived from the upper limit of 0.02 Mpc^{-3} on H I clouds of mass $\sim 5 \times 10^8\,\mathfrak{M}_\odot$ in at least one group (Lo and Sargent 1979). Since the space density of these compact dwarfs is $n_d \approx 0.008\ \mathrm{Mpc}^{-3}$ (Sargent and Searle 1970), we infer that

$$\tau_d \gtrsim \frac{1}{2}\frac{1}{H_0}\frac{n_d}{n_c} \approx 4 \times 10^9\ \mathrm{yr}\,, \tag{6}$$

where the Hubble constant H_0 is taken to be 50 km s^{-1} Mpc^{-1}.

Such a great age is difficult to understand because of the rise in $\mathfrak{M}_{\mathrm{H\,I}}/L$ with decreasing luminosity found for dwarfs. In fact, the faint, gas-rich dwarfs ($M_{\mathrm{pg}} \approx -10$) may predominantly consist of H I, and yet presumably are as long-lived as the more luminous dwarfs. The discussion of § II shows that H I clouds in groups of sufficient mass to form the dwarfs would be unstable over a collapse time scale considerably less than τ_d. How then are we to understand the apparent longevity of dwarfs?

We propose the following model. In these low-mass systems it is likely that the shallow potential well will inhibit star formation. Energy input from evolving stars will restrict the amount of gas that can collapse and form stars. Larson (1974) has demonstrated that this effect is important for the less massive galaxies, resulting in a diminution of systematic enrichment as stellar evolution proceeds. The mechanism by which this inhibition of gas collapse is believed to occur utilizes a galactic wind. The wind is fueled by supernova remnants in more massive systems. In the low-mass systems of interest here, we can extend this argument to consider systems where (since the supernova rate is too low) we may appeal to mass loss or UV radiation from luminous evolving stars to provide the energy input which drives the wind.

There is a critical value of the gas density in order for a wind to be driven (Mathews and Baker 1971) based on the requirement that the characteristic time for radiation losses at the sonic point be longer than the corresponding flow time. This condition can be expressed as a luminosity-dependent constraint on the ratio of total gas mass to stellar luminosity.

Consider a uniform gas cloud of mass $\mathfrak{M}_{\mathrm{H\,I}}$, radius r_c, and density n_c that is homogeneously interspersed with a stellar population of mean surface brightness Σ. Let the effective kinetic temperature of the gas cloud be given by the virial relation

$$T = \frac{G(\mathfrak{M}_{\mathrm{H\,I}} + \mathfrak{M}_*)m_{\mathrm{H}}}{kR_c} = \frac{(4\pi\Sigma)^{1/2}Gm_{\mathrm{H}}}{k}\frac{\mathfrak{M}_*}{L}\left(1 + \frac{\mathfrak{M}_{\mathrm{H\,I}}}{\mathfrak{M}_*}\right)L^{1/2}\,.$$

Assume that the energy input to the gas is $\epsilon L/\mathfrak{M}_{\mathrm{H\,I}}$ ergs g^{-1} s^{-1} and that the radiative cooling rate is $n_c\Lambda(T)/m_{\mathrm{H}}$ ergs g^{-1} s^{-1}. Then no net outflow (or wind) will occur if the cooling is sufficiently efficient, or $n_c\Lambda(T_c)/m_{\mathrm{H}} > \epsilon L/\mathfrak{M}_{\mathrm{H\,I}}$. Using the relation

$$n_c = \frac{\mathfrak{M}_{\mathrm{H\,I}}}{(4/3)\pi m_{\mathrm{H}} R_c^{\,3}} = \frac{3(4\pi)^{1/2}\mathfrak{M}_{\mathrm{H\,I}}\Sigma^{3/2}}{m_{\mathrm{H}}L^{3/2}}\,,$$

we can write down a minimum value of $\mathfrak{M}_{\mathrm{H\,I}}/L$ as a function of L above which there is no net outflow:

$$\frac{\mathfrak{M}_{\mathrm{H\,I}}}{L} > \frac{\epsilon^{1/2}m_{\mathrm{H}}L^{1/4}}{3^{1/2}(4\pi)^{1/4}\Lambda^{1/2}\Sigma^{3/4}} = 0.35\left(\frac{\epsilon}{\Lambda_{22}}\right)^{1/2}\frac{L_7^{\,1/4}}{\Sigma_1^{\,3/4}}\,, \tag{7}$$

where $\Lambda_{22} \equiv \Lambda/10^{-22}$ ergs cm^3 s^{-1}, $L_7 = L/10^7\,L_\odot$, and $\Sigma_1 \equiv \Sigma/1\,L_\odot\,\mathrm{pc}^{-2}$ (equivalent to a visual surface brightness of 26.5 mag arcsec^{-2}). Suppose that Σ is reasonably uniform over a range of L for these dwarf galaxies, as is inferred for somewhat more luminous and higher-surface-brightness galaxies (Kormendy 1977). It now follows that as L is decreased, since $T = 2 \times 10^4(\Sigma_1 L_7)^{1/2}$ K, Λ will rapidly decrease because of the strong diminution in the rate of Lα cooling below 2×10^4 K. Even if the gas is enriched, one may expect Λ to decline to less than $\sim 10^{-25}$ ergs cm^3 s^{-1} below 10^4 K. This extreme sensitivity to Λ over a narrow temperature range that also corresponds to the binding energy of the smallest gas-rich dwarfs suggests that the gas content of these galaxies could be self-regulated by means of the energy input from the winds. In order to yield strong cooling and subsequent throttling or quenching of the winds, $\mathfrak{M}_{\rm HI}/L$ should rise dramatically for $L \lesssim 10^7 L_\odot$. Such an effect is indicated by the observations of Lo and Sargent (1979), and perhaps suggests that the gas content of gas-rich dwarfs is regulated by winds.

More luminous systems ($L \gtrsim 10^7 L_\odot$) are likely to have insufficient gas to quench a wind if $\mathfrak{M}_{\rm HI}/L \lesssim 1$. Here we emphasize that in *gas-rich* systems there is a natural explanation for the apparent sharp rise in $\mathfrak{M}_{\rm HI}/L$ due to the quenching of galactic winds. One might expect oscillations in the rate of star formation (and in L) to occur in the vicinity of the critical value of $\mathfrak{M}_{\rm HI}/L$. Note that such oscillations are self-stabilizing: as L rises, the wind strength increases and the available gas supply is reduced. This effect may account for the extreme blueness of these systems by leading to a series of bursts of star formation. The characteristic duration between bursts is determined by the hydrodynamic flow time, or $\sim 10^8$ yr. The high space density inferred for gas-rich dwarfs (~ 30 Mpc^{-3}) in groups indicates that their interaction with larger galaxies may play a role in galactic evolution. We shall explore this possibility below.

