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SCOTT, John Stewart, 1950-
TURBULENCE AND PARTICLE ACCELERATION.

The University of Arizona, Ph.D., 1975
Physics, astronomy and astrophysics

Xerox University Microfilms, Ann Arbor, Michigan 48106

TURBULENCE AND PARTICLE ACCELERATION

by

John Stewart Scott

A Dissertation Submitted to the Faculty of the

DEPARTMENT OF ASTRONOMY

In Partial Fulfillment of the Requirements
For the Degree of

DOCTOR OF PHILOSOPHY

In the Graduate College

THE UNIVERSITY OF ARIZONA

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THE UNIVERSITY OF ARIZONA

GRADUATE COLLEGE

I hereby recommend that this dissertation prepared under my
direction by John Stewart Scott
entitled TURBULENCE AND PARTICLE ACCELERATION

be accepted as fulfilling the dissertation requirement of the
degree of Doctor of Philosophy

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ACKNOWLEDGMENTS

The author acknowledges gratefully the assistance of Drs. W. J. Cocke, J. E. Felten, R. Kirschner, D. N. Schram, P. A. Strittmatter, J. C. Theys, and N. J. Woolf in the preparation of this work. The author owes Drs. R. A. Chevalier and A. Pacholczyk special thanks for collaboration on various aspects of the work. The writer further thanks the following people for help and inspiration during the trying years of graduate school: Bob Carswell, Roger Chevalier, John Cocke, Andrezj Pacholczyk, Rabbi Peter Strittmatter, Rodger Thompson, Bill Tifft, Smilin' Ray Weymann, John Whelan, Ray White, Bob Williams, and Pacific Pete Worden.

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ABSTRACT

A model for the production of high energy particles in the supernova remnant Cas A is considered. The ordered expansion of the fast moving knots produces turbulent cells in the ambient interstellar medium. The turbulent cells act as magnetic scattering centers and charged particles are accelerated to large energies by the second order Fermi mechanism. Model predictions are shown to be consistent with the observed shape and time dependence of the radio spectrum, and with the scale size of magnetic field irregularities. Assuming a galactic supernova rate at $\frac{1}{50} \text{ yr}^{-1}$, this mechanism is capable of producing the observed galactic cosmic ray flux and spectrum below 10^{16} eV/nucleon. Several observed features of galactic cosmic rays are shown to be consistent with model predictions.

A model for the objects known as radio tail galaxies is also presented. Independent blobs of magnetized plasma emerging from an active radio galaxy into an intracluster medium become turbulent due to Rayleigh-Taylor and Kelvin-Helmholz instabilities. The turbulence produces both in situ betatron and 2nd order Fermi accelerations. Predictions of the dependence of spectral index and flux on distance along the tail match observations well. Fitting provides values of physical parameters in the blobs. The relevance of this

method of particle acceleration for the problem of the origin of X-ray emission in clusters of galaxies is discussed.

CHAPTER 1

INTRODUCTION

The need for a mechanism capable of naturally accelerating elementary particles to ultra-relativistic energies first became apparent with the discovery of cosmic rays in our own galaxy.

This need for a viable acceleration process was intensified when it became clear that most of the radiation from extragalactic radio sources was synchrotron radiation (Shklovsky 1960). The existence of large numbers of ultra-relativistic electrons may therefore be inferred.

Fermi (1949) first pointed out that particles with initial velocities near the speed of light could be accelerated in high conductivity plasmas. Magnetic field irregularities frozen into the plasma can scatter fast moving particles. In order to maximize entropy in the particle-irregularity system, the particle energies must approach equipartition with the kinetic energy at the moving irregularities. This implies acceleration of the particles. Since the necessary conditions are believed to be common both in our galaxy and in others, Fermi's mechanism (which he applied specifically to the acceleration of galactic cosmic rays) was attractive.

Soon after the publication of the Fermi mechanism it was realized (Cocconi 1951, Terletsii and Logunov 1951) that turbulence in an astrophysical plasma provides a situation just such as that described by Fermi, i.e., randomly moving magnetic irregularities. The idea of turbulence as a phenomenon which results in Fermi acceleration was refined by Kardashev (1962) and Ginzburg, Pickelner, and Shklovsky (1955).

Tsyтович (1970) has also proposed that plasma turbulence may result in particle acceleration through "plasma reactor" type mechanisms. As this type of mechanism requires that the plasma be optically thick to a substantial fraction of its own radiation (a situation which does not occur in any of the applications considered here), plasma reactors will not be dealt with in this work.

It is the intention of the writer to apply the simpler form (i.e., Fermi acceleration) of turbulent particle acceleration to specific astrophysical problems. Phenomena such as supernova explosions and mass ejection by active galaxies, which involve violent stirring of a plasma, are particularly interesting. This is because the observed motions of the ejected mass allow estimates of turbulent velocities and scale sizes to be made. Given these parameters, the Fermi acceleration coefficients may be determined. This enables the investigator to make specific predictions about the relativistic particle spectrum which may be

compared with observations. The lack of such predictions has been a weakness in previous discussions of acceleration mechanisms.

The need for viable acceleration theories is, at present, acute for the particular applications of galactic cosmic ray production and high energy electron acceleration in extragalactic double radio sources.

The most widely discussed theories of cosmic ray production have been: Fermi's (1949) acceleration by interstellar clouds; Colgate and Johnson's (1960) supernova shockwave hypothesis; the pulsar hypothesis (e.g., Goldreich and Julian 1969); and the strong wave acceleration mechanism of Gunn and Ostriker (1969). Each of these theories has been shown to suffer major difficulties (e.g., Landstreet and Angel 1971).

If it is assumed that the plasma, magnetic fields, and relativistic electrons which are responsible for the radio emission in the canonical extragalactic double radio sources are ejected from the nuclei of active galaxies, one can, by measuring the relative sizes of the radio lobes and the associated galactic nuclei, estimate the adiabatic particle energy losses. Because of such estimates Longair, Ryle, and Scheuer (1973) have concluded that the observed radio brightness in the double lobed extended sources can be most easily understood if some form of in situ particle acceleration occurs.

In many ways, the ejection of these independent plasma entities (which contain the particles and fields) into the intergalactic matter is phenomenologically similar to the expansion of the supernova ejecta into the interstellar medium. Because the structure of the particular double sources known as radio tail galaxies allows detailed study of a succession of different aged plasma ejection events, the time evolution of a typical ejected plasma entity can be inferred. This time evolution will be investigated here for evidence of turbulent particle acceleration.

CHAPTER 2

COSMIC RAY ACCELERATION IN THE CAS A SUPERNOVA REMNANT

There are several indications that particles are being accelerated to ultra-relativistic energies in supernova remnants. As an example, the writer will consider the Cas A remnant, which shows both a time decreasing radio luminosity and a general flattening of its power law spectrum. The general appearance of Cas A, both optical and radio, is similar to that found in several of the remnants which are of age 0-1000 years and have total energy 10^{51} - 10^{52} ergs. Of these remnants, Cas A is by far the best studied observationally. A backwards extrapolation in time of the radio surface brightness of Cas A may be performed either by assuming that the present rate of flux change was constant throughout the remnant's history, or by extrapolating along the Σ -D relation (the Σ -D relation is simply an observationally determined surface brightness vs. remnant size relation, c.f., Woltjer 1972). Such an extrapolation predicts a brightness considerably larger than that observed in most young (0-85 years) supernova remnants (de Bruyn 1973).

If Cas A is a fairly typical supernova remnant, as is assumed here, this indicates that the radio flux from

these objects builds up on a timescale of order 100 years. This increase in flux is probably due to acceleration of the relativistic electrons which are emitting the observed synchrotron radiation. Particle acceleration is also indicated in a general way by the flattening of the spectrum, since most loss processes tend to preserve or steepen a power law.

Gull (1973a) has constructed a model which predicts particle acceleration in supernova remnants, and has compared the model predictions with observations of Cas A and Tycho's SNR (Gull 1973b). This model, however, depends on the deceleration of the supernova ejecta, and this seems to be contrary to the evidence for Cas A (van den Bergh and Dodd 1970). Van den Bergh and Dodd found that the space motions of the ejecta, determined from proper motion and radial velocity data, indicate a common origin (both in space and time) only if the ejecta velocities have been constant throughout the remnant's lifetime.

Therefore the observed features of Cas A will in this work be compared with a model wherein particles are being continuously accelerated to high energies by the second order Fermi (1949) mechanism. In this model, the requisite magnetic scattering centers are regions of tangled magnetic fields caused by turbulent vortex cells in the plasma of the emitting region. These cells are produced in the ambient interstellar medium by the motion of high

density knots, which are streaming away from the center of the remnant with velocities of 4000-8000 km/sec (van den Bergh and Dodd 1970). It is assumed the knots formed early in the expansion of the remnant by a process such as described by Chevalier (1975). Chevalier proposed that the formation of the knots is a result of thermal instability in the dense shell behind the cooling shock which one expects to be traveling through the envelope of a star whose interior has exploded. This effect should produce knots with a range of velocities very close to the range observed in Cas A. This, of course, implies a geometry just such as that seen, i.e., a spherical shell of outflowing knots with a smooth positive velocity gradient as function of distance from the center of the remnant.

