AGN AND GALACTIC SITES OF COSMIC RAY ORIGIN

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ABSTRACT

In the following we review recent progress in our understanding of the physics of energetic particles in our Galaxy, in active galaxies such as starburst galaxies, in active galactic nuclei and in the jets and radio hot spots of powerful radio galaxies and radioloud quasars. We propose that cosmic rays originate mainly in three sites, a) normal supernova explosions into the approximately homogeneous interstellar medium, b) supernova explosions into stellar winds, and c) hot spots of powerful radio galaxies. The predictions based on this proposition have been tested successfully against airshower data over the particle energy range from 10 TeV to 100 EeV. The energy input from supernova explosions in starburst galaxies is so large, as to produce breakouts from the galactic disks, observable most prominently in X-ray and nonthermal radio emission. Normal galaxies are also likely to have winds. The low energy cosmic ray particles carried out by such galactic winds to the intergalactic medium may ionize and heat this medium consistent with all known constraints. Furthermore we propose that radioloud quasars accelerate nuclei to high energies in shocks in the inner sections of their jets, and thus initiate hadronic cascades which produce the γ-ray emission observed with GRO. Radio observations of the Galactic center compact source suggest that there is a jet similar to that in quasars, but very much weaker in its emission; simple modelling suggests that the power carried by this jet is of order similar to the accretion disk luminosity. By analogy we speculate that also radioweak quasars have jets which carry a large proportion of the accretion power, and which accelerate particles to high energy. These energetic particles might serve to heat the inner regions of the host galaxies to produce the strong farinfrared dust emission, and could also initiate hadronic and electromagnetic cascades in the inner accretion disk close to the presumed black hole in the center, to produce a large part of the X-ray emission observed from radioweak AGN.

1. Introduction

Energetic particles can have an observable effect in many astronomical objects, from heating and ionization of interstellar material to TeV γ-rays from Active Galactic Nuclei (AGN). Key to all such arguments and an important test for our physical insight is the origin of the observed cosmic ray particles, their energy range, spectrum and chemical composition. There has been recent progress in these areas and in this review we will describe some of it.

1.1. Galactic Cosmic Rays

There are few well accepted arguments about the origins of cosmic rays (for recent reviews of cosmic ray physics, see, e.g. Drury (1983), Hillas (1984), Blandford & Eichler (1987), Berezhko & Krymsky (1988), and Jones & Ellison (1991): a) The
cosmic rays below about $10^4$ GeV are predominantly due to the explosion of stars into the normal interstellar medium (Lagage & Cesarsky 1983). b) The cosmic rays from near $10^4$ GeV up to the knee, at $5 \times 10^6$ GeV, are very likely predominantly due to explosions of massive stars into their former stellar wind (Völk & Biermann, 1988). The consequences of this concept have been checked by calculating the cosmic ray abundances and comparing them with observations (Silberberg et al. 1990); the comparison suggests that up to the highest energy where abundances are known, this concept successfully explains the data. c) For the energies beyond the knee there has been no consensus; Jokipii & Morfill (1987) argue that a galactic wind termination shock might be able to provide those particles, while Protheroe & Szabo (1992) argue for an extragalactic origin, although in either case the matching of the flux at the knee from two different source populations remains somewhat difficult. Reacceleration of the existing cosmic ray population has been investigated by various researchers, most recently by Wandel et al. (1987), Wandel (1988), and Ip & Axford (1992); in such models the chemical abundances are modified at high energy from the interstellar medium abundances in favour of heavy nuclei. It would be an interesting task to test these models against air shower data. d) For the cosmic rays beyond the ankle at about $3 \times 10^9$ GeV an extragalactic origin is required because of the extremely large gyroradii of such particles.

In all such arguments (also, e.g., Bogdan & Völk, 1983, Drury et al. 1989, Markiewicz et al. 1990) the spectrum of the cosmic rays remains unexplained. And yet the observations of the cosmic rays themselves, and of the nonthermal radioemission from our Galaxy as well as from all other well observed galaxies (Golla 1989) strongly suggests, that in all galactic environments studied carefully, the cosmic rays have a universal spectrum of about $E^{-2.7}$ below the knee at $5 \times 10^6$ GeV (electrons only $< 10$ GeV); direct air shower experiments show the spectrum beyond the knee to well approximated by $E^{-3}$ (Stanev 1992). This overall spectrum is clearly influenced by propagation effects, since particles at different energies have different chances to escape from the disk of the galaxy. It appears to be a reasonable hypothesis to approximate the interstellar turbulence spectrum by a Kolmogorov law, which leads to an interstellar diffusion coefficient proportional to $E^{1/3}$. Such an energy dependence then requires a source spectrum of cosmic rays of approximately $E^{-2.4 \pm 0.1}$ below the knee, and approximately $E^{-2.8 \pm 0.1}$ above the knee. We note, that Ormes & Freier (1978) have argued already for a source spectrum of $E^{-2.3 \pm 0.1}$.

Recently we proposed a fully developed physical concept to account for the cosmic ray particle energies, spectrum and chemical abundances (Biermann 1993, paper CR I; Biermann & Cassinelli 1993, paper CR II; Biermann & Strom 1993, paper CR III; Rachén & Biermann 1993, paper UHE CR I). We suggest, that the cosmic rays observed near earth arise from three sites:

1) Supernova explosions into the nearly homogeneous interstellar medium; these cosmic rays reach particle energies of about $100 Z$ TeV, and have an injection spectrum of near $E^{-2.4}$, where $Z$ is the charge of the nucleus considered.

2) Supernova explosions into a former stellar wind produce a slightly flatter spectrum of near $E^{-2.3}$ at injection below a bend near $700 Z$ TeV, beyond which
the spectrum steepens to near $E^{-2.7}$ up to a maximum energy of about 100 Z PeV. This means that heavy nuclei dominate at large particle energies. Since here the acceleration is in a shock that runs through the wind of an evolved star, the heavy elements are enriched (Silberberg et al. 1990). This contribution may dominate even at low energies for some elements. Air shower data have been successfully fitted with this model by Stanev et al. (1993, paper CR IV).

3) The hot spots of powerful radio galaxies provide the third site of origin, to be discussed in the next subsection.

We emphasize that the proposal rests on a plausible but nevertheless speculative assumption about the nature of the transport of energetic particles in perpendicular shock waves, namely that there are large fast convective motions across the shock interface; this notion is, however, supported by radio polarization observations. The model has *predictive power*. The predictions for the galactic cosmic rays underwent their most severe tests with the airshower data in paper CR IV, and the predictions for the extragalactic component were tested with airshower data in Rachen et al. (1993, paper UHE CR II). Cosmic ray acceleration and propagation is also discussed by Jokipii in chapter 5 of this book.

We will describe the essential features of the theory for the Galactic contribution in sections 2.1 and 2.2. In section 2.3 we will apply it to starburst and normal galaxies.

1.2. *Extragalactic Cosmic Rays*

The third site of cosmic ray production is the hot spots of powerful radio galaxies, as proposed already in 1990 (Biermann 1993) at the celebration meeting for M.M. Shapiro in Washington, DC. More detailed models were presented at various meetings, at Kofu in 1990 (Biermann 1991), and at Bartol in 1991 (Rachen & Biermann 1992). The extensive model prediction is in Rachen (1992) and Rachen & Biermann (1993, paper UHE CR I). The full tests against the new airshower data from Fly’s Eye (Gaisser et al. 1993) are presented in Rachen et al. (1993, paper UHE CR II).

The relevant observations are discussed in chapter 6 of this book by Sokolsky. We will discuss the theory and its test in section 3.