IV. INTERACTION WITH GALACTIC WINDS AND GASEOUS HALOS

If the actual density of gas-rich dwarfs is comparable to the estimated value $N_{c,30} \equiv N_c/30$ Mpc$^{-3} \approx 1$, a galaxy of mass $\mathfrak{M}_g = 10^{12}\,\mathfrak{M}_\odot$ will undergo frequent interactions with these objects. The effective cross section of the galaxy is increased by the effect of gravitational focusing. If the distance of minimum approach is the characteristic halo radius, we assume that the gas-rich dwarfs are trapped by interaction with the halo medium. The effective accretion radius of such a massive galaxy is the geometric mean of the halo radius and the classical accretion radius (Hoyle and Lyttleton 1939), and is given by $R_A = 2G\mathfrak{M}_g v_{gc}^{-1}\sigma_g^{-1} = 300\mathfrak{M}_{12} v_{gc,2}^{-1}\sigma_{g,2.5}^{-1}$ kpc, where $\mathfrak{M}_{12} = \mathfrak{M}_g/10^{12}\,\mathfrak{M}_\odot$, the effective galaxy-cloud collision velocity $v_{gc} = (v_g^2 + v_c^2)^{1/2}$ is the velocity of a galaxy moving at v_g relative to a cloud of mass $\mathfrak{M}_c$ moving at v_c and $v_{gc,2} = v_{gc}/100$ km s^{-1}, and the internal velocity dispersion of the galaxy $\sigma_{g,2.5} = \sigma_g/250$ km s^{-1}. The ensuing accretion rate of these discrete clouds is therefore $\dot{\mathfrak{M}}_g = 4\pi R_A^2 N_c v_{gc} = 0.07\mathfrak{M}_{c,7} N_{c,30}\mathfrak{M}_{12}\sigma_{g,2.5}^{-2} v_{gc,2}^{-1}\,\mathfrak{M}_\odot$ yr^{-1}, and $\mathfrak{M}_{c,7} = \mathfrak{M}_c/10^7\,\mathfrak{M}_\odot$. The characteristic time scale for a galaxy to collide with a cloud is now expressed as

$$t_{g,c} \approx \frac{\mathfrak{M}_c}{\dot{\mathfrak{M}}_g} = 2 \times 10^8 v_{gc,2}\sigma_{g,2.5}{}^2\mathfrak{M}_{12}{}^{-1}N_{c,30}{}^{-1}\ \text{yr}\,. \tag{8}$$

We now describe a mechanism for trapping these gas-rich dwarfs at large radial distances from the dominant galaxy. Galactic winds (or gaseous halos) can exert sufficient drag on the gaseous component to trap the dwarfs if $\mathfrak{M}_{\rm HI}/\mathfrak{M}_*$ is sufficiently large. The fundamental constraint that the stopping time t_s be less than the galaxy halo crossing time $t_{g,\rm cross}$ is

$$t_s/t_{g,\rm cross} = \frac{4}{3}\left(\frac{n_c R_c}{n_a R_h}\right) \ll 1\,, \tag{9}$$

if the average ambient density is n_a and R_h is the halo radius.

For $\mathfrak{M}_{\rm HI}/\mathfrak{M}_* > 1$, there is a critical condition on the entry velocity v_{gc} of the dwarf such that, if

$$\frac{G\mathfrak{M}_{\rm HI}}{R_c v_{gc}{}^2} \lesssim \frac{3}{2}\left(\frac{n_a}{n_c}\right), \tag{10}$$

the gas and stars will separate. The required density contrast for this to hold is

$$\frac{n_a}{n_c} > \frac{2}{3}\left(\frac{v_{c,s}}{v_{gc}}\right)^2,$$

where we define the internal velocity dispersion of the cloud to be $v_{c,s} = (G\mathfrak{M}_{\rm HI}/R_c)^{1/2}$. In general, we infer that for separation to occur, $n_a/n_c \gtrsim 0.1$. The radius at which separation occurs, obtained by equating the drag force with that exerted by the gas on the stars, is

$$\frac{v_{c,s}}{v_w}\left(\frac{3\dot{\mathfrak{M}}v_w R_c}{8\pi G\rho_c\mathfrak{M}_{\rm HI}}\right)^{1/2} = 89\,\frac{v_{c,s}}{v_w}\left(\frac{v_{w,2}\dot{\mathfrak{M}}_1 R_1}{n_{c,3}\mathfrak{M}_{c,7}}\right)^{1/2}\ \text{kpc}\,, \tag{11}$$

where the mass flux in the wind $\dot{\mathfrak{M}}_1 \equiv \dot{\mathfrak{M}}_w/1\,\mathfrak{M}_\odot$ yr^{-1}, the wind velocity $v_{w,2} \equiv v_w/100$ km s^{-1}, and the cloud density $n_{c,3} \equiv n_c/10^{-3}$ cm^{-3}.

We shall henceforth consider separately the subsequent fate of the infalling gas clouds and the stellar component. Most of our discussion will concentrate on gas cloud infall; we shall return to discuss the fate of the stellar component in § VIII.

Consider the nature of the ensuing gaseous interaction. We shall assume that galaxies have winds, or at least possess gaseous halos. Two cases may be distinguished.

a) Large Clouds

Suppose first that the size of the accreted gas cloud is greater than the radius of the galaxy. This is expected to be a rare occurrence since massive H I clouds are relatively infrequent. Capture will occur if the time scale for the galaxy to cross the cloud exceeds the free-fall time for the cloud to fall into the galaxy. The ratio of the relative encounter velocity to the average internal velocity dispersion of the galaxy must be significantly less than the ratio of the galaxy and cloud radii. In this case, a shell will develop around the galaxy, and be supported against gravity by the pressure of the wind or interstellar medium. The scale height H of the gas shell is $R_0{}^2 v_{a,s}{}^2/G\mathfrak{M}_g$, where $\mathfrak{M}_g$ is the galaxy mass, R_0 is the shell radius at which pressure balance occurs, and $v_{a,s}$ is the sound velocity in the (assumed isothermal) shell. Such a system will be Rayleigh-Taylor unstable, with the hot underdense wind supporting the cooler denser shell against gravity. It is expected that the nonlinear development of this instability leads to the infall of smaller clouds from the surface of the shell. The ensuing fate of these small clouds is analogous to that for similar clouds that have been directly accreted.

b) Small Clouds

Consider next the interaction of a galaxy with clouds of radius smaller than the galactic radius. We have in mind both direct capture of such clouds as well as infall resulting from the instability described above. The condition for infall of a cloud toward the central regions of the galaxy is that the gravitational force on a cloud, mass $\mathfrak{M}_c$, exceed the net outward pressure exerted by the ambient medium. We shall now assume either that an extensive gaseous halo is present or that there is a wind generated predominantly in the central core. In the former case, the condition for infall is that the halo density be less than the cloud density (namely, the buoyancy criterion). To consider a wind, we adopt a simple wind model with constant velocity v_w on the supersonic branch of the wind solution outside the sonic point (Mathews and Baker 1971). If the mass flux in the wind is $\dot{\mathfrak{M}}_w$ and the wind density is ρ_w, we have by mass conservation $\dot{\mathfrak{M}}_w = 4\pi\rho_w r^2 v_w$. The infall condition yields a critical condition on the cloud column density:

$$\rho_c R_c > \langle\rho_c R_c\rangle_{\rm crit} \equiv \frac{3}{16\pi}\left(\frac{\dot{\mathfrak{M}}_w v_w}{\mathfrak{M}_g G}\right) = 10^{-7} v_{w,2} \mathfrak{M}_{12}{}^{-1} \dot{\mathfrak{M}}_1 \text{ g cm}^{-2} . \quad (12)$$