The turbulence induced in the interstellar gas by the passage of the knots amplifies the ambient interstellar magnetic field on a very short time scale. The constantly expanding spherical shell of knots, then, is a region where particles can be accelerated by the betatron mechanism. This serves as a source of fast ($V \sim C$) particles, a prerequisite for the aforementioned Fermi acceleration. The Fermi process is thus continually resupplied with particles from the interstellar medium through which the knots sweep.

In addition to inducing hydrodynamic motions in the interstellar gas, the passage of the knots through the ambient medium shock heats the medium (Shklovsky 1973). The

heated gas undergoes adiabatic expansion with a typical differential expansion velocity, V_{exp} . This expansion velocity, which is responsible for high energy particle adiabatic losses, is shown in Chapter 4 to be much smaller than the turbulent velocity, V_t , which is a typical relative velocity of the scattering centers and is responsible for the acceleration. Therefore, net particle energy gain is expected. We note that V_{exp} should not be confused with the velocity, V , of the fast moving knots. The quantity V defines the rate of increase of the total dimension of the remnant.

In our model the predicted turbulent scale size, amount and spectrum of high energy electrons, and flattening of the electron spectrum are shown to be in agreement with observations of Cas A.

This mechanism is capable of producing $\sim 10^{51}$ ergs of cosmic rays over the lifetime of the Cas A supernova remnant. The amount of mass in the fast moving knots (Minkowski 1968) shows that this is energetically feasible. We, therefore, compare model predictions concerning production of high energy nuclei with known features of galactic cosmic rays.

The Magnetic Scattering Centers

It is here proposed that the directed motion of the optically observed knots (van den Bergh and Dodd 1970)

"stirs" the tenuous gas which forms most of the volume of the remnant. The stirring of the interstellar gas creates vortices therein, which twist and amplify the frozen-in interstellar magnetic field, B , causing irregularities capable of scattering charged particles. Since these vortices are a form of turbulent motion, the scattering of particles becomes a stochastic process.

We expect the initial motion of the stirred gas to be characterized by a velocity, V_t , and a scale size, λ , comparable to the speed and size of the knots. Since the subject of strong plasma turbulence is not at present well understood, we do not know whether the initially formed vortices will break up into smaller cells in a manner analogous to the more familiar, homogeneous fully developed turbulence. It is clear, however, that the gyroradii of all relevant particles are considerably smaller than their mean free collision distances. Plasma motions, therefore, will be influenced by the presence of the magnetic field and the forces which result from the twisting of the field. Specifically, the field will transport momentum in such a way as to effect the allowed turbulent scale sizes. To see this we calculate an effective Reynolds number, R , defined by $R \equiv \text{Reynolds stress/magnetic stresses}$. Thus:

$$R = \frac{\rho (\vec{V}_t \cdot \vec{\nabla}) \vec{V}_t}{\vec{V} B^2 / 8\pi + \frac{1}{4\pi} (\vec{B} \cdot \vec{\nabla}) \vec{B}} \approx \frac{\rho V_t^2}{B^2 / 8\pi} \quad (1)$$

In other words, an analogy has been drawn with conventional hydrodynamics. There the Reynolds number is the ratio of Reynolds stress to viscous stress; i.e., it is the ability of the viscosity to transport momentum which determines the flow pattern. (This assumes that the flow is determined by the Reynolds number; c.f., Batchelor 1967.) The value of B to be used in the equation for R should reflect the strength of the magnetic field after the field has been amplified by the turbulent region. The problem of the amplification of a magnetic field in a turbulent medium has been studied by several authors. In summarizing the works of Saffman (1964), Pao (1963), and Parker (1963), Kraichnan and Nagarajan (1967) conclude that regardless of the ratio of magnetic diffusivity to fluid viscosity, an initially unamplified field will die away after first experiencing amplification or, at best, be increased to a steady state value short of equipartition of magnetic and turbulent kinetic energies. In the case of the supernova remnant we face the additional complication of the general expansion of the medium.

The e-folding time for field amplification, τ_{amp} , may be found from the dynamo term in Ohm's law:

$$\frac{dB}{dt} \text{ (dynamo)} = \vec{V} \times (\vec{V}_t \times \vec{B}) \quad (2)$$

therefore

$$\tau_{\text{amp}} = \frac{B}{\dot{B}} \approx \frac{\lambda}{V_t} \quad (3)$$

where λ is the turbulent scale size. Since the maximum

value of λ is of the same order as the knot size, $\tau_{\text{amp}} \leq 2$ years. The initial interstellar field is $\sim 10^{-3}$ times the equipartition value and therefore it should take only ~ 15 years for amplification to equipartition. It therefore seems likely that after a point very early in the history of the remnant, the value of the average field at a point in the remnant which has previously become turbulent will be slowly decreasing.

It should be noted that the radio brightness profile of the Cas A remnant is consistent with the above discussion. The outer parts of the remnant are brighter, relative to the central region, than would be expected in terms of a simple shell model (Rosenberg 1970). The author suggests that the excess brightness is due to a larger value of the magnetic field in the outer regions; i.e., the regions which have become turbulent more recently. This suggested spatial dependence of the magnetic field would occur naturally given the above proposed time dependence of the average field, B , in each local volume of a turbulent medium. Estimates of B near the outside edge of the remnant obtained by Rosenberg also indicate a higher field in this region.

Thus, since the observations support the hypothesis of a locally rapidly increasing, followed by a slowly decreasing, magnetic field, we maintain that the transport of momentum, which implies in this case suppression of small scale motions, by the field will have been at least as

effective as it is at present throughout most of the remnant lifetime. Therefore, to estimate this effect by consideration of our effective Reynolds number, R , we may use a value of B obtained from present observations. Woltjer (1972) has obtained a mean value of $B \sim 3 \times 10^{-4}$ Gauss by the usual assumption of equipartition of the energies in the magnetic field and in the relativistic electrons. Using $\rho = 1.6 \times 10^{-24}$ gm/cm³, $V_t = 6000$ km/sec (it is assumed that the knots, moving with velocity V , induce gas movements such that $V_t \sim V$) and Woltjer's value of B , one obtains $R \sim 170$. Values of R in this range indicate flow patterns with scales comparable to the size of the objects, in this case the knots, disturbing the gas (cf., Batchelor 1967). Thus we expect vortices in the interstellar medium with velocities and sizes typical of knot parameters.

Radio polarization data by Downs and Thompson (1972) show evidence for the scale of magnetic field irregularities in the remnant being the same as the smallest knot size (1" arc), in good agreement with our estimate. A rotation measure was determined from the frequency dependence of the polarization. The observed rotation measure was much less than the expected value calculated from estimates of the field strength and electron density. Since the rotation measure is actually proportional to the line of sight component of the field, the effective field can be reduced by irregularities which randomly change the sign of the

projected field. In order to bring the calculated rotation measure into agreement with observations, Downs and Thompson (1972) propose that irregularities with a typical scale size the same as the knot size exist throughout the remnant. We therefore have observational evidence that implies the existence of the large scale vortex cells. Thus one expects $\lambda \approx 4 \times 10^{16}$ cm (1" arc at a distance of 2.8 kpc). Since, as we have seen, the efficient transport of momentum by the magnetic field produces large scale motions more readily than small disturbances, the thermal dissipation rate is expected to be relatively slow. Thus the turbulence will not decay rapidly. Observations showing the small radio polarization throughout Cas A support the hypothesis that large scale irregularities in the magnetic field have a long lifetime.

Gain Processes

The turbulent cells move randomly relative to each other with the typical velocities imparted by the knots; i.e., $V_t \approx 6000$ km/sec. This velocity, which is near the sound velocity in the shock heated medium behind the knots, and the cell scale size, λ , are therefore the relative velocity and the mean distance between the scattering centers. Fast moving particles scattered in a stochastic process experience an energy gain by the second order Fermi mechanism (Fermi 1949), such that the average gain per

collision is $\Delta E_{\text{col}} \approx E \left(\frac{v_t}{c} \right)^2$. This results in a systematic particle energy increase with e-folding time, τ_g , such that

$$\tau_g = E/\dot{E} = \frac{\lambda c}{v_t^2}. \quad (4)$$

With the values of the parameters derived previously we obtain $\tau_g \approx 80$ years.

The gain time, τ_g , is correct only for particles whose initial velocities were near the speed of light. The thermal electrons, whose temperature should range from the observed X-ray temperature of 1.5×10^7 °K to the proton shock temperature of 5×10^8 °K (Gorenstein, Harnden, and Tucker 1974) provide an adequate source of fast electrons. The thermal protons have typical velocities of only ~ 6000 km/sec.

One must, however, consider the effect of the betatron acceleration mechanism on these non-relativistic protons (or any other nucleus with an original hydrodynamic velocity of 6000 km/sec). By the first Alfven invariant $E \propto B$; as each cell intensifies and de-intensifies its magnetic field, the thermal ion energies increase and decrease. However, since the average field, \bar{B} , in each new region of interstellar material which becomes turbulent as the knots stir it increases by a factor $\sim 10^3$ in only 15 years. Thus charged particles (electrons or ions) whose initial thermal velocity is ~ 6000 km/sec are accelerated to

a velocity near C , and energy $\equiv E_0$. These particles would then be moving fast enough to be efficiently accelerated to ultra-relativistic energies ($\sim 10^{13}$ eV/nucleon) during a time equal to the present age of the Cas A supernova remnant.