1.3. *Radioloud quasars*

Radioloud quasars have long been recognized to provide a unique laboratory for studying high energy particle interactions, since their emission across nearly the entire electromagnetic spectrum could readily be interpreted as synchrotron and inverse Compton emission from particles presumably accelerated in shock waves in relativistic jets. The spatial configuration of these shocks appears to be a double helix in both radio and HST images at high spatial resolution (see, e.g., Macchetto et al. 1991); this appearance is even suggested in radiojets with apparent superluminal motion of the various features (e.g. Krichbaum et al. 1990). Clearly, this calls into question whether we really have shockwaves or not some other mechanism to enhance radio emission at the location of a double helix throughout a jet. Radio quasars with a flat
radio spectrum \( i.e. \) with flux density \( S_\nu \sim \nu^\alpha \), and the spectral index \( \alpha \) being near zero) have been recognized to be a most interesting observationally defined class of radio quasars; in a complete sample of flat spectrum radio quasars all objects have been demonstrated to show evidence for bulk relativistic motion (Witzel et al. 1988).

The new GRO observations of flat spectrum radio quasars demonstrate (Hartman et al. 1992a, Hermsen et al. 1992, Montigny et al. 1992, Bertsch et al. 1993, and many IAU circulars: Kanbach et al. 1992, Fichtel et al. 1992, Michelson et al. 1992, Hartman et al. 1992b, 1992c, Hunter et al. 1992, Dingus et al. 1993), that they have higher luminosities in the \( \gamma \)-ray range than at any other wavelength, with one object even showing emission near TeV photon energies (Whipple collaboration: Weekes et al. 1992, Punch et al. 1992). In chapters 1 and 4 of this book Fichtel & Thompson and Fegan review \( \gamma \)-ray data and possible interpretations.

We will review the hadronic interaction concept in relativistic jets to explain these observations in section 4. Neutrino observations provide an important test for such concepts and are discussed by Stanev in chapter 7 of this book.

1.4. The Galactic Center and Radioweak Quasars

The Galactic center has been demonstrated by recent mm-VLBI observations to show elongated radio emission akin to a compact jet (Krichbaum et al. 1993); in simple models this jet can be shown to carry nearly as much power as is in the accretion disk luminosity (Falcke et al. 1992, 1993). These arguments will be summarized in section 5.1. For radioweak quasars it has only recently become clear (Chini et al. 1989a) that their farinfrared emission is due to dust, and is thus thermal radiation. The origin of this energy is not clear; several possibilities exist, and one of these possibilities is the hypothesis, that the active nucleus provides a large proportion of its power in the form of energetic particles - \( e.g. \) arising from shocks in a putative jet thus following the physical interpretation of the Galactic center made here - which in turn heat the molecular clouds which contain the dust. Various tests of this hypothesis will be described in section 5.2, including some consequences for the emission from the inner accretion disk (X-rays and neutrinos) as well as the farinfrared emission spectra of quasars.

1.5. Background

The following review relies heavily on the recent work of my graduate students, postdocs and collaborators elsewhere, especially the following papers: 1. Biermann 1993, paper CR I, 2. Biermann & Cassinelli 1993, paper CR II, 3. Biermann & Strom 1993, paper CR III, 4. Stanev, Biermann, Gaisser 1993, paper CR IV, 5. Rachen & Biermann 1993, paper UHE CR I, 6. Rachen, Stanev, Biermann 1993, paper UHE CR II, 7. Mannheim 1993a, Proton Blazar, 8. Mannheim 1993b, 3C273, 9. Nellen, Mannheim, Biermann 1993, RQQ.Neutrinos, 10. Falcke et al. 1993 and Falcke, Mannheim, Biermann 1992, Sgr A*, 11. Niemeyer & Biermann 1993, RQQ.FIR, 12. Linden, Duschl, Biermann 1993, and Biermann, Duschl, Linden 1993, G.C. Mol.clouds, 13. Nath & Biermann 1993, IGM, as well as further work.

1.6. History and next steps
Many of the essential conceptual elements used in the following have a long history: After the discovery of cosmic rays by Hess (1912) and Kohlhörster (1913), Baade & Zwicky (1934) already proposed that supernova explosions produce cosmic rays; this viewpoint is upheld in the work reported here up to particle energies of about $3 \times 10^9$ GeV. Alfvén (1939) argued early for a local origin in our Galaxy, a viewpoint which is confirmed by the age determinations of the cosmic rays (Garcia-Munoz et al. 1977). Fermi (1949, 1954) proposed the basic concept of acceleration still being used. Shklovsky (1953) and Ginzburg (1953) made a convincing case for particle acceleration in supernova remnants, again a notion being used in the following, in somewhat modified form. Ginzburg (1953) already emphasized the interesting role of novae; and indeed, the nova GK Per is an important test case. Hayakawa (1956) proposed that stellar evolution gives rise to an enhancement of heavy elements and pointed out the importance of spallation in the interstellar medium; and again, the enrichment found in the cosmic ray contribution from supernova explosions into winds is due to just this enrichment. And finally, Cocconi (1956) already argued convincingly that the most energetic cosmic ray particles are from an extragalactic origin, an argument made quantitative here.

There are many obvious tasks to be done next; we will outline 12 specific next steps that appear feasible in the near future. Additional work may take much longer to follow through with checking the proposal made here in all details and also to produce comprehensive alternatives.

2. Galactic Cosmic Rays

When discussing galactic cosmic rays, the first question to ask is: What is the effect of the leakage out from the galactic disk on the spectrum? There is ample evidence from a variety of sources that the turbulence spectrum in the interstellar medium is of a Kolmogorov character. As a consequence the transport coefficient for cosmic ray particles out of the galactic disk has a $E^{1/3}$ dependence.

First, if one assumes that the turbulence in the interstellar medium responsible for scattering cosmic rays has a single powerlaw over the entire range of energies observed, then we have the limit that even at the highest particle energies, the transit time through the disk should be larger than the simple light crossing time, because we observe no anisotropy. This means that the powerlaw in leakage time vs. particle energy dependence is weaker than 0.4 in the exponent.

Second, the Chicago data already by themselves suggest that over a more limited particle energy range the exponent is close to 1/3 (Swordy et al. 1990).

Third, all evidence from the interstellar medium itself (e.g. Larson 1979, 1981, Jokipii 1988, Narayan 1988, Rickett 1990) suggests that the turbulence is of a Kolmogorov character under the important proviso that the interstellar medium can be thought of as various turbulent phases each with a power law turbulence spectrum: Here, the most definitive measurements come from Very Long Baseline Interferometry (VLBI) at radiowavelength; with this technique one can observe water masers, pulsars and quasars (Blandford et al. 1986, Gwinn et al. 1988a, 1988b; Spangler & Gwinn...
1990; Britzen 1993) and always finds properties of the interstellar medium consistent with the Kolmogorov spectrum.

Fourth, and finally, detailed plasma simulations demonstrate that in the regime where the magnetic field is of a strength similar to the thermal pressure the natural turbulence scaling is indeed Kolmogorov (Matthaeus & Zhou 1989), and only in the limit of strongly dominating magnetic field do we get a different behaviour (the Kraichnan law).

Summarizing, it appears clear that a large body of evidence points to a Kolmogorov law.