Clouds satisfying this condition will be accreted by the galaxy. For the shell considered previously, a similar condition holds and the internal radius of the shell (where the wind pressure supports the shell against gravity) is $R_0{}^2 = \mathfrak{M}_{\rm shell}/4\pi\langle\rho H\rangle_{\rm crit}$, where $\mathfrak{M}_{\rm shell}$ is the mass of the shell and H is the shell scale height. (For typical numbers $\mathfrak{M}_{\rm shell} = 10^8\,\mathfrak{M}_\odot$ and $R_0 = 30$ kpc.) Henceforth, it is immaterial whether we consider the case of a wind or an isothermal gaseous halo. In either case, it is a reasonable approximation away from the nuclear regions to assume that the ambient density varies as r^{-2}.

Clouds whose column density exceeds the critical column density are initially cooler and denser than their surroundings. To illustrate their likely fate, let us assume that the infalling clouds evolve adiabatically and remain in pressure balance with the ambient gas. Combining appropriate expressions for the adiabatic and pressure balance conditions yields an expression for the cloud density as a function of galactocentric radius,

$$\frac{\rho_c}{\rho_{c0}} = \left(\frac{r_{c0}}{r}\right)^{6/5},$$

where

$$r_{c0} = \left(\frac{T_w \dot{\mathfrak{M}}_w}{T_{c0}\rho_{c0} v_w}\right)^{1/2} .$$

This expression is given for a wind. To transform the result to the case of a gaseous halo, we simply replace $\dot{\mathfrak{M}}_w$ by $4\pi\rho_{a0} r_0{}^2 v_w$, where the halo density ρ_a satisfies $\rho_a = \rho_{a0}(r_0/r)^2$ and the subscript zero denotes a reference point (evaluated at the halo core). The density contrast decreases until a critical radius is reached at which buoyancy forces will halt any further infall. This radius, at which $\rho = \rho_a$, is given by

$$r_b = \left(\frac{\dot{\mathfrak{M}}_w T_w}{4\pi v_w}\right)^{1/2} \left(\frac{T_{c0}}{\rho_{c0}{}^{2/2}}\right)^{3/4} . \quad (13)$$

Note that ρ_0 and T_0 enter the definition of r_b in a form which is constant along an adiabat. However, r_b is proportional to $T_{c0}{}^{3/4}$, indicating that initially colder clouds will fall further in toward the galactic center.

Because buoyancy forces halt their infall, the clouds are subject to the Rayleigh-Taylor instability. Note that, since the drag forces on the cloud are appreciable, the Rayleigh-Taylor instability will act on the cloud throughout the infall. However, the characteristic growth time scale for the Rayleigh-Taylor instability to develop is

$$\tau_{\text{R-T}} = \left(\frac{16\pi}{3}\frac{n_c}{n_a}\right)^{1/2}\frac{R_c}{v_c}\,. \tag{14}$$

This can be longer than the time to reach the buoyancy limit, or the radius at which all clouds will be buoyant and eventually go Rayleigh-Taylor unstable. Infalling clouds will initially tend to overshoot their equilibrium positions, resulting in a situation where lighter cloud material is being decelerated by denser ambient gas. This leads to mixing with the ambient gas and consequent enrichment in heavy elements.

V. CLOUD EVOLUTION

There are two parameters that are crucial to the evolution of an infalling cloud. These are the initial level of ionization and the degree of heavy-element enrichment. We first consider two initial cloud configurations corresponding to the extreme cases of neutral or ionized clouds, and then discuss how enrichment enhances cloud cooling and instability.

a) Neutral Clouds

We consider first the fate of primordial neutral gas clouds. Compressional heating raises the temperature to $\sim 10^4$ K. If Lα cooling is significant, the cloud remains effectively isothermal. It is necessary to consider the trapping of Lα photons to evaluate the actual net cooling rate. The Lα photon diffusion time can be expressed as $\sim \tau R_c/c$, where the optical depth at line center $\tau = 5.9 \times 10^{-14}(T_{c,4})^{-1/2} n_c R_c$ (Adams 1971). This estimate assumes that the cloud has no systematic differential motion; if such a motion parametrized by velocity $v_{c,s}$ is present, the condition for photons to escape is $t_{\text{diff}} < R_c/v_{c,s}$. This is equivalent to $\tau < c/v_{c,s}$, or column density less than $6 \times 10^{17}(T_{c,4})^{1/2}(10 \text{ km s}^{-1}/v_{c,s})$ cm^{-2}. Evidently, if the cloud column density is less than $\sim 6 \times 10^{17}$ cm^{-2}, photons escape by frequency redistribution into the line wings in a time scale t_{diff}. If the column density exceeds this critical value, Lα photons escape by Doppler shift out of the line core in a time scale $R_c/v_{c,s}$. We conclude that neutral clouds will evidently remain isothermal. It is apparent from the cloud trajectories in the (T, n)-diagram (Fig. 1) that H I clouds of mass less than $\sim 10^8$ $\mathfrak{M}_\odot$ will be able to undergo considerable compression before

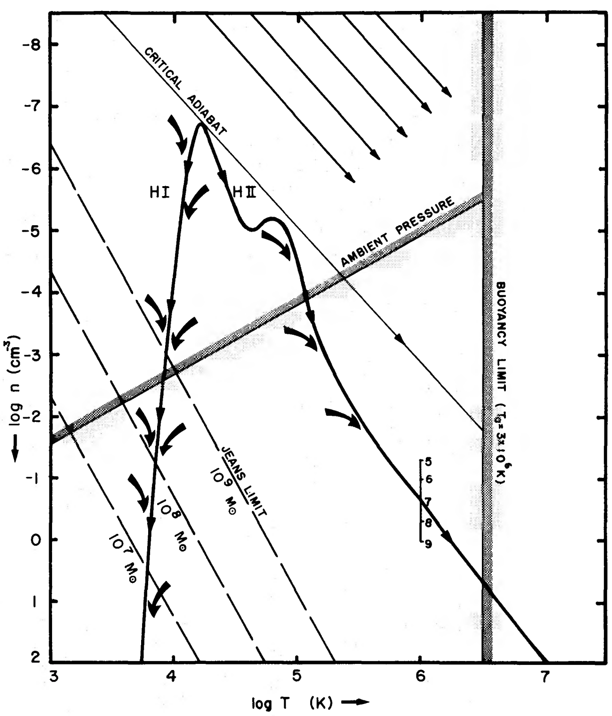

FIG. 1.—The (T, n)-diagram. Arrows denote cloud evolution trajectories as described in text. Heavy curve is cooling constraint for clouds of 10^7 $\mathfrak{M}_\odot$ in pressure equilibrium; vertical scale shows displacement for 10^5, 10^6, 10^8, and 10^9 $\mathfrak{M}_\odot$ clouds. Also shown are lines of constant Jeans mass (marked "Jeans limit"), corresponding to 10^7, 10^8, and 10^9 $\mathfrak{M}_\odot$, the critical adiabat, such that clouds lying above this are unable to cool, and hatched lines corresponding to typical values of the ambient halo pressure and temperature (this latter yields the buoyancy limit for pressure-confined clouds).