Loss Processes

Shklovsky (1966) has suggested that adiabatic expansion losses are important in Cas A and are responsible for the general decrease in radio luminosity. The expansion losses may be thought of as Fermi deceleration. Since the expansion velocity of the ambient, stationary medium (and hence of the magnetic scattering centers), V_{exp} , is a directed rather than a random relative motion the loss is first order in V_{exp}/c ; i.e., $\Delta E_{\text{coll}} = E(V_{\text{exp}}/c)$, giving a loss time

$$\tau_L = \lambda/V_{\text{exp}} \quad . \quad (5)$$

In order to estimate V_{exp} , one needs to estimate the systematic relative velocity of two points in the remnant a distance λ apart. Besides creating turbulence in the medium, the shockwaves produced by the knots heat the medium. The expansion of a volume of hot gas will approach the Sedov (1959) solution, which describes the evolution of a point explosion in a uniform medium. Sedov finds that the gas velocity is approximately proportional to the distance from the center, i.e., $V_f \sim r V_s/r_s$ where V_f is the flow velocity, r the radial distance from the center of the

explosion, and V_s and r_s the shock velocity and position. Although not strictly applicable due to their assumption of a smooth shell, the calculations of Rosenberg and Scheuer (1973) and Gull (1973a) indicate that larger velocity gradients are to be expected in the transient phases before the Sedov phase is established. If the Sedov velocity profile does apply, the relative velocity of two adjacent vortex cells is the same in the radial direction as in the two tangential directions and we have $V_{\text{exp}} = \Delta V_f = \frac{\lambda V_s}{r_s} \approx \frac{1}{100} V_s = 60-80 \text{ km s}^{-1}$. If the velocity gradient is larger, the relative velocity in the radial direction is larger, and V_{exp} is increased. Here, V_s and r_s are taken to be the mean velocity and radius of the shell of fast moving knots. This seems reasonable since the radius of the X-ray source $\sim r_s$.

The differential expansion of the ambient medium is responsible for the first order losses. This velocity, V_{exp} , is much smaller than the shock velocity, V_s , which is typical of the overall expansion of the volume enclosed by the fast moving knots. The turbulent velocity, V_t , is responsible for the second order acceleration and since $V_t \sim V_s \gg V_{\text{exp}}$ acceleration can occur. In fact,

$$V_{\text{exp}} = V_s \frac{\lambda}{r} \approx V_t \frac{\lambda}{r} \quad (6)$$

where r is the size of the remnant. Thus $V_{\text{exp}}/c \sim \frac{V_t}{c} \frac{\lambda}{r}$ which gives numerically:

$$\frac{V_{\text{exp}}}{c} \sim \left(\frac{V_t}{c}\right)^2 . \quad (6a)$$

Therefore, τ_L , the loss time, is of the same order as τ_g .

The synchrotron losses of all relevant particles are insignificant.

The Particle Spectrum

Given the gains and losses described in the preceding sections, the appropriate equation of continuity in particle (both electrons and ions) energy space is given by (Morrison 1961, Cocke 1975):

$$\frac{\partial N}{\partial t} = \frac{1}{\tau_g} \frac{\partial}{\partial E} (E^2 \frac{\partial N}{\partial E}) + \left(\frac{1}{\tau_g} + \frac{1}{\tau_L}\right) \frac{\partial}{\partial E} (EN) + q\delta(E-E_0) \quad (7)$$

where N is the particle density in energy space, and the writer is assuming a constant supply of particles (at some low energy E_0) $q\delta(E-E_0)$. The time scale for the change in shape of the synchrotron electron spectrum (Dent, Aller, and Olsen 1974), is observed to be almost an order of magnitude greater than the time scales involved in the above equation. Therefore one may, as a first approximation, solve the above equation by a separation of the time and energy variables. The solution has the form

$$N \propto e^{-at} E^m \quad (8)$$

for $E > E_0$ where

$$m = -1 - \frac{\tau_g}{2\tau_L} - \frac{\tau_g}{2} \sqrt{\tau_L^{-2} - 4a/\tau_g} . \quad (9)$$

If it is assumed that the observed flux decrease from Cas A is indeed a result of adiabatic loss then the separation constant, a , is given by Dent et al. (1974) as $a \approx .01 \text{ yr}^{-1}$. Thus, within the accuracies of our estimates of the quantities τ_g and τ_L we may easily match the observed power law in Cas A, $\bar{m}_{\text{obs}} \approx -2.5$. For example, if $\tau_g = \frac{1}{a} = 2.2 \tau_L$, then $m = -2.56$.

As discussed previously, it seems likely that the average magnetic field should be slowly decreasing as the remnant expands. As \bar{B} decreases the tendency for the medium to become fully turbulent is enhanced. This lowers the mean scattering distance, which decreases τ_g and increases τ_L (through V_{exp}), resulting in a flattening of the spectrum, in agreement with the observations.

Cosmic Rays

If one assumes a galactic supernova rate of $1/50 \text{ yr}^{-1}$, each supernova need only produce $\sim 6 \times 10^{50}$ ergs in fast particles in order to supply the observed (through galactic synchrotron background) galactic cosmic ray energy density ($\sim 1 \text{ eV/cm}^3$).

If one assumes that the proton injection rate, q , is the same as the electron injection rate, and that the observed electron spectrum extends down to energy E_0 , one may estimate the amount of energy in the fast protons in Cas A. Here E_0 is the energy of the particles after

experiencing betatron acceleration, and it depends on the particles' original energy. We integrate Equation (8) to find the total energy in protons

$$E_T = \int_{E_0}^{\infty} N E dE \approx q \tau_L E_0^{m+2} \quad (10)$$

(the constants q , τ_L arise from the solution of Equation [7] but are irrelevant here) which along with our assumption that $q(\text{proton}) = q(\text{electron})$ yields a total proton energy which depends on the ratio $[E_0(\text{proton})/E_0(\text{electron})]^{m+2}$, and on the total electron energy, which is inferred from observations. This ratio, of course, depends only on $\frac{m_p}{m_e}$ and on the ratio of the velocities of each species of particle prior to acceleration by the betatron mechanism; i.e., it depends on whether or not the electrons are thermalized in a time short compared to the 15 year betatron acceleration time. If the electrons are completely thermalized before the betatron acceleration, the total proton energy content is $\sim 2 \times 10^{50}$ ergs. If there is no electron thermalization the proton energy is $\sim 2 \times 10^{51}$ ergs. If the thermalization time is greater than the ~ 15 years over which the betatron process occurs, we expect the proton energy to be closer to the second case, i.e., $\sim 10^{51}$ ergs.

The upper limit is determined by the available kinetic energy of the ejecta, which Minkowski (1968) has estimated to be 2×10^{51} ergs. We expect this energy to be divided between thermal motions, turbulent motions, and

relativistic particles, giving an upper limit of order 6×10^{50} ergs for the particles. The magnetic energy may be less than the above mentioned energy as it only acts as an intermediary between the turbulent motions and relativistic particles. The author therefore believes that the production of the galactic cosmic rays by the mechanism outlined in this paper is energetically feasible. Furthermore, the spectrum in the remnant ($m \sim -2.5$) is very close to the galactic cosmic ray spectrum ($m \sim -2.6$). It is therefore interesting to compare model predictions with several other observed features of galactic cosmic rays

The lifetime of the knots in a typical supernova remnant is of order 1000 years (Chevalier 1975). The most important effects seems to be evaporation of the knots caused by frictional heating as the knots pass through the interstellar gas. This remnant "particle production" lifetime is many particle energy gain times. Thus particles are accelerated continuously until the particle gyroradii become larger than the mean scattering distance, λ . Using parameters mentioned earlier, this occurs at an energy 3.6×10^{15} eV/nucleon. Hillas (1974) has reported a steepening of the cosmic ray spectrum at 4×10^{15} eV. Note that this model predicts slightly different turnover energies for varying species' charge to mass ratios. From the arguments presented in this model, one predicts a spectrum in substantial agreement with that observed for

galactic cosmic rays up to and including the spectral break at $\sim 5 \times 10^{15}$ eV. Significantly higher energy particles cannot be produced by this mechanism, but cosmic ray isotropy arguments indicate these higher energy particles are extragalactic in origin (cf., Meyer, Ramaty, and Webber 1974).

This method of cosmic ray production has implications for cosmic ray composition. The galactic cosmic rays have considerably more hydrogen than would be expected in a star which has processed material to iron. The composition is in fact consistent with such stellar material having been mixed with several times its mass in interstellar material (Schramm 1974) as would be expected in the model described here.

The production of cosmic rays in the relatively tenuous medium of supernova remnants is consistent with the observed cosmic ray path length $\sim 5 \text{ g cm}^{-2}$ (cf., Arnett and Schramm 1973). Furthermore, higher energy cosmic rays have a lower daughter to parent spallation ratio, indicating that the high energy cosmic rays have traversed less material in their lifetimes (Cesarsky and Audouze 1973). This is a natural consequence of our model since the particles attain higher energies late in the evolution of the supernova remnant, when the remnant material occupies more volume and is less dense.

Summary

The basic model presented in this paper may be represented schematically as follows: fast moving knots → turbulence → randomly moving magnetic scattering centers → 2nd order Fermi acceleration.

The model quantitatively fits observed features of the radio spectrum for the Cas A supernova remnant. The predictions in the model are not inconsistent with the galactic cosmic rays having been produced by this mechanism.