However, there is a major difficulty: The secondary products of nucleus-nucleus collisions have a spectrum compared to the primaries which in a simple stationary one zone leaky box model for the galactic disk suggests a law for the effective turbulence which is different. Such an analysis suggests that the energy dependence of the leakage time is closer to 2/3 rather than 1/3. However, the data on giant molecular clouds, where surely a large part of the cosmic ray interaction happens, suggest that simple stationary one zone models are not adequate. Large molecular clouds assemble out of bits of small clouds in spiral arms and then slowly disperse again; the time scales for assembly and dispersal are comparable to the time scale of cosmic ray leakage from the galactic disk, making it obvious that the interaction of cosmic rays with clouds in their structure and their slow dispersal and diffusion out of clouds has to be investigated in order to answer the question what the secondary to primary nuclei ratio ought to be as a function of particle energy. This problem is not solved at present, and remains to be worked out; this constitutes a specific next step.

We will assume in the following that the correction from observed cosmic ray spectrum to source spectrum is a change in spectral slope by exactly 1/3, so that we are looking at source spectra in the range \( E^{-2.4 \pm 0.1} \) below the knee, and \( E^{-2.8 \pm 0.1} \) above the knee.

In the following we will first discuss the explosions into the interstellar medium in section 2.1 and then the explosions into a former stellar wind, in section 2.2. Generalizations for a discussion of starburst and normal galaxies follow in section 2.3.

2.1. Supernova explosions into the interstellar medium

Already for some time it has been argued that normal supernova explosions provide the lower particle energy range of the cosmic rays observed (e.g. Lagage & Cesarsky 1983). However, there were several open questions:

1) The standard diffusive particle acceleration at shocks (Drury 1983) works only in the approximation that the shocks are parallel, i.e. that the shock normal is parallel to the prevailing direction of the magnetic field. This is, however, normally not true, and already Jokipii (1987) argued that this at the very least gives an acceleration time very much shorter than the standard approach. The problem of perpendicular shocks remained.

2) The spectrum of the cosmic rays at source could not be derived reliably from
physical arguments; this problem stems partially from the lack of understanding of
the spectral correction between source and observer (here we use $1/3$, see above).

3) What is the maximum particle energy that can be reached?

4) Why do we have at low particle energies such a large discrepancy between
the Hydrogen and Helium to heavy nuclei ratio in the cosmic rays compared to the
interstellar medium, out of which these particles are presumed to be accelerated?

We propose to discuss here the acceleration of particles in the blastwaves
caused by supernovae in the normal interstellar medium.

Observational arguments suggest that the simple Sedov solution is adequate
(Biermann & Strom 1993, paper CR III); we will use the standard Sedov solution:
The Sedov similarity solution for an explosion into a homogeneous medium can be
written as a function of shock radius $r$ and shock speed $U_1$; as a function of time $t_4$,
in units of $10^4$ years; explosion energy $E_{51}$, in units of $10^{51}$ ergs; and density of the
interstellar medium $n_o$ in units of $\text{cm}^{-3}$. This gives (Cox 1972)

$$ r = 13.7 t_4^{2/5} (E_{51}/n_o)^{1/5} \text{pc} \quad (1) $$

and

$$ U_1/c = 1.8 \times 10^{-3} t_4^{-3/5} (E_{51}/n_o)^{1/5} 
= 0.090 \left( \frac{r}{\text{pc}} \right)^{-3/2} (E_{51}/n_o)^{1/2}. \quad (2) $$

2.1.1. Particle Acceleration

Shock acceleration in its standard form just takes the Lorentz transformation
for a energetic particle at a velocity $v$ much larger than the shock velocity as it
goes between scatterings in a weakly turbulent magnetic field from downstream to
upstream and back. In practice we will assume $v$ to be close to $c$. One such cycle
gives an momentum $p$ gain of

$$ \left( \frac{\Delta p}{p} \right)_{LT} = \frac{4}{3} \frac{U_1}{v} \left( 1 - \frac{U_2}{U_1} \right). \quad (3) $$

We assume the shock to be subrelativistic and so the phase space distribution
of the particles to be nearly isotropic. Then downstream (see Drury 1983 for the exact
derivation) the particles have a finite chance to escape. They cross the shock with
with the flux $n v/4$, where $n$ is the particle density, $v$ their velocity, and the factor 1/4
comes from $1/2$ of all particles moving in one direction with a velocity component on
average $1/2$ of their absolute velocity parallel to the shock normal. Far downstream
the particle flux is, however, $n U_2$, where $U_2$ is the downstream flow velocity in the
frame of the shock. The ratio $4 U_2/ v = U_1/ v$ is the chance to escape from the shock
region. This probability is per cycle. The combination yields a powerlaw for the
distribution $p^2 f(p)$, where $f(p)$ is the particle distribution function in phase space,
with the powerlaw index
However, this assumes that there are no other energy losses or energy gains during a cycle of a particle going back and forth between upstream and downstream. In a configuration, where there is a component of the magnetic field perpendicular to the shock normal, there can be energy gains by drifts parallel to the electric field seen by a particle moving with the shock system, and also losses due to adiabatic expansion in the expansion of a curved shock. Since drift acceleration is a rate, the resulting particle energy gain is proportional to the time spent on either side of the shock. Furthermore, there can be spectral changes due to the fact, that particles were injected at a different rate in the past, when some particle under consideration now was injected; again, this is likely to happen in a spherical shock.

Thus, the key here is to know what the time scale of a cycle actually is, or in other words, what the transport or diffusion coefficient is for the particles to move from one side of the shock to the other. Both energy losses and energy gains (over and above the Lorentz transformation) are proportional to the time scale of a cycle. This we will discuss therefore in a separate section next.

2.1.2. The basic Hypothesis

The radio observations of young supernova remnants can be a guide here. The explosions into the interstellar medium are explosions where the unperturbed magnetic field is not far from perpendicular to the shock direction over most of $4\pi$ steradians. Radio polarization observations of supernova remnants clearly indicate what the typical local structure of these shocked plasmas is. The observational evidence (Milne 1971, Downs & Thompson 1972, Reynolds & Gilmore 1986, Milne 1987, Dickel et al. 1988) has been summarized by Dickel et al. (1991) in the statement that all shell type supernova remnants less than 1000 years old show dominant radial structure in their magnetic fields near their boundaries. There are several possible ways to explain this: Rayleigh-Taylor instabilities between ejected and swept up material can lead to locally radial differential motion and so produce a locally radial magnetic field (Gull 1973). It could be due to strong radial velocity gradients of various clumps of ejecta, or result from clouds being overrun that now evaporate and cool the surrounding material (all of which are mentioned by Chevalier [1977]). It may seem paradoxical that we will be considering perpendicular shocks where the observations show a parallel magnetic field, but it must be remembered that the observed field is a superposition of the configurations of all emitting regions in the telescope beam. The degree of polarization in all of the objects referred to above is low (typically 10% or less), much less than that expected from a perfectly uniform field (generally 65%), which is consistent with there being both parallel and perpendicular components present.

The radial magnetic field configuration is also found in a few shell remnants somewhat older than 1000 yr, such as RCW 86 (SN 185) and Pup A (with a kinematic age of some 3700 yr, Winkler et al. 1988). However, in most “mature” remnants the
field is tangential, or complex (Milne 1987). While a radially-directed magnetic field provides the most compelling evidence for large-scale differential motions parallel to the shock direction, there are other indications as well. It is generally believed that the oxygen-rich fast moving knots (fmk's) in Cas A are ejecta from the progenitor (Kamper & van den Bergh 1976). Oxygen-rich knots have also been seen in Pup A (Winkler et al. 1988) and the similarly aged G 292.0+1.8 (Braun et al. 1986), and may also be present in a number of extragalactic SNRs (Winkler and Kirshner 1985). Radially-directed radio structures have been mapped in Cas A, and their motion has also been determined (Braun et al. 1987). Recent Rosat observations of the older (10000 yr) Vela SNR show similar features in the X-ray emission which have apparently penetrated the main SNR shell (seminar by Trümper at the Texas conference December 1992), and are believed to result from fmk-like ejecta. (Their very soft X-ray spectra would make similar features difficult to detect in more distant remnants.)