becoming Jeans unstable. The range of initial cloud parameters is constrained by the ambient pressure of the intragroup medium. Neutral clouds will be at $n_c > 10^{-3}$ cm^{-3} if $n_i T_i \gtrsim 10$ cm^{-3} K. Infall into the central regions (~ 10 kpc) from the accretion radius (~ 1 Mpc) requires compression by a factor 10^4 and can be achieved only by H I clouds with masses less than $\sim 10^6$–10^7 $\mathfrak{M}_\odot$. Because the H I clouds remain colder and denser than the ambient medium, they never become buoyant.

Gravitational instability occurs at sufficiently high density, and consequent star formation will effectively halt the infall, unless dynamical friction provides an important drag force in the high-stellar-density galactic nucleus. Since $\mathfrak{M}_J \propto r$ for isothermal infall, it is actually the lowest-mass clouds that fuel the central regions. *Large numbers of low-mass clouds therefore provide a much more important accretion source than a single massive cloud.*

b) Ionized Clouds

If the infalling clouds are initially ionized to the extent that the Lα (and to some extent the He II) cooling barriers can be overcome, they will subsequently tend to heat adiabatically. To demonstrate this, we introduce in Figure 1 adiabatic trajectories for clouds of fixed mass. A cloud will be unable to cool if its adiabat lies above a cooling curve defined by the condition that the cloud cool locally within a sound-wave crossing time $t_{c,\mathrm{cross}}$. For typical parameters, we find that

$$t_{c,\mathrm{cross}} = 10^8 \mathfrak{M}_{c,7}^{1/3}(10^5\ \mathrm{K}/T_c)^{1/2}\ \mathrm{yr}\ , \tag{15}$$

whereas the cooling time at 10^5 K is

$$t_{\mathrm{cool}}(10^5\ \mathrm{K}) = 10^{7.5} n_{c,3}^{-1}\ \mathrm{yr}\ . \tag{16}$$

It is clear that for a wide variety of initial conditions in the (T, n)-plane that initially satisfy $t_{\mathrm{cool}} > t_{\mathrm{cross}}$ (for example, at pressures determined by the ambient medium), the infalling clouds will heat adiabatically. Note that only massive ionized clouds ($\gtrsim 10^{11}$ $\mathfrak{M}_\odot$) can be gravitationally unstable and collapse directly (cf. Silk 1977). Less massive clouds, if initially below the Jeans mass, will remain so as they evolve adiabatically ($T \propto n^{2/3}$). Such infalling clouds are halted only by buoyancy, after having fallen in a considerable distance. Buoyancy results in Rayleigh-Taylor instability and mixing (§ IV*b*). The resulting rapid cooling will lower the Jeans mass and initiate cloud collapse.

c) Enrichment

We have described in § IV how primordial gas clouds can become enriched by means of a specific mechanism for mixing with ambient material. A quantitative estimate of the cooling due to the enrichment process can be obtained from the energy equation $(3/2)nk\dot{T} = -n^2\Lambda(T, n, Z)$, where Z is the mass fraction of heavy elements. If we schematically write the cooling function $\Lambda(T) = A_\Lambda Z(t) T^\mu$ and describe the increase in heavy-element abundance by $\dot{Z} = \dot{Z}_0(t_0/t)$, then the cloud temperature is given by

$$\frac{T_c}{T_0} = \left\{1 - (1 - \mu)\frac{2n_c}{3k}\frac{A_\Lambda Z_0}{T_0^{1-\mu}}[t_0 - t + t\ln(t/t_0)]\right\}^{1/(1-\mu)}, \tag{17}$$

with initial conditions $T_c = T_0$ and $Z = Z_0$ at some initial time $t = t_0$. The initial value of Z_0 is chosen such that the bremsstrahlung and heavy-element cooling rates are comparable. This result indicates that the temperature drops rapidly, over a time scale comparable to the initial cooling time at t_0. The rapid cooling guarantees that gravitational instability occurs, and fragmentation and star formation will immediately follow enrichment.

VI. CLOUD KINEMATICS

Infalling clouds will conserve angular momentum until the frictional drag exerted by the ambient medium becomes appreciable. Subsequently, as the drag force increases with decreasing galactocentric distance, the infall rate will increase and the clouds will spiral inward.

While it seems that low-angular-momentum material will fall in on radial orbits, the bulk of the accreting material with higher specific angular momentum may well have a more flattened distribution. Enhanced accretion rates may actually be a consequence of the interaction of a wind or gaseous halo with infalling clouds, and gas will accumulate in the central regions. To demonstrate this, consider first the spiraling-in time t_f of a cloud of radius R_c due to friction with the ambient medium. This can be expressed relative to the angular frequency $\Omega = (G\mathfrak{M}_g/R_g^3)^{1/2}$ as

$$\Omega t_f = \frac{\mathfrak{M}_c}{\dot{\mathfrak{M}}_w}\frac{v_w}{r}\left(\frac{r}{R_c}\right)^2 = 6\left(\frac{N_c}{10^{19}\ \mathrm{cm}^{-2}}\right)\left(\frac{v_w}{300\ \mathrm{km\ s}^{-1}}\right)\left(\frac{r}{100\ \mathrm{kpc}}\right)\left(\frac{1\ \mathfrak{M}_\odot\ \mathrm{yr}^{-1}}{\dot{\mathfrak{M}}_w}\right), \tag{18}$$

where N_c is the cloud column density (which must also satisfy the critical infall condition [12]). Combining equations (12) and (18) yields $\Omega t_f \geq v_w^2 r/2G\mathfrak{M}_g \approx 1$, and we infer that clouds are in essentially Keplerian orbits.

Under this assumption, the radial velocity of a cloud is given by (recall that R_c decreases with decreasing r)

$$\frac{v_r}{r} = -2\pi\left(\frac{R_c}{r}\right)^2\left(\frac{\mathfrak{M}_w}{\mathfrak{M}_c}\right)\left(\frac{G\mathfrak{M}_g}{v_w{}^2 r}\right)^{1/2}$$
$$= 5 \times 10^{-10}\left(\frac{N_c}{10^{19}\ \mathrm{cm}^{-2}}\right)\left(\frac{\dot{\mathfrak{M}}_w}{1\ \mathfrak{M}_\odot\ \mathrm{yr}^{-1}}\right)\left(\frac{100\ \mathrm{kpc}}{r}\right)^2\left(\frac{G\mathfrak{M}_g}{v_w{}^2 r}\right)^{1/2}\ \mathrm{yr}^{-1}$$
$$\equiv \left(\frac{-H}{1+\Theta}\right) r^{-\Theta-1}\,, \tag{19}$$

where $\Theta = 0.7$ if the clouds compress adiabatically during infall and

$$H = 2\pi(1+\Theta)\left(\frac{G\mathfrak{M}_g}{v_w{}^2}\right)^{1/2}\left(\frac{\mathfrak{M}_w}{\mathfrak{M}_c}\right) R_{c0}{}^2 r_{c0}{}^{-8/10}\,. \tag{20}$$

The wind-driven accretion mechanism therefore allows low-surface-density clouds ($N_c \approx 10^{19}$–10^{20} cm^{-2}) to fall into the nuclear region.