A similar model with knots replaced by galaxies may be responsible for the acceleration of particles in clusters of galaxies, and this possibility will be discussed in Chapter 3.

CHAPTER 3

RADIO TAIL GALAXIES

Radio galaxies with tails (see Fig. 1) were observed by Ryle and Windram (1968) in the Perseus cluster (NGC 1256 and IC 310); by Willson (1970) in the Coma cluster (5C4.81); and by MacDonald, Kenderdine, and Neville (1968) and Hill and Longair (1971) in the cluster 3C129. The original interpretation of the tail phenomenon, forwarded by Ryle and Windram (1968), was made in terms of a stream of relativistic particles emitted from a Seyfert galaxy (NGC 1275) interacting with the gaseous component of radio galaxies. This point of view seemed to be supported by the fact that in both cases studied (NGC 1265 and IC 310) the tails were pointing roughly away from the direction of NGC 1275. Further observations (Miley et al. 1972; Miley 1973; Wellington, Miley, and van der Laan 1973) however, revealed double structure of the radio tails; also a weak source CR 15 was detected in the Perseus cluster with a tail pointing roughly toward NGC 1275 rather than away from it. Thus a picture of a source independent on any other galaxy in the neighborhood emerged; the tailed source was assumed to be analogous to a double-lobed extended source, but

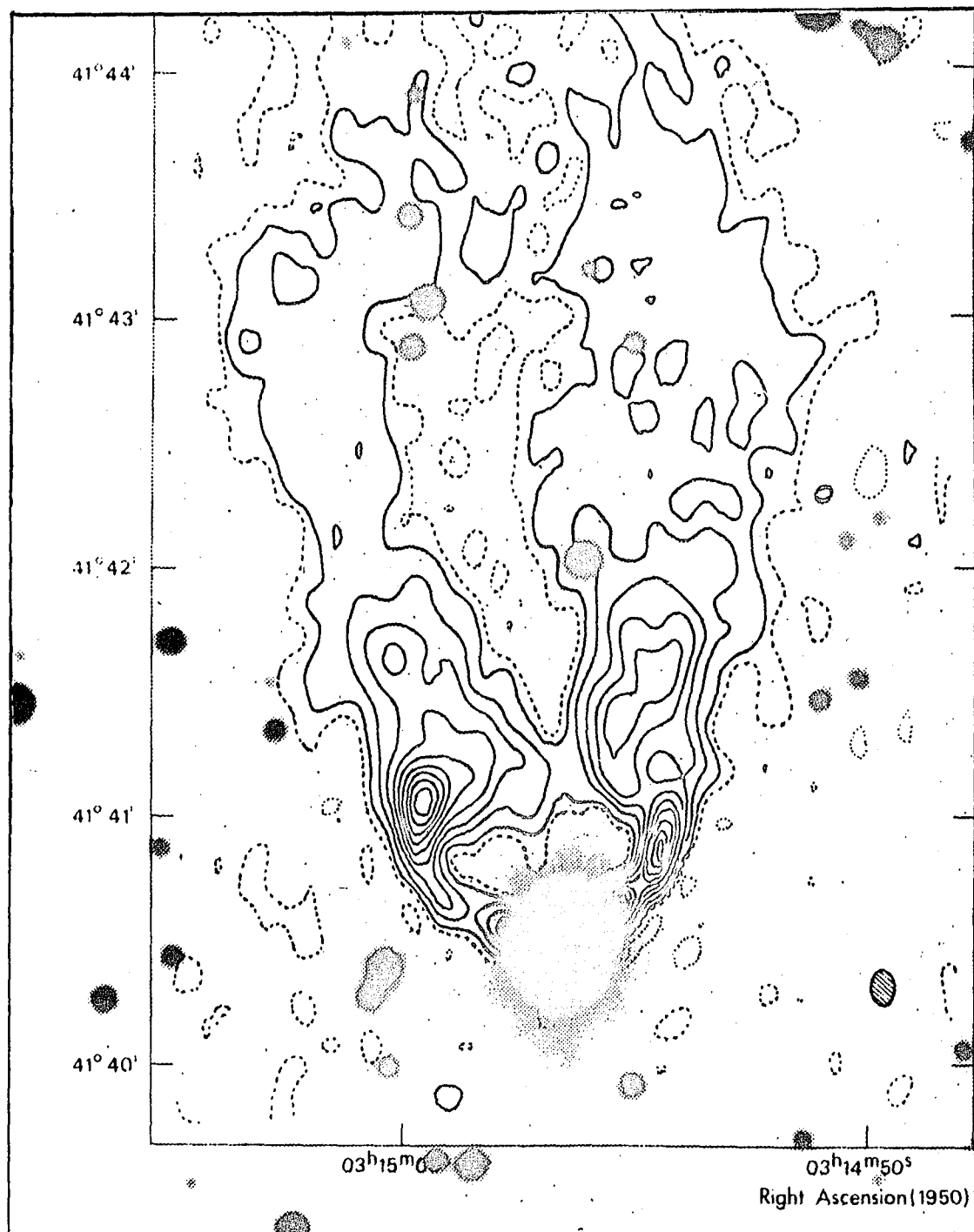


Fig. 1. 1415 MHz radio map of NGC 1265 (Miley 1973).

distorted into a double tail by the motion of the galaxy relative to an intracluster medium.

High resolution maps of NGC 1265 (Wellington et al. 1973) demonstrate clearly that the radio tail is made of discrete components, increasing in size with increasing distance from the head. Jaffe and Perola (1973) considered a model of a head-tail source in a cluster, wherein independent "blobs" of magnetized plasma, containing relativistic electrons, are ejected from the active galaxy. The "blobs" interact with an intracluster gas through which the galaxy is moving, and this produces the observed pattern of emission. In this model the acceleration of electrons was assumed to have taken place prior to the ejection of the blob from the parent galaxy. Subsequently the electrons were subject only to processes leading to losses of energy; i.e., adiabatic and synchrotron losses.

Adiabatic losses caused by the expansion of the plasma "blobs" theoretically lead to much steeper radio spectra than those observed. Because of this disagreement with the observational data, Jaffe and Perola (1973) abandoned this model, although they realized that the discrepancy with the observations might be eliminated if some type of particle acceleration mechanism were operable during the expansion of the blobs.

Rather than considering various acceleration processes, they proposed another model in which relativistic

particles were ejected into an essentially rigid galactic magnetosphere. The success of this second model depends entirely on the assumption that these free particles do not suffer expansion losses. This assumption is incorrect. Jaffe and Perola state that the group of particles injected into the magnetosphere retain the appearance of a blob because plasma waves prevent the particles from dispersing at the speed of light. We point out that this particle-wave scattering causes the particles to undergo adiabatic losses, since the volume wherein the particles are contained by the waves expands by an observationally determined amount; the particles lose energy to the receding plasma waves by first order Fermi deceleration. The magnetosphere model therefore does not alleviate the problem of the disagreement between the calculated and observed spectral indices. Therefore Jaffe and Perola's (1973) magnetosphere model will not be considered further here.

In the present paper we will adopt the basic idea of the expanding "blob" model, but in addition, we will consider the effects of second order Fermi acceleration in the magnetized turbulent plasma of the blobs. In our model, the requisite scattering centers are presumed to be magnetic irregularities caused by turbulent vortices in the plasma of the blob. As the blobs are ejected from the radio galaxy, the intracluster gas impinges upon the blob causing a Rayleigh-Taylor instability and giving rise to turbulence.

Thus the strength of the turbulence is determined by the relative velocity of the blob and the intracluster medium; i.e., the vector sum of the blob ejection velocity and the relative velocity of the galaxy and the intracluster medium. This turbulent velocity, V_t , which along with a scale size determines the particle energy gain time, is an order of magnitude larger than the blob expansion velocity, V_{exp} , which is responsible for the particle adiabatic losses. It is thus possible for the particles to be accelerated.

The Observations

The following are the main observationally determined characteristics of a typical tail (Miley 1973; Miley and van der Laan 1973; Hill and Longair 1971):

1. Geometry of the tail indicates that the blobs have a typical radius, r_0 , at the point of exit from the parent galaxy of a few kiloparsecs ($r_0 \sim 2$ kpc for 3C129). Subsequently the blobs expand by a factor ~ 3 . The blob reaches its maximum radius at a point which we define to be x_2 (60 kpc for 3C129) after which the blob size does not change significantly.
2. The two-frequency spectral index α_j^i ($i = 1400$ or 2700 MHz and $j = 400$ MHz) decreases slightly following the blob exit from the galaxy to a minimum which occurs near x_1 . The index then rises sharply

until the blob reaches a point near x_2 ; thereafter α_j^i remains relatively constant up to a large distance from the galaxy x_3 ($x_3 \sim 160$ kpc for 3C129). At values of x larger than x_3 , α_j^i rises again (Fig. 2).

3. The flux density F_ν increases sharply within a small distance from the parent galaxy, reaching a maximum at the point x_1 (this defines x_1), and in all observed cases $x_1 \leq x_2$. For 3C129, $x_1 \sim 20$ kpc (Miley 1973). The flux then declines fairly rapidly to a value at the point x_2 lower by a large factor (7.5 for 3C129); after the point x_2 the flux follows a fairly slow decline until it disappears into the background.
4. The linear polarization also dips reaching a minimum value near x_1 ; for 3C129 $\pi_L(x_1) \sim 1\%$, then increases in the region between x_2 and x_3 and becomes rather high ($\sim 60\%$) beyond x_3 .