These examples are for supernova explosions into the interstellar medium; there is also an observation demonstrating the same effect for an explosion into a wind: Seaquist et al. (1989) find for the spatially resolved shell of the nova GK Per a radially oriented magnetic field in the shell, while the overall dependence of the magnetic field on radial distance \( r \) is deduced to be \( 1/r \) just as expected for the tightly wound up magnetic field in a wind. The interpretation given is that the shell is the higher density material behind a shock wave caused by the nova explosion in 1902 and now travelling through a wind. Seaquist et al. (1989) note the similarity to young supernova remnants.

The important conclusion for us here is that there appear to be strong radial differential motions in perpendicular shocks which provide the possibility that particles get convected parallel to the shock direction. We assume this to be a diffusive process, and note that others have also pointed out that this may be a key to shock acceleration (e.g. Falle 1990). Our task here will be to derive a natural velocity and a natural length scale, which can be combined to yield a transport coefficient. A classical prescription is the method of Prandtl (1925): In Prandtl's argument an analogy to kinetic gas theory is used to derive a diffusion coefficient from a natural scale and a natural velocity of the system. Despite many weaknesses of this generalization Prandtl's theory has held up remarkably well in many areas of science far beyond the original intent. We will use a similar prescription here.

Consider the structure of a layer shocked by a supernova explosion into a homogeneous medium in the case that the adiabatic index of the gas is \( 5/3 \) and the shock is strong. Then there is an inherent scale in the system, namely the thickness of the shocked layer, in the spherical case for a strong shock \( r/12 \). In a spherical strong shock into a wind this natural length scale is \( r/4 \). There is also a natural velocity scale, namely the velocity difference of the flow with respect to the two sides of the shock. Both are the smallest dominant scale in velocity and in length; we will use the assumption that the smallest dominant scale is the relevant scale several times subsequently to derive diffusion coefficients and other scalings. This choice can be thought of also as follows: The smallest dominant scale in a transport coefficient
provides the longest time scale for the transport mechanism, and it is well understood that in a combination of transport processes the slowest one always wins.

Our basic conjecture, **postulate 1**, based on observational evidence, is that the convective random walk of energetic particles perpendicular to the unperturbed magnetic field can be described by a diffusive process with a downstream diffusion coefficient $\kappa_{rr,2}$ which is given by the thickness of the shocked layer and the velocity difference across the shock, and is independent of energy:

$$
\kappa_{rr,2} = \frac{1}{9} \frac{U_2}{U_1} r (U_1 - U_2)
$$

(5)

The upstream diffusion coefficient can be derived in a similar way, but with a larger scale. We make here the second critical assumption, **postulate 2**, namely that the upstream length scale is just $U_1/U_2$ times larger, and so is $r/3$. This, obviously, is the same ratio as the mass density and the ratio of the gyroradii of the same particle energy. Since the magnetic field is lower by a factor of $U_1/U_2$ upstream, that means that the upstream gyroradius of the maximum energy particle that could be contained in the shocked layer, is also $r/3$. Hence $r/3$ is the natural scale. And so the upstream diffusion coefficient is

$$
\kappa_{rr,1} = \frac{1}{9} r (U_1 - U_2)
$$

(6)

For these diffusion coefficients, it follows that the residence times (Drury 1983) on both sides of the shock are equal and are

$$
\frac{4 \kappa_{rr,1}}{U_1 c} = \frac{4 \kappa_{rr,2}}{U_2 c} = \frac{4}{9} \frac{r}{c} \frac{U_2}{U_1} (1 - \frac{U_2}{U_1})
$$

(7)

Adiabatic losses then cannot limit the energy reached by any particle since their time scale is proportional to the acceleration time, both being independent of energy. So the limiting size of the shocked layer limits the energy that can be reached to that energy for which the gyroradius just equals the thickness of the shocked layer, provided the particles can reach this energy. We assume here that the average of the magnetic field $\langle B \rangle$ is not changed very much by all this convective motion, but leave the possibility open that the root mean square magnetic field $\langle B^2 \rangle^{1/2}$ is increased; this implies that a magnetic dynamo does not work as fast as the timescales given by the shock (see, e.g., Galloway & Proctor 1992 for arguments on the shortest possible dynamo time scale). This then leads to a possible maximum energy of

$$
E_{max} = \frac{1}{3} \frac{U_2}{U_1} Z e B_2 = \frac{1}{3} Z e B_1
$$

(8)

where $Ze$ is the particle charge and $B_{1,2}$ is the magnetic field strength on the two sides of the shock. This means, once again, that the energy reached corresponds to the maximum gyroradius the system will allow. It also means that we push the diffusive picture right up to its limit where the diffusive scale becomes equal to the mean free path and the gyroradius of the most energetic particles. For supernova explosions
into stellar winds the corresponding argument produces the limiting condition, while for supernova explosions into the interstellar medium we will derive a more restrictive condition below.

2.1.3. The Particle Spectrum

Consider particles which are either upstream of the shock, or downstream; as long as the gyrocenter is upstream we will consider the particle to be there, and similarly downstream.

In general, the energy gain of the particles will be governed primarily by their adiabatic motion in the electric and magnetic fields. The expression for the energy gain is well known and given in, e.g., Northrop (1963, equation 1.79), for an isotropic angular distribution where a first term arises from the drifts and a second from the induced electric field. The expression is valid in any coordinate frame. We explicitly work in the shock frame, separate the two terms and consider the drift term first. The second term is accounted for further below.

The \( \theta \)-drift can be understood as arising from the asymmetric component of the diffusion tensor, the \( \theta r \)-component. The natural scales there are the gyroradius and the speed of light, and so we note that for (Forman et al. 1974)

\[
\kappa_{\theta r} = \frac{1}{3} r_g c,
\]

the exact limiting form derived from ensemble averaging, we obtain the drift velocity by taking the proper covariant divergence (Jokipii et al. 1977); this is not simply (spherical coordinates) the \( r \)-derivative of \( \kappa_{\theta r} \). The general drift velocity is given by (see, e.g., Jokipii 1987)

\[
V_{d,\theta} = \frac{c E}{3Ze} \frac{\text{curl}_\theta B}{B^2}.
\]

The \( \theta \)-drift velocity is thus zero upstream and only finite downstream due to curvature:

\[
V_{d,\theta} = \frac{1}{3} c r_g \frac{r_g}{r},
\]

where \( r_g \) is now taken to be positive.

It must be remembered that there is a lot of convective turbulence which increases the curvature: The characteristic scale of the turbulence is \( r/12 \) for strong shocks, and thus the curvature is \( 12/r \) maximum. Taking half the maximum as average we obtain then for the curvature \( 6/r \) which is six times the curvature without any turbulence; this increases the curvature term by a factor of six thus changing its contribution from \( 1/3 \) to \( 2 \) in the numerical factor in the expression above. Hence the total drift velocity is thus

\[
V_{d,\theta} = \frac{1}{2} \left( \frac{U_1}{U_2} \right) c r_g \frac{r_g}{r},
\]
now written for arbitrary shock strength. It is easily verified that the factor in front
is two for strong shocks where \( U_1/U_2 = 4 \). With \( \Delta E_1 = 0 \) we have then downstream

\[
\frac{\Delta E_2}{E} = \frac{2 U_1}{c} \left( 1 - \frac{U_2}{U_1} \right)
\]

(13)

which is the total energy gain from drifts.