However, cloud-cloud collisions may prevent infall if they are sufficiently disruptive. The condition for a collision to lead to disruption rather than coalescence can be expressed as the requirement that the cooling time of the shocked cloud gas exceed the shock crossing time. We express the cooling time over the temperature range 2×10^5–3×10^7 K (Raymond, Cox, and Smith 1976) as

$$t_c = 10^{7.85} T_6{}^{1.75}(10^{-3}\ \mathrm{cm}^{-3}/n_c)\ \mathrm{yr}\,, \tag{21}$$

where the immediate post-shock temperature for clouds colliding at relative velocity v_c is $T_6 = 10^{0.1}(v_c/300\ \mathrm{km\ s}^{-1})^2$. The shock crossing time is $10^{7.5}(R_c/10\ \mathrm{kpc})(300\ \mathrm{km\ s}^{-1}/v_c)$ yr, and we infer that cloud disruption occurs outside a critical radius

$$\left(\frac{R_c}{v_c t_c}\right)^{3/4} r_h = 0.2\lambda^{3/4}\left(\frac{N_c}{10^{19}\ \mathrm{cm}^{-2}}\right)^{3/4}\left(\frac{300\ \mathrm{km\ s}^{-1}}{v_c}\right)^{\beta} r_h\,, \tag{22}$$

where r_h is the halo radius and λ equals the heavy-element abundance relative to that of the Sun. Equation (22) is valid for heavy-element cooling with $\beta = 27/8$; if λ is sufficiently small ($\lesssim 0.01$ at 10^6 K), the thermal bremsstrahlung cooling rate must be used, with $\beta = 3/2$. It is clear that if there is sufficient enrichment ($\lambda \approx 1$), as expected if the infalling gas clouds originate in dwarf galaxies, coalescence can occur within 20 kpc, where we would therefore expect H I to accumulate.

VII. STAR FORMATION

We have hitherto assumed the existence of a wind or gaseous halo independently of the accretion rate. Steady mass loss from evolving stars in an elliptical galaxy at a rate $\sim 0.15\ \mathfrak{M}_\odot\ \mathrm{yr}^{-1}\ (10^{10}\ L_\odot)^{-1}$, coupled with a standard supernova rate (Tammann 1974), is generally considered sufficient to result in a wind. If the mass injection rate is appreciably increased, radiative cooling will set in and steady outflow will cease (Mathews and Baker 1971). Continuing gas infall will lead to massive star formation. Thus, if the cloud accretion rate is sufficiently great, the resulting supernovae will drive and amplify the wind. However, as the wind strength increases, infall will be inhibited. In this manner, accretion can regulate a galactic wind and produce bursts of star formation. We now describe a simple schematic model that illustrates how the feedback can occur.

a) Bursts of Star Formation

Let us assume that a mass fraction k_* of the infalling clouds forms stars with a universal mass function $N_*(m)$ given by

$$\frac{dN_*}{dm} = A_* m^{-1-x} \quad (m_l \le m \le m_u)\,. \tag{23}$$

The slope x and the upper and lower mass cutoffs will be taken to be independent of time. Only the integrated mass will be allowed to vary. The total mass in stars is

$$\mathfrak{M}_* = \frac{A_*}{1-x} m_l{}^{1-x}\,, \qquad \text{if } x > 1\,.$$

If all stars more massive than $m_{\rm SN}$ are assumed to become supernovae with typical evolution time $t_{\rm SN}$, the supernova rate is given by

$$\left\langle \frac{1}{t_{\rm SN}} \right\rangle = \int_{m_{\rm SN}}^{m_u} \frac{dN_*}{dm} \frac{1}{t_{\rm SN}(m)} dm$$

$$= \frac{A_*}{t_{\rm SN}} m_{\rm SN}^{-x} . \tag{24}$$

The mass of infalling clouds is given by

$$k_*^{-1}\mathfrak{M}_* = \frac{A_*}{k_*} \frac{m_l^{1-x}}{1-x} . \tag{25}$$

The cloud mass infall rate can also be expressed in terms of the flux of clouds accreted by the galaxy across a fixed boundary (taken to be the accretion radius), namely,

$$\pi R_A^2 v_{g,c} \int_{\mathfrak{M}_{\rm crit}}^{\infty} \frac{dN(\mathfrak{M})}{d\mathfrak{M}} \mathfrak{M} d\mathfrak{M} , \tag{26}$$

where $N(\mathfrak{M})$ is the cloud mass spectrum and the lower limit $\mathfrak{M}_{\rm crit}$ is derived from the infall criterion (12). This condition is actually a lower limit on the column density for infall to occur, expressed in terms of the wind strength. We shall write this as a functional dependence of $\dot{\mathfrak{M}}$ on cloud mass, with the minimum cloud mass that can accrete written

$$\mathfrak{M}_{\rm crit} = A \dot{\mathfrak{M}}_w^{\beta} . \tag{27}$$

For uniform-density clouds accreted across a fixed boundary (at R_A), we explicitly have

$$\mathfrak{M}_{\rm crit} = \frac{9\pi}{16} \left(\frac{v_w}{G\mathfrak{M}_g} \right)^3 \rho_c^{-2} \dot{\mathfrak{M}}_w^3 . \tag{28}$$

Now, the wind mass-loss rate can be expressed in terms of the supernova rate as

$$\dot{\mathfrak{M}}_w = \frac{E_{\rm SN}}{v_w^2} \langle t_{\rm SN}^{-1} \rangle . \tag{29}$$

If we adopt a simple power-law cloud mass spectrum of the form

$$\frac{dN(\mathfrak{M})}{d\mathfrak{M}} = D\mathfrak{M}^{-\alpha} , \tag{30}$$

the resulting equation for $\mathfrak{M}_*$ is

$$\frac{\dot{\mathfrak{M}}_*}{m_u} = \frac{1}{\tau_*} \left[1 - \left(\frac{\mathfrak{M}_*}{m_u \Psi} \right)^{\beta(2-\alpha)} \right] , \tag{31}$$

where $\tau_*^{-1} = k_* \pi R_A^2 v_{gc} (\bar{\rho}/m_u)$ and $\bar{\rho}$ is the mean mass density in clouds. In the general case (eq. [27]), the parameter Ψ is given by

$$\Psi^{-1} = \left[\left(\frac{A}{m_u} \right)^{1/\beta} m_u \frac{1-x}{t_{\rm SN}} \frac{E_{\rm SN}}{m_l v_w^2} \left(\frac{m_{\rm SN}}{m_l} \right)^{-x} \right] ,$$

and in the uniform cloud case (eq. [28]) by

$$\Psi^{-1} = \left(\frac{9\pi}{16} \right)^{1/3} \left[(1-x) \left(\frac{m_u}{\rho_c} \right)^{2/3} \frac{1}{G\mathfrak{M}_g} \frac{v_w}{t_{\rm SN}} \frac{E_{\rm SN}}{m_l v_w^2} \left(\frac{m_{\rm SN}}{m_l} \right)^{-x} \right] .$$

Inspection of these equations shows that, since a physical solution requires $\dot{\mathfrak{M}}_* > 0$, we must have $\alpha < 2$. In this case, star formation rises linearly with time and is abruptly quenched after a time scale greater than $\sim \tau_*$ when $\mathfrak{M}_* \approx \Psi m_u$.