The writer proposes that the flux, spectral index, and polarization data in the portions of the tail near the galaxy are consistent with the hypothesis that turbulence is formed in the emerging blobs. The following qualitative arguments are intended to demonstrate this.

Fig. 2. The observed spectral index α_j^i as a function of distance along the tail, x_j . We show α_{408}^{1415} for NGC 1265 (Riley 1973), α_{408}^{1415} and α_{408}^{2700} for 3C129 (Riley 1973), and α_{408}^{1415} for 5C4.81 (Willson 1970). Note the following pattern: a dip, followed by a plateau or slowly rising region, followed by a sharply rising region.

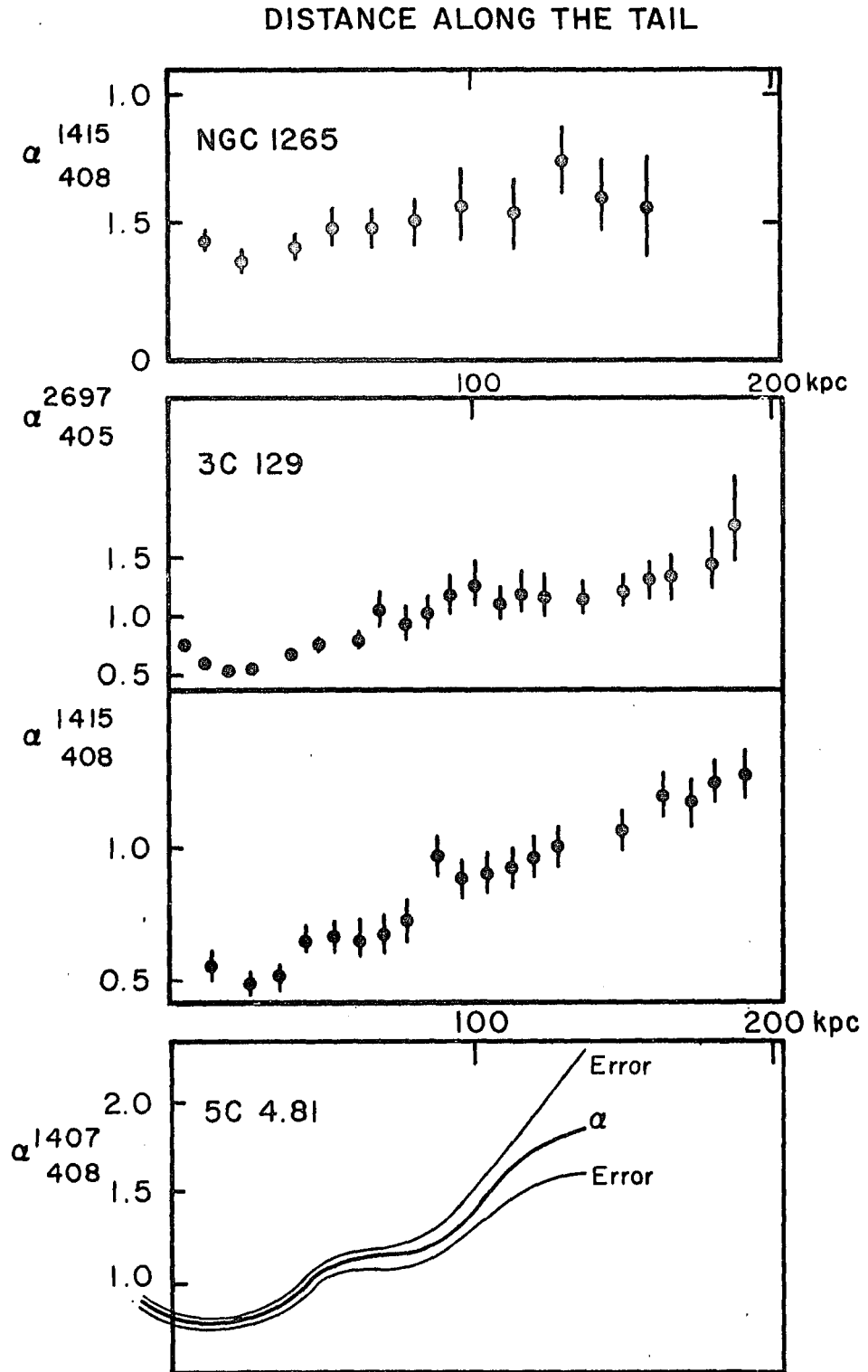


Fig. 2. The observed spectral index α_j^i .

The Model

The author adopts the essential idea of an expanding blob model and draws relevant observational parameters from Jaffe and Perola (1973).

Blake (1972) has shown that the ratio of instability (either Rayleigh-Taylor or Kelvin-Helmholz) e-folding time, t , to the ram pressure stopping time of the blob, τ_s , is determined only by the ratio ρ_b/ρ_{ic} , i.e., the ratio of the densities in the blob and in the intracluster medium. Blake also points out that $\rho_b/\rho_{ic} \approx D/h$ where D is the stopping distance and h is the pressure scale height in the blob. Jaffe and Perola (1973) give $D = 55$ kpc for 3C129, and assuming $h \approx r$, where r is the blob radius, we obtain $\rho_b/\rho_{ic} \approx 16$. We apply Blake's (1972) formula:

$$\frac{t_{RT}}{\tau_s} \approx \left(\frac{\rho_{ic} k_{min}}{\rho_b k} \right)^{1/2} \quad \text{and} \quad \frac{t_{KH}}{\tau_s} \approx \frac{\rho_{ic}}{\rho_b} \eta \frac{k_{min}}{k} \quad (1)$$

where k is the wave number of a blob surface perturbation, $k_{min} = \frac{1}{r}$, and η is the ratio of the density of the unshocked intracluster gas to the density of the gas which has passed through the bow shock of the blob. We thus find that even for $k = k_{min}$ the blob is both Rayleigh-Taylor and Kelvin-Helmholz unstable with e-folding times short compared to the blob stopping time.

In fact, the velocities in the instabilities will reach the sound velocity of the blob front in a time $\approx \left(\frac{\rho_{ic}}{\rho_b} \right)^{1/2} \tau_s$. The sound velocity in this region is of the

same order as the relative velocity V_0 , of the blob and the intracluster gas. Thus the strength of any turbulence which develops in the blob as a result of the instability motions is determined by V_0 .

The plasma parameters derived by Jaffe and Perola (1973) indicate that typical ion and electron collision lengths are much greater than the size of the blobs, rendering the usual plasma viscosities meaningless. In such a situation, however, we may define an effective Reynolds number $\frac{\rho_b V_0^2}{H^2/8\pi}$ (Chapter 2, also, Scott and Chevalier 1975). Using Jaffe and Perola's (1973) values of these parameters and a value $H = 7 \times 10^{-6}$ Gauss derived from the usual minimization of total energy, we find $R_{\text{eff}} = 10^3$ in the emerging blobs. Thus, the motions induced in the blob by the instabilities will result in well-developed turbulence in the blob plasma. We see, then, that the assumption that there is no turbulence (manifested through the use of Jaffe and Perola's derived dynamical quantities) leads to a contradiction.

Furthermore, the turbulent mixing of blob material and intracluster gas should cause the blobs to gradually lose their individual identities as the parent galaxy moves further away. That this is the case may easily be seen from the observations of Miley (1973). We therefore conclude that turbulence is formed in the blobs. The turbulent pressure may be added as a correction to Jaffe and Perola's

(1973) dynamics, allowing derivation of relevant physical parameters. Using the corrected values of these parameters we may re-derive R_{eff} and use it as a consistency check on the assumption that turbulence is formed.

The presence of turbulence in the plasma implies amplification of the magnetic field in the blob (Kraichnan and Nagarajan 1967). The amplification of the magnetic field in turn implies in situ acceleration of the relativistic particles by the betatron mechanism. Both of these effects increase the synchrotron emission in the plasma on a time-scale equal to the timescale for turbulent amplification of the magnetic field. Since the actual amplification is a result of the dynamo effect, we may estimate the e-folding time for field amplification, $\tau_{\text{amp}} \equiv H/\dot{H} \sim L/V_0$ where L is the size of the blob. Kraichnan and Nagarajan (1967) believe that the amplification will continue until the magnetic and turbulent kinetic energy densities are near equipartition, after which the field will remain constant or slowly decrease. For 3C129 $\tau_{\text{amp}} \sim 8 \times 10^5$ years, a value which is considerably smaller than the expansion time. We believe this explains the sharp flux increase between the point where the blobs emerge and the point x_1 . In fact, all of the objects studied by Riley (1973) and Miley (1973) exhibit this sharp rise to flux maximum at x_1 showing a qualitative agreement with the model presented here.

If this interpretation of the sharp flux rises is correct, the betatron acceleration ceases near x_1 . The expansion of the blobs continues after x_1 , however, and this should cause the flux to decrease, as observed.