Next, we consider the influence of adiabatic losses and the injection history on
the particle spectrum.

Let us consider then one full cycle of a particle remaining near the shock and
cycling back and forth from upstream to downstream and back. Adding the energy
gain in the relativistic approximation due to drifts to the energy gain from the simple
Lorentz transformation we obtain

\[
\frac{\Delta E}{E} = \frac{4 U_1}{3 c} \left( 1 - \frac{U_2}{U_1} \right) x
\]

(14)

where

\[
x = 1 + \frac{1}{6}.
\]

(15)

Consider how long it takes a particle to reach a certain energy:

\[
\frac{dt}{dE} = \left\{ 8 \frac{\kappa_{rr,1}}{U_1 c} \right\} \left\{ \frac{4 U_1}{3 c} \left( 1 - \frac{U_2}{U_1} \right) x E \right\}.
\]

(16)

Here we have used \( \kappa_{rr,1}/U_1 = \kappa_{rr,2}/U_2 \). Since we have

\[
r = \frac{5}{2} U_1 t
\]

(17)

this leads to

\[
\frac{dt}{t} = \frac{dE}{E} \frac{3 U_1}{U_1 - U_2} \frac{2 \kappa_{rr,1}}{x r U_1} \frac{5}{2}
\]

(18)

and so to a dependence of

\[
t(E) = t_o \left( \frac{E}{E_o} \right)^\beta
\]

(19)

with

\[
\beta = \frac{3 U_1}{U_1 - U_2} \frac{2 \kappa_{rr,1}}{x r U_1} \frac{5}{2}
\]

(20)

Particles that were injected some time \( t \) ago were injected at a different rate,
say, proportional to \( r^b \). Also, in a \( d \)-dimensional space, particles have \( \sim r^d \) more space
available to them than when they were injected, and so we have a correction factor
which is

\[
\left( \frac{E}{E_o} \right)^{-\frac{2}{3} (b+d) \beta}.
\]

(21)
The combined effect is a spectral change by

$$\frac{3U_1}{U_1 - U_2} \frac{2}{x} (d + b) \frac{\kappa_{rr,1}}{r U_1^2}. \quad (22)$$

We note that the factor $2/5$ from the Sedov expansion drops out again. For a detailed argument on the sign conventions used here, see paper CR I. This expression can be compared with a limiting expansion derived by Drury (1983; eq. 3.58), who also allowed for a velocity field; Drury (1983) generalized earlier work on spherical shocks by Krymskii & Petukhov (1980) and Prishchep & Ptuskin (1981). Drury's expression agrees with the more generally derived expression given here for $x = 1$. The comparison with Drury's work clarifies that for $\kappa \sim r$ the inherent time dependence drops out except, obviously, for the highest energy particles, discussed further below; the same comparison shows that the statistics of the process are properly taken into account in our simplified treatment. The $r^d$-term describes the adiabatic losses from the linear expansion of the shock layer in their effect on the spectrum, and thus accounts for the second term in Northrop's expression mentioned above. This also describes then the adiabatic losses, due to the general expansion of the shock layer. Hence the total spectral difference, as compared with the plane parallel case, is given by

$$\frac{3U_1}{U_1 - U_2} \left\{ \frac{U_2}{U_1^2} \left( \frac{1}{x} - 1 \right) + \frac{2}{x} (b + d) \frac{\kappa_{rr,1}}{r U_1^2} \right\}. \quad (23)$$

If this expression is positive the spectral index is steeper.

Furthermore, we have to discuss the effect of the expansion of the material in the shell, in this context already introduced by Drury (1983). This expansion is only that beyond the linear expansion of the overall flow already taken into account. We define for the flow field $U(\xi r)$ as function of radius $\xi r$ interior to the shock of radius $r$ the velocity by

$$U(\xi r) = V(\xi) \xi U_1. \quad (24)$$

As shown by Drury (1983), this velocity field introduces a steepening by

$$\frac{3U_1}{U_1 - U_2} \frac{\kappa_{rr,2}}{r U_2} \frac{V'(1)}{V(1)}, \quad (25)$$

neglecting drifts and in the limit of very small diffusion coefficient. However, in our context where we consider a finite shell, we have to use a non-local approximation (see paper CR III). From a detailed discussion we have an extra term from this postshock adiabatic loss due to the shell expansion of

$$\frac{3U_1}{U_1 - U_2} \frac{1}{x} \frac{\kappa_{rr,2}}{r U_2} \frac{V'(1)}{V(1)} 1.970. \quad (26)$$

which is then slightly less than double the effect which Drury discusses. For a strong shock then this particular effect adds 0.563 to the spectral index.
For strong shocks in stellar winds, for which the natural length scale is \( r/4 \), the linear part of the expansion is three times higher and so the remaining difference to the postshock speed of sound is only \( 2.25 - \sqrt{5} \), which gives an effect of only \( 2 \times 10^{-3} \) in the spectral index, justifying our neglect of this effect in the case of a shock travelling through a stellar wind (see papers CR I and CR II and below).

The exponent \( b \) describes the injection as a power of the radius; in a Sedov solution the injection, assumed to be proportional to \( \rho U^2_1 \), is given by

\[
\rho U^2_1 \sim r^{-3}.
\]

Hence we have \( b = -3 \), so that \( b + d = 0 \). Thus we have for a pure Sedov solution, using the velocity field term introduced above, a spectral difference to the plane parallel case of 0.420 and thus an injection spectrum of

\[
E^{-2.420}.
\]

However, we note that this spectrum is relevant for electrons obviously only if they are injected into the acceleration process at all. This spectrum derived here agrees with the electron spectrum deduced from the radio synchrotron emission in the radio shell of the supernova remnant Cas A. Adding the term from diffusive losses from the Galaxy the final observable spectrum expected is

\[
E^{-2.753},
\]

very close to what is observed, both in Galactic cosmic rays (protons and nuclei), as through the nonthermal radioemission of other galaxies (Golla 1989, electrons).

Clearly, there is uncertainty in our method to treat the non-linear flow field. Further non-linear corrections, such as also averaging the non-linear flow field itself, lead to spectral indices within 0.04 of the answer given above.

2.1.4. Observational Tests

Calculating now the radio emission contribution from the shell of \( r U_2/(3 U_1) \), we note that in the Sedov solution the ram pressure is proportional to \( U^2_1 \), which in turn is proportional to \( r^{-3} \). The total radio emission from the shell is then constant with time, if the efficiency of injecting electrons \( \eta \) is a constant. The luminosity is then given by

\[
L_\nu = 5.4 \times 10^{23} B_{-5.3}^{1.710} \eta^{-1} E_{51}^{1.90} \nu_{10}^{-0.710} \text{erg/sec/Hz}
\]

and so does not depend on the interstellar medium density; since the adiabatic expansion is a condition of constant energy, all such dependencies drop out. We have used here as reference 0.1 for \( \eta \), 5\( \mu \)Gauss for the interstellar unperturbed magnetic field strength, and 1 GHz for the emission frequency observed. We note, that the emission here is coming from a shell of thickness \( r/12 \); additional emission further inside might arise if there is a reverse shock resulting from the transition of free expansion to adiabatic expansion. Furthermore, there is observational evidence that the
synchrotron emission systematically samples the higher magnetic field substructures, by a factor of up to 5 and so we can expect the synchrotron emission to be higher by factors up to 15, leading to an implied radio luminosity of up to $8.5 \times 10^{24} \text{ erg/sec/Hz}$, everything else being equal. This luminosity is in agreement with the upper limit of the distribution of observed luminosities, suggesting that in their initial phases the efficiency of putting kinetic energy into an energetic electron population is indeed of order 0.1. An important next step is to actually model all the selection effects that dominate the observations (see especially the discussion in Berhuijsen 1986).