Our solution shows that star formation occurs in bursts, with the star formation effectively throttled by the rise in wind strength. Typical time scales for renewal of a burst of star formation are determined by the gas replenishment time of $\sim 10^9$ yr.

It should be noted that this feedback occurs at the accretion surface. As clouds fall in, their column densities rise because of compression and the value of the critical mass (above which infall can occur) decreases with decreasing distance from the galactic nucleus. For clouds that are compressed adiabatically, this critical mass for infall corresponding to a cloud of mass $\mathfrak{M}_0$ can be expressed as a function of galactic radius in the form

$$\frac{\mathfrak{M}_{\rm crit}}{\mathfrak{M}_0} = \pi^{-3}\left(\frac{\dot{\mathfrak{M}}_w R_{c0}}{v_w \mathfrak{M}_0}\right)^{9/5}\left(\frac{r}{R_A}\right)^{12/5}\left(\frac{v_w}{v_g}\right)^{6}\left(\frac{R_{c0}}{R_A}\right)^{3/5}\left(\frac{3T_{c0}}{T_w}\right)^{6/5}, \quad (32)$$

where R_{c0} is the initial cloud radius. Because of the strong radial dependence, the rise in wind strength fueled by the feedback and the compression as clouds fall in compete in determining whether $\mathfrak{M}_{\rm crit}$ actually could rise and therefore eventually halt the infall. We have assumed that conditions at the accretion radius determine the mass of clouds that can overcome the wind and that infall dominates until clouds fall into the buoyancy radius. This can be justified if the time scale for $\dot{\mathfrak{M}}$ to change is larger than (or comparable to) the cloud infall time, a condition which should be well satisfied.

b) *Radial Dependence of the Star Formation Rate*

According to the buoyancy constraints, ionized clouds of different initial entropy will first become unstable at different radii. Neutral clouds will become Jeans unstable at radii that depend on their initial density. If, as before, we assume that cloud instability results in star formation, we can, in principle, infer the radial dependence of the star formation rate in terms of parameters that describe the initial spectrum of clouds.

If the cloud mass spectrum is such that most of the mass in H I clouds is at the low end of the spectrum, it is clear that the star formation rate in general will increase with decreasing galactocentric distance.

VIII. EVOLUTIONARY CONSIDERATIONS

It is reasonable to infer that there might have been a considerably greater density of gas-rich dwarfs and clouds in the past. In what follows we shall not differentiate between gas-rich dwarfs and clouds at earlier epochs. The characteristic lifetime of a cloud is determined by the time for a cloud to collide with another cloud or with a galaxy. At the present time, the latter process dominates, and the corresponding time scale

$$\tau_{cg} = (N_g \sigma_g v_{cg})^{-1} = 3 \times 10^{10} N_{g,1}\left(\frac{R_A}{0.3\ \text{Mpc}}\right)^{-2} v_{cg,2}^{-1}\ \text{yr}, \quad (33)$$

where the density of massive galaxies in groups $N_{g,1} = N_g/(1\ \text{Mpc})^{-3}$. Note that the galaxy-cloud collision time relevant for estimates of the accretion rate is a factor of ~ 100 shorter than the cloud-galaxy collision time; see equation (8). The approximate coincidence between this characteristic cloud lifetime and a Hubble time ($\sim 2 \times 10^{10}$ yr) suggests that we are merely witnessing the remnant of a more extensive primordial distribution. If this is indeed the case, accretion by galaxies of H I contained in gas-rich dwarfs now proceeding at the modest rate of $\sim 0.1\ \mathfrak{M}_\odot\ \text{yr}^{-1}$ will have been considerably greater in the relatively recent past. This can have profound implications for galactic evolution, as we shall now demonstrate.

The evolution of the space density of clouds can be described by the equation

$$\frac{dN_c}{dt} = -N_c N_g \sigma_g v_{cg} - N_c^2 \sigma_c v_{cc},$$

which has two limiting solutions: (*a*) cloud-cloud collisions dominate,

$$N_c(t) = \frac{N_c(0)}{1 + t/\tau_{cc}}, \quad (34)$$

with characteristic time $\tau_{cc}^{-1} = N_c(0)\sigma_c v_{cc}$; and (*b*) cloud-galaxy collisions dominate,

$$N_c(t) = N_c(0) \exp(-t/\tau_{cg}). \quad (35)$$

If $\tau_{cc} \ll \tau_{cg}$, cloud-cloud collisions will dominate at early times and give a t^{-1} decay in the cloud density. After a time τ_{cg}, an exponential decay sets in as cloud-galaxy collisions prevail. If $\tau_{cc} \gtrsim \tau_{cg}$, the exponential decay over a time scale τ_{cg} will dominate. To decide which of these limits is applicable, we write

$$\frac{\tau_{cc}}{\tau_{cg}} = \frac{N_g \sigma_g}{N_c(0)\sigma_c} = N_{g,1}\left(\frac{10^3\ \text{Mpc}^{-3}}{N_c(0)}\right)\left(\frac{R_A}{0.1\ \text{Mpc}}\right)^2\left(\frac{10\ \text{kpc}}{R_c}\right)^2.$$

Our calculation of R_A (§ IV) may underestimate R_A in the limit of sufficiently high cloud densities at early times, when cloud-cloud collisions will occur and the classical accretion radius (Hoyle and Lyttelton 1939) (equal to $R_A \sigma_g / v_{gc}$) should be used. The actual value of the ratio τ_{cc}/τ_{cg} is unimportant in consideration of the observable effects of cloud collisions, which are dominated by the exponential decay. The remarkable point about potential applications of collisions to galactic evolution is that τ_{cg} *is comparable to a Hubble time for moderately massive galaxies*. This means that *strong evolutionary effects can occur at low redshift*.