The presence of the betatron acceleration for $x < x_1$ should result in a decrease in α_j^i . It will be shown in the next section that in this region all processes, except for synchrotron losses for very high energy electrons, are completely dominated by the betatron mechanism. The entire spectrum, including the break due to synchrotron losses is therefore simply translated to higher frequency, clearly resulting in the observed decrease in α_j^i . The dependence of α_j^i on x past the point x_1 , is determined by the Fermi acceleration and by adiabatic and synchrotron losses. It is demonstrated in the next section that the theoretical dependence of α_j^i on x for various combinations of these processes also agrees well with the observations.

The development of "noisy" turbulence postulated in this model would naturally result in a dip in polarization due to randomization of the magnetic field. This too is consistent with observations.

The Calculations

From the data presented previously one may define three regions: Region I, where expansion occurs, and the spectral index rises; Region II, where the blob radius and

spectral index are roughly constant; and Region III, where the spectral index and polarization rise sharply (Fig. 3).

We feel that these observationally defined regions result from the three likely combinations of processes affecting particle energies. Thus, Region I includes both Fermi gains (with acceleration coefficient f) and adiabatic losses (i.e., $f \neq 0$, $V_{\text{exp}} \neq 0$). For the purposes of the calculations we do not include the values of x such that $x < x_1$. Because of the dominance of the betatron mechanism in this region one cannot determine interesting physical parameters from a study of α_j^i . Therefore Region I will include values of x in the range $x_1 < x < x_2$, where most of the expansion occurs, and where adiabatic losses and Fermi gains are nearly balanced. Region II starts at x_2 , the point where blob expansion stops. Since the polarization is still fairly low, we conclude that in Region II Fermi acceleration is still occurring ($V_{\text{exp}} = 0$, $f \neq 0$). In Region III, which begins at the point x_3 , the polarization rises sharply to large values, indicating that the turbulence has dissipated. Thus the change in the spectral index is governed only by synchrotron losses ($f = V_{\text{exp}} = 0$).

In order to calculate the behavior of the spectral index with x we have to solve the equations of continuity for electrons in energy space (Kardashev 1962)

$$\frac{\partial N}{\partial t} = \frac{\partial}{\partial E} [(-\alpha E + \beta E^2)N] \quad (2)$$

Fig. 3. A "typical" radio tail -- This is a composite of observations of several tails showing the features we felt must be explained by our model. All graphs are plotted as a function of x . The processes occurring in each region are written at the bottom of the figure.

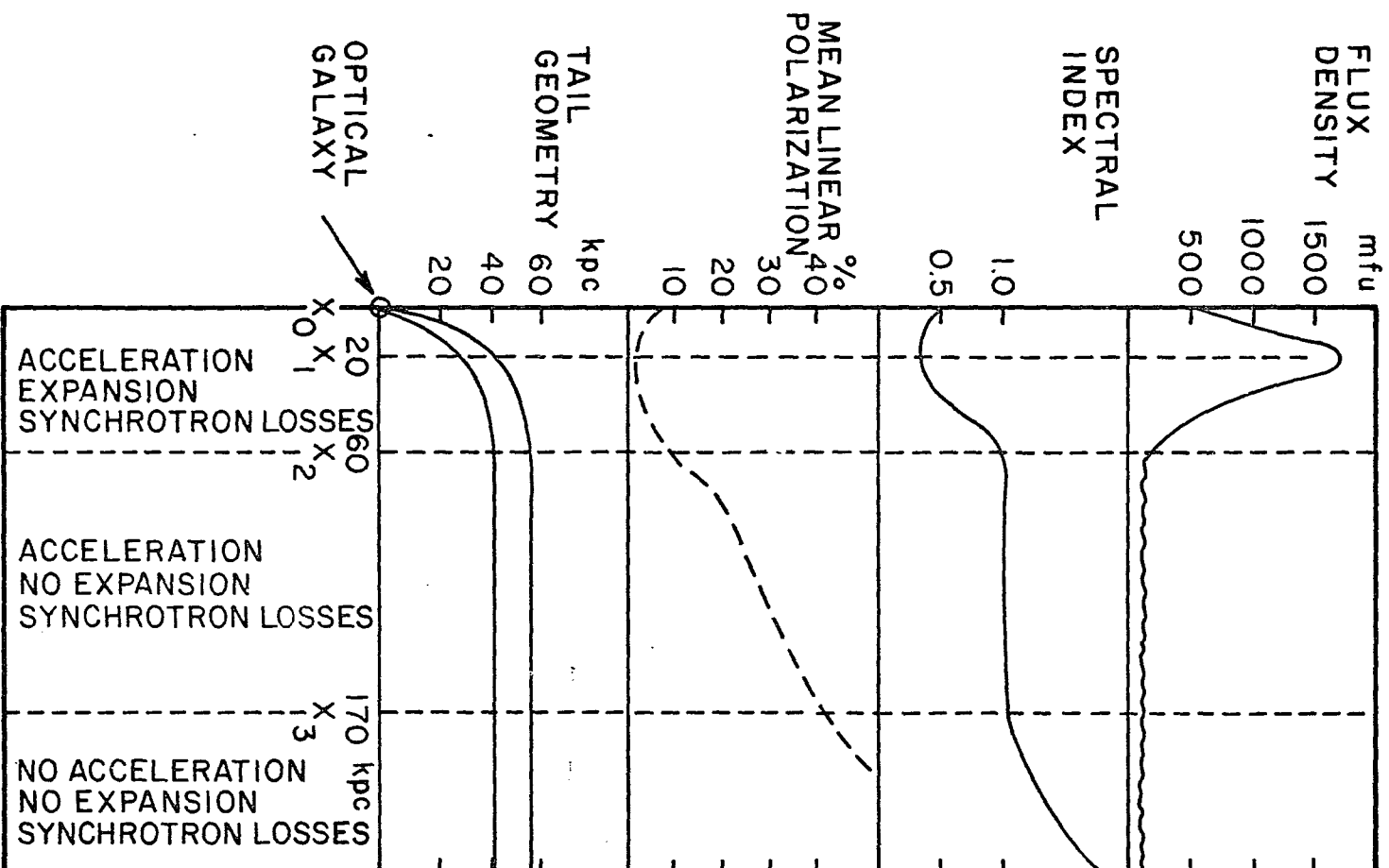


Fig. 3. A "typical" radio tail.

where $N = N(E, t)$ is the number density of relativistic electrons in the blob, $\beta E^2 = h r^{-4} E^2$ is the synchrotron energy loss rate, and where (Ginzburg et al. 1955, see Appendix A)

$$\alpha = \frac{V_t^2}{cL} - \frac{V_{\text{exp}}}{r} \approx \frac{V_B^2}{cL_1} \left(\frac{r_1}{r}\right)^{5/3} - \frac{V_{\text{exp}}}{r} \equiv f r^{-5/3} V_{\text{exp}} r^{-1}. \quad (3)$$

Here, L and V_t are the scale and velocity of turbulent elements responsible for the acceleration and $V_B \sim V_0$ is the turbulent velocity induced in the blob (i.e., largest scale motion). We are ignoring here a diffusion term which would introduce a smoothing effect on the features in energy spectrum but which does not affect significantly the spectral index itself. Assuming $N(E, t) = K E^{-\gamma}$, Kardashev's (1962) solution of Equation (2) is

$$N(E, t) = \text{const } r^3 E^{-\gamma} \exp\left[(\gamma-1) \int_{t_1}^t \alpha dt\right] \cdot$$

$$\left[1 - E_c e^{-\int_{t_1}^t \alpha dt} \int_{t_1}^t \beta e^{\int_{t_1}^t \alpha dt} dt\right] \quad (4)$$

where the cut-off energy E_c is given by

$$E_c = \frac{e^{\int_{t_1}^t \alpha dt}}{\int_{t_1}^t \beta e^{\int_{t_1}^t \alpha dt} dt}. \quad (5)$$

In Region I the integration of Equation (5) yields the

expression for E_c ,

$$E_c = \frac{f}{h} \frac{1}{z^{-7/3} - z e^{-\kappa(1-z^{-2/3})}}, \quad (6)$$

where $\kappa \equiv 3/2 (f/V_{\text{exp}})$ and $z = r/r_1$, and since $\kappa \gg 1$, only first order terms in κ^{-1} were retained. Also, to make the integration tractable, we made the approximation $V_{\text{exp}} = \text{const}$. Since V_{exp} increases at most by a factor of 2 from the average value in Jaffe and Perola's (1973) solution for V_{exp} , we let $V_{\text{exp}} = \bar{V}_{\text{exp}} \sim 3 \times 10^7$ cm/sec.

In terms of frequencies we have

$$\frac{1415}{v_c(I)} = \frac{7.4 \times 10^{-10} h^2}{H f^2 r_1^{14/3}} z^4 [z^{-10/3} - e^{-\tilde{\kappa}(1-z^{-2/3})}]^2, \quad (7)$$

where $v_c = 1.8 \times 10^{18} H E_c^2$ (see e.g., Pacholczyk 1970), and $\tilde{\kappa} = \kappa r_1^{-2/3}$.

We consider the situation in Region II, where the expansion losses are not playing an important role, but where the turbulent processes are continuing to accelerate the electrons. In this region synchrotron losses are balanced by acceleration gains leading to a formation of a plateau in the run of the spectral index with x . Indeed, in this region, the solution of Equation (5) is constant

$$E_c = \frac{f z_2^{7/3}}{h} = \text{const} \quad (8)$$

where $z_2 = z(x = x_2)$ and, consequently,

$$\frac{1415}{v_c(II)} = \frac{7.4 \times 10^{-10}}{H(r_2)} \frac{h^2}{f^2} r_2^{-14/3} = \text{const} \quad (9)$$

where $r_2 = r(x = x_2)$. This solution, of course, agrees with that of Equation (7) at $x = x_2$.