The evolution of a supernova remnant (see, e.g., for the basic physical concepts employed Cox [1972]) is then is divided into four phases:

First, there is free expansion until the interstellar medium overrun by the explosion shock is of about the same mass as the mass ejected. This happens at a radius of

$$R_e = 1.92 \text{ pc} \left( \frac{M_{ej}}{M_\odot} \right)^{1/3} n_o^{-1/3},$$

where $M_{ej}$ is the ejected stellar mass and $n_o$ is again the environmental density in particles per cc. Numerical simulations demonstrate that the expansion is noticeably slowed down long before this critical radius is reached.

Second, we have a Sedov phase with adiabatic expansion and steady fresh electron injection. The observations of radio emitting stars (see paper CR II) suggest that there is a critical electron injection velocity for a shock, below which electron injection is substantially weakened. Clearly, the critical velocity in the interstellar medium is likely to be quite different from that in stellar winds. We do not wish to dwell on electron injection physics here (the nonthermal radio emission of other stars can be a guide here), but adopt somewhat arbitrarily as a critical velocity in the interstellar medium with an Alfvénic Mach number roughly similar to the massive star wind case, with 1000 km/sec and obtain then

$$R_{crit} = 9.0 \text{ pc} \left( \frac{E_{51}}{n_o} \right)^{1/3} U_{crit,-2.5}^{-2/3}.$$  

An observational comparison of the evidence for the value of the critical velocity in different sites is an important next step, as for instance on the basis of the radio and optical observations of the classical nova GK Per.

Third, when the expansion velocity has decreased below the critical velocity, the electron injection ceases (or, to be more accurate, decreases by many orders of magnitude) and we have further adiabatic expansion, but with the electron population only modified by adiabatic losses (Shklovsky 1968). Adiabatic losses of a relativistic particle decrease its energy by the ratio of the initial and final radius scale.

And finally, fourth, we have the break-up phase due to the various instabilities in the cooling layer (McCray et al. 1975, Chevalier & Imamura 1982, Smith 1989). Clearly, these instabilities do not stop the expansion of the dense layer, but lead to a rather chaotic slowdown of the various shell fragments (see, e.g., the beautiful new color photos of the Cygnus loop from HST and ROSAT circulating in the popular press). The cooling radius is (Cox 1972)
\[ R_{\text{cool}} = 15.6 \, \text{pc} \, E_{51}^{3/11} \left(10^{22} L\right)^{-2/11} n_o^{-5/11}, \quad (33) \]

where \( L \) is the cooling coefficient; the cooling curve has been compared with recent detailed calculations by Schmutzler & Tscharnuter (1993) and is still a very good approximation at the temperature range of interest here. During this time the energy of any individual relativistic electron decreases by adiabatic expansion as the ratio of the radii from the time when injection ceases to the time of break up. Thus the energy density of the electron population (from conservation of the adiabatic moment only and disregarding spatial dilution, see below) decreases by the factor \((p = 2.420)\)

\[ \left(\frac{R_{\text{cool}}}{R_{\text{crit}}}\right)^{p-1} = 2.18 \, E_{51}^{-0.09} n_o^{-0.17} \left(10^{22} L\right)^{-0.26} U_{\text{crit}, -2.5}^{0.95}. \quad (34) \]

showing a very weak dependence on the environmental density \( n_o \). The dependence on the assumed specific value of the critical velocity is very nearly linear; hence the numerical factor in this expression would go down with a smaller value for the critical velocity. The simple spatial dilution of the energy density is the same for both protons and electrons and drops out in the ratio, as long as we are in the Sedov phase (see Shklovsky [1968], eq. 7.27). For a tenuous interstellar medium of density about 0.01 Hydrogen atoms per cc the factor is increased from 2.18 to nearly 4.8.

This intermediate switch from electron injection with steady acceleration to a simple adiabatic loss regime determines the net scaling of the power of the electron population to that of the proton population in the cosmic rays. The observations suggest that from 1 GeV the electron number is only about 0.01 (e.g. Wiebel 1992) of the protons. From this observed ratio the energy density ratio integrated over the entire relativistic part of the particle spectrum, protons relative to electrons, for the spectral index of \(-2.42\), is given by 4.3. Assuming then, that the ratio of protons and electrons is not influenced by propagation effects, this suggests an expansion of a factor in radius of order 3 between the time when electron injection ceases and the time when cooling breaks up the shell of the supernova remnant; here we assume that electrons and protons originally have comparable energy densities of their relativistic particle populations. We may have thus identified the origin of the observed electron/proton ratio in cosmic rays.

The maximum energy particles can possibly reach is given above on spatial arguments, and depends linearly on the magnetic field. Near the symmetry axis we have diffusion parallel to the magnetic field, and, given the strong turbulence induced by the shock we take there the Bohm limit for the upstream diffusion coefficient \( \kappa || = \frac{1}{3} \frac{E}{ZeB_1} c \). At the equator we use our radial upstream diffusion coefficient derived earlier, here for a strong shock, of \( \kappa \perp = \frac{r U_1}{12} \). Combining these two upstream diffusion coefficients (Jokipii 1987) to obtain an acceleration time scale and setting that equal to the upstream flow time scale of \( r/(3U_1) \) then yields a condition on the maximum particle energy that can be reached. This condition is

\[ \frac{r}{3U_1} = \frac{4}{3} \frac{E}{ZeB_1} \frac{c}{U_1^2} \mu^2 + \frac{r}{3U_1} \left(1 - \mu^2\right) \frac{1}{x}. \quad (35) \]
This leads to

$$E_{\text{max}} = \frac{1}{4} Z e B_1 r \frac{U_1}{c}$$

(36)

in the approximation of \(x \simeq 1\), which is true here. In stellar winds, on the other hand, \(x \simeq 2\), and this approximation cannot be used (see paper CR I). A more detailed consideration of the strength of the magnetic field as a function of latitude does not change this result. Since the angular coordinate does not enter this expression, it is valid for any subsection of the sphere, and thus also for an inhomogeneous magnetic field.

Using normal interstellar magnetic fields of 5 \(\mu G\), and a standard Sedov explosion with an arbitrary radius then yields a maximum particle energy. At breakup, when we suggest the actual mixing of the particle population into the interstellar medium begins, these maximum energies are

$$E_{\text{max}}(\text{protons}) = 2.6 \times 10^4 F_{CR}\text{GeV},$$

(37)

and

$$E_{\text{max}}(\text{iron}) = 7.4 \times 10^5 F_{CR}\text{GeV},$$

(38)

where we use

$$F_{CR} = E_{51}^{0.364} (10^{22} L)^{0.091} n_o^{-0.273}.$$

(39)

The protons ought to dominate over heavier nuclei due to the normal abundances in the interstellar medium, and higher energies for protons are clearly reached in the tenuous hot part of the interstellar gas where the densities are of order 0.01 particles/\(\text{cc}\), and so maximum proton energies are possible to about 1.1 \(10^5\) GeV. These numbers are similar to earlier estimates of the maximum particle energy in a Sedov expansion phase of a supernova. Below (see paper CR IV) we obtain from a fit to the air shower data an estimate for this number of 1.2 \(10^5\) GeV, in very good agreement with these arguments, and suggesting that the cosmic ray injection is most effective in the tenuous part of the interstellar medium, an argument which has been made before on quite different grounds.