We now consider several specific applications.

a) *Radio Source Evolution*

The recent survey by de Ruiter (1978) has shown that the space density evolution of radio sources has an extremely steep rise proportional to $(1 + z)^{11}$ in the vicinity of $z \approx 0.5$ and subsequently flattens out to give the $(1 + z)^6$ evolution law obtained by previous workers. This strong evolutionary effect begins at a lookback time of $\tau_R \approx 4 \times 10^9$ yr. We identify τ_{cg} with $t_0 - \tau_R$, where t_0 is the total time elapsed since galaxy and group formation. Thus galaxies which satisfy

$$1 = \delta \int_0^\infty N_g(\mathfrak{M}_g)\mathfrak{M}_g^2 v_{cg}^{-3} d\mathfrak{M} = \left(\frac{\tau_{cg}}{10^{10}\ \mathrm{yr}}\right)^{-1}(1\ \mathrm{Mpc}^{-3})(5 \times 10^{11}\ \mathfrak{M}_\odot)^2(100\ \mathrm{km\ s}^{-1})^{-3}\,, \quad (36)$$

where $N_g(\mathfrak{M}_g)$ is the differential mass spectrum (per unit volume) of galaxies in groups and δ is the density enhancement of the group relative to the background galaxy distribution, will exhibit an exponential decrease in cloud accretion rate over the required time scale. If we assert that cloud accretion fuels powerful radio galaxies, we then have a natural, although qualitative, explanation of the observed evolution. The implied space density of clouds is a factor $\sim 10^3$ greater at $z \gtrsim 1$ than at the present epoch.

A natural consequence of this interpretation is that radio source evolution will be accompanied by bursts of star formation and the associated galaxies will be significantly bluer than non–radio galaxies. Our model of cloud accretion works effectively for galaxies in groups, where the low velocity dispersion yields large accretion rates. A similar effect may also occur in rich clusters for central galaxies or for very massive systems. However, since such galaxies are relatively rare compared to galaxies in groups, we expect that the bulk of the observed radio evolution will occur in galaxies outside rich clusters.

b) *Quasar Evolution*

Observations of steep-radio-spectrum quasars indicate that their space density evolution can be fitted with an evolution law of the form $\exp[-10(t - t_0)/t_0]$, where t is cosmic epoch (Schmidt 1975). It is intriguing to note that, according to equation (36), relatively dense groups will accrete clouds with a characteristic time scale $\tau_{cg} \lesssim 10^9$ yr. If $\tau_{cg} \approx 10^9$ yr,

$$N_c(t) = N_c(0) \exp(-10t/t_0)\,,$$

where the initial cloud density $N_c(0) = 10^4 N_c(t_0)$ and $N_c(t_0)$ is the present density of clouds in the vicinity of these massive galaxies. In general, the range in δ is considerable; Bhavsar (1978) finds that δ is proportional to f_g^{-2}, where f_g is the fraction of galaxies in groups and $0.01 \leq f_g \leq 1$. Hence $\tau_{cg} \propto (\delta f_g)^{-1}$ and has a range of 10^8–10^{10} yr. It is natural to expect that the greatest cloud density will occur in the densest groups, which we identify with the phenomenon of quasar evolution. Occurrence of quasars in groups, as predicted by our model for their density evolution, is generally consistent with their apparent anticorrelation with rich clusters at $z < 0.5$ (Bahcall 1978).

c) *Quasar Absorption Lines*

The large extent of the accretion sphere around massive galaxies guarantees that at $z \gtrsim 0.5$ these regions will contain many clouds of mass $\sim(10^6$–$10^7)\ \mathfrak{M}_\odot$. At the present epoch they contain of order 1–10 clouds, whereas at earlier epochs our radio source evolution model implies an increase of $\sim 10^3$. The cross section for the line of sight to a quasar at a redshift $z \gtrsim 1$ to intersect a cloud in the halo of an intervening galaxy is accordingly enhanced relative to that for intersection with the radius $R_{g,\mathrm{opt}}$ inferred from the optical image by a factor $\sim 10^4(R_c/R_{g,\mathrm{opt}}) \gtrsim 100$. This suffices to account for the considerable number of absorption-line systems found in quasars which could be produced in intervening systems. Specific observations that support this general picture of absorption-line formation in clouds within the extended halos of intervening galaxies are of narrow 21 cm (Haschick and Burke 1975) and of broad Ca II (Boksenberg and Sargent 1978) absorption lines in the spectrum of 3C 232. The quasar is at a projected distance of ~ 60 kpc from the center of the galaxy NGC 3067 and at the same redshift. This indicates the possible presence of at least two components in the intervening galaxy that may be identified with a discrete gas cloud embedded in a more dilute and extensive ionized medium. It has been previously noted that extended gaseous halos provide a possible site for the formation of quasar absorption lines (Bahcall and Spitzer 1969). However, both the observations cited of 3C 232 and of excited fine-structure Si II absorption at several redshifts in PKS 0237−23 (Boroson *et al.* 1978) indicate that the absorbing medium must be nonuniform if the intervening halo hypothesis is adopted.

d) *Stellar Halos*

Separation of gas and stars in gas-rich dwarf galaxies can leave the stellar component with positive energy. This provides an attractive mechanism for populating the galactic halo with stars. A more frequent occurrence would be the expansion of the stellar component to a more relaxed spheroidal configuration. This could conceivably result in the formation of globular clusters and dwarf spheroidal galaxies. It is tempting to speculate that a high frequency of gas-rich dwarfs at an earlier epoch would have evolved by interaction with a massive parent galaxy into a disk and halo system similar to that of the Milky Way (cf. Fall and Rees 1977, who discuss the tidal and evaporative disruption of primordial galactic subclusters).

IX. CONCLUSIONS

The lack of any significant evidence for the existence of massive intergalactic H I clouds has motivated us to consider the interaction of gas-rich dwarf galaxies with early-type galaxies. Such gas-rich dwarfs may be relatively numerous in groups, with space density $\sim 30\ \mathrm{Mpc}^{-3}$, and a massive galaxy will accrete such a dwarf every $\sim 10^8$ yr. The dwarf galaxy gas content now plays a crucial role. We point out that $\mathfrak{M}_{\mathrm{H\,I}}/L$ must significantly increase for dwarf systems if they are to retain any gas at all. This effect enables us to understand the extremely high gas content of low-luminosity dwarf irregular systems. Interaction with a galactic wind or an extensive gaseous corona results in orbital drag and subsequent trapping of gas-rich dwarfs in the halo. Gas stripping will occur, forming dwarf spheroidal systems that remain in the halo and gas clouds that dissipate and suffer further infall. The implied association of dwarf spheroidal companions around massive parent galaxies results in the production of systems that can resemble the hypergalaxies of Einasto. We assume that the stripping process results in the injection into the outermost halo of a spectrum of clouds confined by external thermal pressure with a wide range of masses. General considerations of cooling and stability enable us to distinguish two cases. Initially neutral clouds in pressure equilibrium with the ambient halo medium will fall in isothermally until they become gravitationally unstable. Initially ionized clouds which are colder than the ambient medium will fall in adiabatically until the ambient density is comparable to the internal cloud density, when the clouds achieve buoyancy and become Rayleigh-Taylor unstable. In this case, mixing of enriched material will accelerate cloud cooling and tend to initiate gravitational instability. Thus we generally conclude that clouds will fall in a substantial radial distance before becoming unstable and forming stars. This provides a natural mechanism for developing a radial abundance gradient and for injecting matter into the central regions of galaxies.