In Region III the turbulence has died out and the acceleration ceased to operate. The spectrum is dominated by synchrotron and inverse Compton losses which cause a further increase of the spectral index. This is described by the solution

$$E_c = \frac{1}{hr_1^{-4} \left(\frac{x-x^*}{V_G} \right)^2} \quad (10)$$

and

$$\frac{1415}{v_c(III)} = \frac{7.4 \times 10^{-10}}{H(r_2)r_2^8} h^2 \frac{(x-x^*)^2}{V_G^2}, \quad (11)$$

where V_G is the relative velocity of the galaxy and the intracluster gas and where x^* is a constant that can easily be determined by matching (9) at x_3 with (7); $x^* = 50$ kpc.

At x_3 (=165 kpc) the observed spectral index $\alpha = 1.0$. The observed spectral index at the end of the tail ($x_4^{\text{obs}} = 185$ kpc) is $\alpha = 1.3$. Calculations based on Equation (11) indicate that this value of the spectral index should be attained at a distance $x_4^{\text{calc}} = 150$ kpc.

As can be seen from Equations (7), (9), and (11), the predicted α_j^i should qualitatively fit the picture of a typical tail (see Fig. 3); i.e., it rises in Region I, is

flat in Region II, and rises again in Region III. A specific fit is made for 3C129, and is displayed in Fig. 4.

The dependence of the spectral index, α_j^i , on x in the three regions is derived from the above equations for the quantity $1415/\nu_c$ through the application of Jaffe and Perola's (1973) numerical tabulation $F(\nu_c)$, which is a function describing the shape of the cutoff for $\nu > \nu_c$. From their Figure 3 one may transform directly from $1415/\nu_c$ and $408/\nu_c$ to α_j^i .

The absolute value of α_j^i cannot be obtained without knowledge of the exact plasma turbulence parameters which determine α . These and other physical parameters both in the blobs and in the intracluster material may be derived by fitting the flux curve and the equations for α_j^i to the specific observations relevant to 3C129. We calculate first the change of the flux density during the expansion of the blob in Region I. The change of energy E of a particle during the adiabatic expansion of a turbulent blob can be described by the equation

$$\frac{dE}{dt} = \alpha E. \quad (12)$$

Integrating Equation (12) yields the result that the energy of an individual particle depends on the blob's radius as follows:

$$E = E_1 \frac{r_1}{r} e^{-\kappa r^{-2/3}}; \quad (13)$$

here we again assume $r \propto t$, i.e., $V_{\text{exp}} \sim \text{const.}$ The

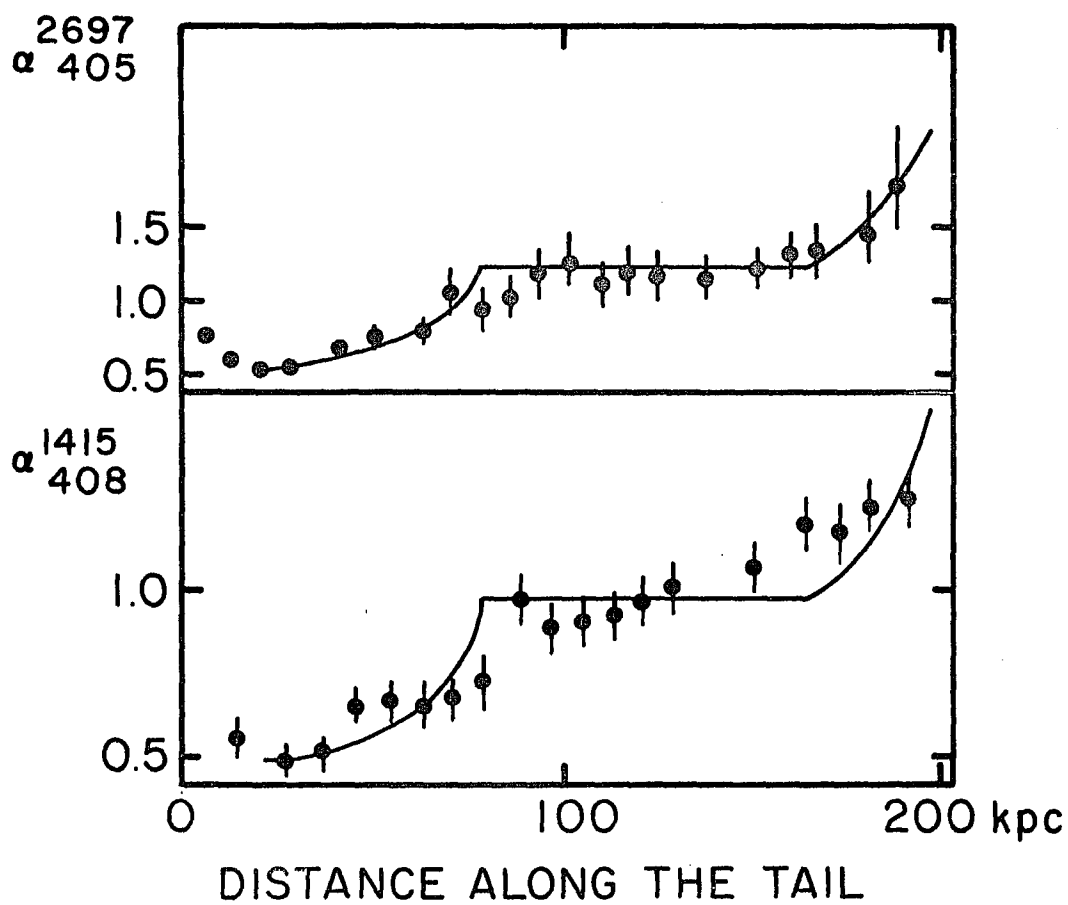


Fig. 4. Observations of Miley (1973) and Riley (1973) of $\alpha_{2697/405}$ and $\alpha_{1415/408}$ for 3C129 -- Our theoretical curve is the solid line.

conservation of particles during expansion yields the following expression for N_0 , the coefficient in the power-law energy distribution of electrons ($N = N_0 E^{-\gamma}$):

$$N_0 = N_0^1 \left(\frac{r}{r_0}\right)^{-\gamma-2} e^{(\gamma-1)\kappa(r_1^{-2/3} - r^{-2/3})}. \quad (14)$$

Assuming conservation of magnetic flux during the expansion we obtain for an optically thin source the following expression for the flux density at a frequency below that of synchrotron cut-off, ν_c :

$$F_\nu \propto F_\nu^1 z^{-2\gamma} e^{(\gamma-1)\tilde{\kappa}(1-z^{-2/3})} \quad (15)$$

where $\tilde{\kappa} \equiv \kappa r_1^{-2/3}$ and $z \equiv r/r_1$. At the end of Region I $z \equiv z_2 \approx 3$ and the flux density evidences a decrease by a factor of about 7.5 from its maximum value for 3C129. We can, by solving Equation (15) for the quantity $\tilde{\kappa}$, estimate the efficiency of turbulent acceleration in an expanding blob: $\tilde{\kappa} = 3.51$ and, therefore $\tilde{f} \equiv f r_0^{-5/3} = 3.68 \times 10^{-15}$.

The observations (Miley 1973) indicate a value near 0.7 for $1415/\nu_c$ at $z = 3$. Using this number and the values of κ and f obtained from the flux density changes, we can estimate the maximum magnetic field of the blob by using Equation (11). We get

$$\frac{\tilde{f}^2}{H_{\max}^3} = 3.4 \times 10^{-16}; \quad (16)$$

with $\tilde{f} = 3.7 \times 10^{-15}$ we obtain $H_{\max} = H_1 = 1.2 \times 10^{-5}$ Gauss.

Thus H is, for most of Region I, larger than the effective field of 3°K blackbody radiation justifying a posteriori the neglect of inverse Compton losses.

As discussed previously, the maximum value of the magnetic field should be such that magnetic and turbulent kinetic energies are near equipartition. Therefore, the internal pressure of the blob should be $\sim 2 H_i^2 / 8\pi$. The condition that initially the blob is confined by ram pressure, i.e.,

$$2 H_i^2 / 8\pi = \rho_{ic} V_o^2 \quad (17)$$

leads directly to the value $\rho_{ic} \approx 9.1 \times 10^{-29} \text{ g/cm}^3$.

Jaffe and Perola's (1973) and Blake's (1972) equation describing the fact that the blob is decelerated by ram pressure is:

$$M_b = D 2\pi r_1^2 \rho_{ic} \quad (18)$$

This implies the mass of the blob in 3C129; $M_b \approx 3.5 \times 10^{39} \text{ g}$, and the density in the fully expanded blob, ρ_{b2} , is $\rho_{b2} \approx 7 \times 10^{-29} \text{ g/cm}^3$.