There is another observational check, and that is on the electron spectrum of cosmic rays. For particle energies above a few GeV, and that corresponds to synchrotron emission at high radio frequencies normally unobservable in galaxies because of increasing thermal dust emission, the observed electron spectrum is about \(E^{-3.3 \pm 0.1}\), consistent with a steepening by unity from injection due to synchrotron losses, in this energy range faster than leakage losses. This confirms our injection spectrum derived. A detailed discussion of the time scales involved remains to be done.
2.2. Supernova Explosions into a Stellar Wind

The basic hypothesis for the second site of origin of cosmic rays is, that we consider explosions into a stellar wind with a Parker spiral topology (Parker 1958). We remind the reader that in such a wind in the asymptotic regime, where the magnetic field decreases with the radial distance $r$ as $1/r$ and the wind velocity is constant, the (tangential) Alfvén velocity is also constant with radial distance.

We also note that the wind in massive stars has a similar energy integrated over the main sequence life time as the subsequent supernova explosion. The wind bubble is large and is itself surrounded by a dense shell. Hence we can expect the shock of the supernova to disperse this shell and to mix the energetic particle population produced in the shock running through the wind directly into the interstellar medium. Thus, there is no additional energy dependence introduced here to go from the spectrum which we calculate below to the injection spectrum of cosmic rays.

Outside the wind acceleration region stellar winds are likely to be similar to the solar wind, and so we will assume a Parker spiral topology of the magnetic field (e.g. Jokipii et al. 1977). This leads then to a tangential drift of

$$V_{d\theta} = \frac{2}{3} c r_g/r,$$  \hspace{1cm} (40)

where $r_g$ is the Larmor radius of the particle under consideration. The drift is towards the equator for $Z B_s > 0$. This drift - both gradient and curvature drift - here is just due to the unperturbed structure of the stellar wind magnetic field; we will consider the consequences of additional curvature from turbulence.

2.2.1. The Particle Spectrum below the Knee

Consider the structure of a layer shocked by a Supernova explosion into a stellar wind in the case, that the adiabatic index of the gas is $5/3$ and the shock is strong. Then there is an inherent length scale in the system, namely the thickness of the shocked layer, in the spherical case for a shock velocity much larger than the wind speed and in the strong shock limit $r/4$. This leads to very similar expressions for diffusion coefficients and respective time scales as derived above for the case of an explosion into a homogeneous medium, just with a factor of 3 larger length scales.

There is an important consequence of this picture for the diffusion laterally: From the residence timescale and the velocity difference across the shock we find a distance which can be traversed in this time of

$$\frac{4}{3} c \frac{r}{U_1} (1 - \frac{U_2}{U_1}) (U_1 - U_2).$$ \hspace{1cm} (41)

Since the convective turbulence in the radial direction also induces motion in the other two directions, with maximum velocity differences of again $U_1 - U_2$, this distance is also the typical lateral length scale. From this scale and again the residence time we can construct an upper limit to the diffusion coefficient in lateral directions of
\( \kappa_{\theta \theta, \text{max}} = \frac{4}{9} (1 - \frac{U_2}{U_1})^3 (\frac{U_1}{c})^2 r c, \)  
(42)

which is for strong shocks equal to

\( \kappa_{\theta \theta, \text{max}} = \frac{1}{3} (\frac{3}{4} \frac{U_1}{c})^2 r c. \)  
(43)

Again in the spirit of the idea, that the smallest dominant scale wins, this then will begin to dominate as soon as the \( \theta \)-diffusion coefficient reaches this maximum at a critical energy. As long as the \( \theta \)-diffusion coefficient is smaller, it will dominate particle transport in \( \theta \) and the upper limit derived here is irrelevant. When the \( \theta \)-diffusion coefficient reaches and passes this maximum, then the particle in its drift will no longer see an increased curvature due to the convective turbulence due to averaging and the part \( (1/3 \text{ of total for strong shocks}) \) of drift acceleration due to increased curvature is eliminated. This then reduces the energy gain, and the spectrum becomes steeper from that energy on. The critical particle energy thus implied will be identified below with the particle energy at the knee of the observed cosmic ray spectrum, and this reduced energy gain from drifts provides the steepening of the spectrum beyond the knee.

Again, as above the turbulence increases the curvature of the local medium, and so increases the drift velocity. The total drift velocity, combining now again the curvature \( (2/3) \) and gradient \( (1/3) \) terms, is then

\[ V_{d, \theta} = \frac{1}{3} (1 + \frac{U_1}{2U_2}) \frac{cr_g}{r}, \]  
(44)

now written for arbitrary shock strength. It is easily verified that the factor in front is unity for strong shocks where \( U_1/U_2 = 4 \).

The energy gain associated with such a drift is given by the product of the drift velocity, the residence time, and the electric field. Upstream this energy gain is given by

\[ \Delta E_1 = \frac{4}{3} E \frac{U_1}{c} f_d (1 - \frac{U_2}{U_1}), \]  
(45)

where

\[ f_d = \frac{1}{3} (1 + \frac{U_1}{2U_2}). \]  
(46)

Thus, \( f_d = 1 \) for strong shocks. The total energy gain is

\[ \frac{\Delta E}{E} = \frac{4}{3} \frac{U_1}{c} f_d (1 + \frac{U_2}{U_1}) (1 - \frac{U_2}{U_1}). \]  
(47)

The drift energy gain averages over the magnetic field strength during the gyromotion. We emphasize that this energy gain is independent of this average magnetic field, so that even variations of the magnetic field strength due to convective motions do not change this energy gain.
Considering again a full cycle of a particle going back and forth across a curved shock produces the following net energy gain:

\[ \frac{\Delta E}{E} = \frac{4}{3} \frac{U_1}{c} (1 - \frac{U_2}{U_1}) x \] (48)

where

\[ x = 1 + \frac{1}{3} (1 + \frac{U_1}{2 U_2}) (1 + \frac{U_2}{U_1}) \] (49)

which is 9/4 for a strong shock when \( U_1/U_2 = 4 \).

It is easy to show that the additional energy gain flattens the particle spectrum. Consider then adiabatic losses and history of injection as above leads to a spectral correction relative to a planar shock of

\[ \frac{3 U_1}{U_1 - U_2} \left\{ \frac{U_2}{U_1} (\frac{1}{x} - 1) + \frac{2}{x} (b + d) \frac{\kappa_{rr,1}}{r U_1^2} \right\}. \] (50)

The total spectral change is then for \( U_1/U_2 = 4 \) given by 1/3, so that the spectrum obtained is

\[ \text{Spectrum (source)} = E^{-7/3}. \] (51)

After correcting for leakage from the galaxy the spectrum is

\[ \text{Spectrum (earth)} = E^{-8/3} \] (52)

very close to what is observed near earth at particle energies below the knee.

Such an injection spectrum of \(-7/3\) of relativistic particles in strong and fast shocks propagating through a stellar wind leads to an unambiguous radio synchrotron emission spectrum of \(\nu^{-2/3}\) from energetic electrons. Such a radio spectrum has been observed for the classical nova GK Per (Reynolds & Chevalier 1984).