This accretion model provides an adequate source of H I to account for observations of neutral gas in early-type galaxies. For two of the nine ellipticals in which H I has been detected, it seems that the neutral gas has a significantly different angular momentum distribution from that of the parent galaxy. One possibility is that a single massive ($\sim 10^9\ \mathfrak{M}_\odot$) intergalactic H I cloud has recently been accreted. Observational searches for such clouds constrain this to be a relatively rare event. A more frequent occurrence will be accretion of several (perhaps ≥ 10) gas-rich dwarfs within the precession time of a massive gaseous disk ($\sim 10^9$ yr). In this latter case, one can readily demonstrate that if the parent galaxy has a net velocity relative to the incident dwarfs, the accreted material will retain a net angular momentum characteristic of the rms value of the specific angular momentum evaluated at the various radii where trapping of the individual clouds first occurred. If, at early epochs, cloud-cloud collisions are sufficiently dissipative, then disk formation will be a natural consequence. Warping of spiral disks may represent another manifestation of this accretion phenomenon.

Furthermore, there are important cosmological implications. We may just be witnessing the destruction of the tail end of a more extensive primordial distribution of gas-rich dwarf galaxies and intergalactic clouds. This speculation rests on the fact that the characteristic time scale for a gas-rich dwarf galaxy to be accreted and lose its gas is at present of order 10^{10} yr. We argue that this surprising coincidence can be readily understood by supposing that there was a far higher density of such systems at earlier epochs. Such a model implies that accretion rates were greatly enhanced at relatively recent cosmological epochs ($z \gtrsim 0.5$). We propose that gas accretion fuels radio sources, and we associate this enhanced accretion rate with the rapid cosmological evolution inferred for radio galaxies and quasars. The increased frequency of gas clouds in galactic halos at moderate redshift can also account for the observed frequency of quasar absorption-line systems.

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REFERENCES

Adams, T. F. 1971, *Astr. Ap.*, **12**, 280.
Bahcall, J. N. 1978, *Phys. Scripta*, **17**, 229.
Bahcall, J. N., and Spitzer, L. 1969, *Ap. J.* (*Letters*), **156**, L63.
Bhavsar, S. P. 1978, *Ap. J.*, **222**, 412.
Boksenberg, A., and Sargent, W. C. W. 1978, *Ap. J.*, **220**, 42.
Boroson, T., Sargent, W. L. W., Boksenberg, A., and Carswell, R. F. 1978, *Ap. J.*, **220**, 772.
Bottinelli, L., and Gouguenheim, L. 1977, *Astr. Ap.*, **54**, 641.
Canterna, R., and Flower, P. J. 1977, *Ap. J.* (*Letters*), **212**, L57.
Cesarsky, D. A., Falgarone, E. G., and Lequeux, J. 1977, *Astr. Ap.*, **59**, L5.
Cesarsky, D. A., Laustsen, S., Lequeux, J., Schuster, H.-E., and West, R. M. 1977, *Astr. Ap.*, **61**, L31.
Cowie, L. L., and McKee, C. F. 1976, *Ap. J.* (*Letters*), **209**, L105.
de Ruiter, H. 1978, Ph.D. thesis, University of Leiden.
Einasto, J. 1978, in *IAU Symposium No. 79, Large Scale Structure of the Universe*, ed. M. S. Longair and J. Einasto (Dordrecht: Reidel), p. 51.
Ekers, R. D. 1978*a*, in *IAU Symposium No. 77, Structure and Properties of Nearby Galaxies*, ed. E. M. Berkhuijsen and R. Wielebinski (Dordrecht: Reidel), p. 221.
———. 1978*b*, in *IAU Symposium No. 79, Large Scale Structure of the Universe*, ed. M. S. Longair and J. Einasto (Dordrecht: Reidel), p. 149.

Faber, S., and Gallagher, J. S. 1976, *Ap. J.*, **204**, 365.
Fall, S. M., and Rees, M. J. 1977, *M.N.R.A.S.*, **181**, 37P.
Fisher, J. R., and Tully, R. B. 1975, *Astr. Ap.*, **44**, 151.
Fosbury, R. A. E., Mebold, U., Goss, W. M., and Dopita, M. A. 1978, *M.N.R.A.S.*, **183**, 549.
Gallagher, J. S., Knapp, G. R., Faber, S. M., and Balick, B. 1977, *Ap. J.*, **215**, 463.
Gunn, J. 1977, preprint.
Haschick, A. D., and Burke, B. F. 1975, *Ap. J.* (*Letters*), **200**, L137.
Haynes, M. P., and Roberts, M. S. 1979, *Ap. J.*, **227**, 767.
Hoyle, F., and Lyttleton, R. A. 1939, *Proc. Cambridge Phil. Soc.*, **35**, 405.
Hulsbosch, A. 1978, *Astr. Ap.*, **66**, L5.
Knapp, G. R., Faber, S. M., and Gallagher, J. S. 1978, *A.J.*, **83**, 11.
Knapp, G. R., Kerr, F. J., and Williams, B. A. 1978, *Ap. J.*, **222**, 800.
Kormendy, J. K. 1977, *Ap. J.*, **217**, 406.
Larson, R. B. 1974, *M.N.R.A.S.*, **166**, 585.
Laustsen, S., Richter, W., van der Lans, J., West, R. M., and Wilson, R. N. 1977, *Astr. Ap.*, **54**, 639.
Lo, K. Y., and Sargent, W. L. W. 1979, *Ap. J.*, **227**, 756.
Longmore, A. J., Hawarden, T. G., Webster, B. L., Goss, W. M., and Mebold, U. 1978, *M.N.R.A.S.*, **183**, 97P.
Mathews, W. G., and Baker, J. C. 1971, *Ap. J.*, **170**, 241.
Mathewson, D. S., Cleary, M. N., and Murray, J. D. 1975, *Ap. J.* (*Letters*), **195**, L97.
Oort, J. H. 1967, in *IAU Symposium No. 31, Radio Astronomy and the Galactic System*, ed. H. van Woerden (London: Academic Press), p. 279.
Raymond, J. C., Cox, D. P., and Smith, B. W. 1976, *Ap. J.*, **204**, 290.
Reif, K., Mebold, U., and Goss, W. M. 1978, *Astr. Ap.*, **67**, L1.
Sargent, W. L. W., and Searle, L. 1970, *Ap. J.* (*Letters*), **162**, L155.
Schmidt, M. 1975, in *Galaxies and the Universe*, ed. A. Sandage, M. Sandage, and J. Kristian (Chicago: University of Chicago Press), p. 283.
Shane, W. W. 1977, *Bull. AAS*, **9**, 362.
Silk, J. 1970, *Ap. J.*, **160**, 793.
———. 1977, *Ap. J.*, **211**, 638.
Tammann, G. 1974, in *Supernovae and Supernova Remnants*, ed. C. B. Cosmovici (Dordrecht: Reidel), p. 155.
van den Bergh, S. 1968, *J.R.A.S. Canada*, **62**, 1.

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