We should like to comment now on the behavior of polarization of the radiation in the tail with the distance x along the tail. Around x_1 , where the turbulence is well-developed, the polarization is very low. It increases in the region of decaying turbulence (Region II) and it reaches high values in Region III, where no effects of turbulence are present. The magnetic field becomes ordered

by elongation due to the anisotropic effect of ram pressure. This results in the magnetic field acquiring predominant direction along the tail. Knowing the blob's density and magnetic field at distances larger than x_2 , one can compute Faraday rotation measure ($\chi_1 = 2.4 \times 10^{-14} \rho_2 H_2 r_2 / m_H$; Pacholczyk 1970); at 1415 MHz it is $\chi_F = 3.7 \times 10^{-2}$ radians.

We may make several consistency checks on our model:

1. As mentioned in the discussion of the model, $\frac{\rho_b}{\rho_{ic}} \approx D/h \approx 16$ from observations. We now have $\frac{\rho_{ic}}{\rho_b} \text{ calc} = 1.3 \frac{\rho_{ic}}{\rho_b} \text{ obs}$ in good agreement.
2. As was mentioned earlier, it has been claimed (e.g., Kraichnan and Nagarajan 1967 and Longair et al. 1973) that turbulence should amplify an ambient magnetic field to a value near equipartition with turbulent kinetic energy. Thus we should expect $H_1 \lesssim H_{eq}$ where $H_{eq} \approx \sqrt{\rho v_O^2 \cdot 8\pi} \approx 7.2 \times 10^{-5}$ Gauss. As we have seen $H_1 \approx 1.2 \times 10^{-5}$ Gauss $\approx \frac{1}{6} H_{eq}$.
3. Using the e-folding time τ_{amp} calculated earlier, we extrapolate backward from H_1 to find the original field of emerging blob, H_O ; i.e., $H_O = H_1 e^{-x_1/V_g \tau} \approx .75 \times 10^{-6}$ Gauss. Thus $R_{eff} \approx 10^4$ in the emerging blobs and the assertion that turbulence is formed in the blobs leads to a self-consistent model.
4. Our values for the magnetic field, H_2 , and blob density, ρ_2 , in the fully expanded tail should be consistent with no Faraday depolarization since in

Region III, the polarization is $\sim 68\%$. The Faraday rotation of 3.7×10^{-2} radians at 1410 MHz is consistent with the above discussion. It is important to note that this check clearly favors our value $\rho_2 \sim 10^{-29} \text{ gm/cm}^3$ over Jaffe and Perola's (1973) value $\rho_2 \sim 10^{-27} \text{ gm/cm}^3$.

5. All physical parameters were determined by the fitting at the border of Regions I and II. These parameters determine the dependence of $(1415/\nu_c)_{\text{III}}$ on x . As mentioned above, the fitting of our equation for $(1415/\nu_c)$ to the observations between the points x_3 and x_4 , the last observed point, yields $(1415/\nu_c)_{x_4} (\text{predicted}) = 1.1 (1415/\nu_c)_{x_4} (\text{observed})$. See Fig. 4.
6. It is also interesting to point out that McKee (1975) has suggested that the curvature seen in many of the radio tails is due to buoyancy. Our values indeed give $\rho_2 < \rho_{ic}$, the necessary condition for buoyancy. From Equation (18) we see that this ratio depends only on the observed quantities D and r_1 .
7. The energy (per unit volume) acquired by the particles in the turbulent acceleration is of the order of 9×10^{-12} . This constitutes about 2% of the energy of the forward motion of the blob.

CHAPTER 4

X-RAY CLUSTERS

The existence of the radio tail galaxies is relevant to the problem of the origin of the diffuse X-ray emission in clusters of galaxies. Gursky (1973) has pointed out that all models of radio tail galaxies invoke the presence of an intracluster gas. Since the intracluster gas is an obvious requisite of the thermal bremsstrahlung model of X-ray emission, Gursky believes the existence of radio tail systems favors the bremsstrahlung model. Our analysis of the observations of 3C129 indicates an intracluster gas density $\rho_{ic} \approx 7 \times 10^{-29}$, a value which is more than an order of magnitude smaller than the typical densities needed to explain the cluster X-rays in terms of a bremsstrahlung model. Both because of the inaccuracies in the observations and the fitting and because the cluster containing 3C129 is a weaker than average X-ray source (Gursky et al. 1972a, 1972b; Kellogg and Murray 1974), the model is not precluded.

We point out, however, that our model of radio tails qualitatively strengthens the inverse Compton model of X-ray emission. Two points of discussion which previously have been thought to be unexplainable or even inconsistent with

the inverse Compton model are explainable in light of our evidence for particle acceleration in clusters of galaxies.

In the original formulation of the inverse Compton model, Brecher and Burbidge (1972) proposed that the electrons which produce both the diffuse radio and diffuse X-ray sources in clusters leak into the intracluster medium from active radio galaxies. This picture received support from the discovery of a strong correlation between clusters which are X-ray sources and clusters which contain active galaxies (e.g., Owen 1974). Harris and Romanishin (1974) however, have pointed out that the energies of typical electrons found in radio galaxies are about an order of magnitude smaller than the electron energies needed to explain both radio and X-ray observations of the cluster. We feel that the evidence for particle acceleration in radio tails presented here is quite strong, and provides a plausible explanation of this apparent electron energy problem. It is noteworthy that many of the X-ray clusters contain radio tail galaxies: e.g., the Perseus cluster (3 tails), the Coma cluster, the obscured cluster containing 3C129 (2 tails), and Zw2247 + 11. Six of the seven acknowledged tail sources lie in X-ray clusters. Furthermore, it is important to note that this mechanism can be applied to any radio source which expels magnetized blobs of plasma. Thus, in order to be an X-ray source, a cluster need not contain a radio tail, but could produce typical X-ray

emission if it contained a canonical double lobed radio source.

Solinger and Tucker (1972) have reported a strong dependence of cluster X-ray luminosity on velocity dispersion of the galaxies in the cluster. This too was interpreted as being supportive of the thermal bremsstrahlung model. It is clear, however, that the larger the velocity dispersion of a cluster, the higher the probability that an individual active galaxy will be moving at a faster velocity relative to the intracluster gas. Larger V_g implies more efficient acceleration (both betatron and Fermi), and the number of electrons in a given energy band will be a strongly increasing function of the galaxy velocity. Thus the inverse Compton X-ray luminosity in a given frequency range will also be a strong function of the cluster velocity dispersion. The flux emitted in a given frequency band F_ν (I.C.), is proportional to the number of electrons present capable of emission. Thus F_ν (I.C.) $\propto N \propto e v^2$ (from Equations [14] and [3]), a strong dependence.

As more radio data on tailed sources become available, we hope to apply our model to a variety of these interesting objects in order to derive ρ_{ic} for several clusters. If we obtain fittings for a significant number of clusters, this process may help to decide the question of the origin of X-rays in clusters of galaxies.

CHAPTER 5

SUMMARY

We have shown that it is very likely that blobs ejected from active galaxies of the type which give rise to radio tails become turbulent through Rayleigh-Taylor and Kelvin-Helmholz instabilities. Assumptions to the contrary result in contradictions, and we have therefore investigated the following model:

The turbulence, through amplification of the blob magnetic field accelerates particles by the betatron mechanism. Further acceleration of the particles is the result of 2nd order Fermi acceleration, with the turbulent vortices acting as magnetic scattering centers. The electrons also suffer adiabatic and synchrotron losses.

Calculations of the behavior of the spectral index of the synchrotron radiation emitted by the electrons undergoing the above processes result in a predicted spectral index pattern which matches well the observed pattern in several radio tails. Specific fitting of the spectral index, flux, geometry, and polarization data for 3C129 result in a derivation of the plasma parameters in the blobs, and in the density of the intracluster gas.

The value of the intracluster gas density derived here ($\sim 10^{-28}$ g/cm³) is an order of magnitude smaller than densities required to produce typical cluster X-ray sources by the thermal bremsstrahlung mechanism. It is also pointed out that the fact that acceleration of electrons takes place outside the active galaxy is qualitatively supportive of the inverse Compton model of cluster X-ray emission.

It is hoped that further studies of a larger number of tailed sources will be helpful in determining the nature of the cluster X-ray sources.

It is also interesting to note that this type of mechanism may have applications to the more general problem of double lobed radio sources. The instabilities encountered in the radio tail blobs also appear in the plasma ejected from radio galaxies into the intergalactic medium (if such a medium exists and is responsible for lobe confinement; Blake 1972). The resultant in situ acceleration of electrons may help resolve the perplexing problem of the unacceptably large adiabatic losses which are thought to occur in the lobes (Longair et al. 1973).

APPENDIX A

DERIVATION OF ACCELERATION COEFFICIENT

The derivation of the dependence of the acceleration coefficient, $\frac{V_t^2}{cL}$, on blob radius, r , is as follows.

Since L is a mean free path for scattering, $L \propto n^{-1}$ where n is the space density of scattering centers. Thus $L \propto r^3$. Assuming that the turbulent cells form a spectrum which approximates a Kolmagorov spectrum (cf., Batchelor 1967), $\frac{V_t^3}{L} \propto \frac{V_B^3}{r}$. But V_B is a constant ($\sim V_O$) and therefore $\frac{V_t^3}{L} \propto r^{-1}$. Therefore $V_t \propto r^{2/3}$, and $\frac{V_t^2}{cL} \propto r^{-5/3}$; i.e.,

$$\frac{V_t^2}{cL} = \frac{V_B^2}{cL_1} \left(\frac{r_1}{r}\right)^{5/3}.$$

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