Generalizing now for arbitrary wind speed and arbitrary shock strength we obtain for the thickness of the shocked layer

\[ \frac{\Delta r}{r} = \frac{U_2}{U_1} \frac{V_W + U_1}{V_W}. \] (53)

This reduction of the thickness of the layer for a finite wind velocity is due to the fact that the material which is snowplowed together is not all gas between zero radius and the current radius \( r \), but between zero radius and \( r(1 - V_W/(V_W + U_1)) \), since the gas keeps moving while the shock moves out towards \( r \). In this case we have

\[ f_d = \frac{1}{3} \left( 1 + \frac{1}{6} \frac{r U_1}{\kappa_{rr,1} U_2} \frac{U_1}{U_2} (1 - \frac{U_2}{U_1}) \right). \] (54)

It is easily verified that \( f_d = 1 \) for strong shocks and negligible wind velocity. We take here the curvature length scales to be the same on both sides of the shock, since the curvature is induced by the thickness of the shock region.
The energy gain associated with the $\theta$-drift is

$$\frac{\Delta E}{E^*} = \frac{4}{3} \frac{U_1}{c} (1 - \frac{U_2}{U_1}) x$$

where

$$x = 1 + 3 \frac{\kappa_{rr,1}}{r U_1} f_d (1 + \frac{U_2}{U_1})/(1 - \frac{U_2}{U_1}),$$

which is $9/4$ for negligible wind speeds and a strong shock when $U_1/U_2 = 4$; on the other hand, for $V_W/U_1 = 1$ we have $x = 2.042$ and for the limiting case of large wind speeds compared to shock speeds $x = 1.833$.

The spectral difference to the plane parallel case is thus

$$\frac{3U_1}{U_1 - U_2} \left\{ \frac{U_2}{U_1} \left( \frac{1}{x} - 1 \right) + \frac{2}{x} (b + d) \frac{\kappa_{rr,1}}{r U_1} (1 + \frac{V_W}{U_1}) \right\}.$$  \hspace{1cm} (57)

We note that

$$\kappa_{rr,1} (1 + \frac{V_W}{U_1})/r U_1$$

is now independent of the wind speed, and the only effect of the wind which remains is through $x$. For the sequence of $V_W/U_1 = 0, 1.0, \text{and} \gg 1$ we thus obtain particle spectral index differences, in addition to the index of $7/3$, of $0.0, 0.136, 0.303$, corresponding to synchrotron emission spectral index of an electron population with the same spectrum, of $0.667, 0.735, 0.818$. We will use the first two cases as examples below. The work of Owocki et al. (1988) suggests that typical shocks in winds have a velocity in the wind frame similar to the wind velocity in the observers frame itself, which implies that, in the simplified picture here, only spectral indices for the synchrotron emission between $0.667$ and $0.735$ are relevant, with an extreme range of spectral indices up to $0.818$ for strong shocks. Obviously, for weaker shocks with $U_1/U_2 < 4$ the spectrum can be steeper, e.g. for $U_1/U_2 = 3.5$ we obtain an optically thin spectral index for the synchrotron emission of $0.734$ for $V_W \ll U_1$ and $0.815$ for $V_W/U_1 = 1$.

2.2.2. The Polar Cap

We wish to discuss here the bend in the spectrum of cosmic rays at the knee, near $5 \times 10^6$ GeV.

Let us consider the structure of the wind through which the supernova shock is running. The maximum energy a particle can reach is proportional to $\sin^2 \theta$, since the space available for the gyromotion from a particular latitude is limited in the direction of the pole by the axis of symmetry. Hence, clearly the maximum energy attainable is lowest near the poles. Then, consider the pole region itself, where the radial dependence of the magnetic field is $1/r^2$, and the magnetic field is mostly radial. We can make two arguments here: Either we put the upstream diffusive scale $4 \kappa_{rr,1}/(c U_1)$ equal to $r/c$ in the strong shock limit, or we can put acceleration time...
and flow time equal to each other. Both arguments lead to the same result. Using here the Bohm limit in the diffusion coefficient \( \kappa_{rr,1} = \frac{4}{3} e E /Ze B(r) \), since we have a shock configuration near the pole, where the direction of propagation of the shock is parallel to the magnetic field - often referred to as a parallel shock configuration - then leads to a maximum energy for the particles of

\[
E = \frac{3}{4} Ze B(r) r \frac{U_1}{c}, \quad (59)
\]

which is proportional to \( 1/r \) near the pole, where the magnetic field is parallel to the direction of shock propagation; the corresponding gyroradius is then given by \( \frac{3}{4} U_1 r \).

Putting this equal to the gyroradius of particles that are accelerated further out at some cotlitude \( \theta \), where the magnetic field is nearly perpendicular to the direction of shock propagation, gives the limit where the latitude-dependent acceleration breaks down. This then gives the critical angle as

\[
sin \theta_{crit} = \frac{3}{4} \frac{U_1}{c}. \quad (60)
\]

The angular range of \( \theta < \theta_{crit} \) we refer to as the polar cap. The maximum particle energy at the location of the critical angle is then given by

\[
E_{knee} = Ze B(r) r \left( \frac{3}{4} \frac{U_1}{c} \right)^2. \quad (61)
\]

We identify this energy with the knee feature in the cosmic ray spectrum, since all latitudes outside the polar cap contribute the same spectrum up to this energy; from this energy to higher particle energies a smaller part of the hemisphere contributes and also, the energy gain is reduced, as argued below. This is valid in the region where the magnetic field is nearly perpendicular to the shock, and thus this knee energy is independent of radius. The argument described earlier on the limits of the \( \theta \)-diffusion leads to the same critical energy.

All this immediately implies that the chemical composition at the knee changes in such a way, that the gyroradius of the particles at the spectral break is the same, implying that the different nuclei break off in order of their charge \( Z \), considered as particles of a certain energy (and not as energy per nucleon).

Hence it appears plausible to suppose that in the polar cap the effective diffusion coefficient is quite a bit smaller than further down in latitude, and so that the particle spectrum is rather close to \( E^{-2} \). On the other hand, the polar cap is small relative to \( 4\pi \) by about \( (U_2/U_1)^2 \) and only a spectrum like \( E^{-2} \) will make it possible for the polar cap to contribute appreciably, but only near the knee energy, because the spectral flux near the knee is increased relative to 1 GeV by \( (E_{knee}/m_p c^2)^{1/3} \) which approximately compensates for its small area. The combination of the polar cap with the rest of the stellar hemisphere might lead to a situation where up to, say, \( 10^4 \) GeV the entire hemisphere excluding the polar cap dominates, while from \( 10^4 \) GeV up to the knee the polar cap begins to contribute appreciably. Near the knee energy the polar caps might thus contribute equally to the rest of the \( 4\pi \) steradians. Because of spatial limitations most of the hemisphere has to dominate again above the knee,
although with a fraction of the hemisphere that decreases with particle energy. The superposition of such spectra for different chemical elements is tested in paper CR IV and briefly summarized below. These tests with airshower data suggest indeed that near the knee of the overall cosmic ray spectrum the polar cap contributes noticeably, and produces the sharpness of the knee feature.

The expression for the particle energy at the knee also implies by the observed relative sharpness of the break of the spectrum that the actual values of the combination $B(r)rU_1^2$ must be very nearly the same for all supernovae that contribute appreciably in this energy range. Note that $B(r)r$ is evaluated in the Parker regime, and so is related to the surface magnetic field by $B(r)r = B_s r_s^2 \Omega_s/v_W$, where the values with index $s$ refer to the surface of the star and $v_W$ is the wind velocity. Thus the expression

$$B_s r_s^2 \frac{\Omega_s}{v_W} U_1^2$$

is approximately a universal constant for all stars that explode as supernova after a Wolf Rayet phase or red giant phase with a strong wind. It may hold for all massive stars of lower mass as well that explode as supernovae, and might even extend to below the mass the range of supernovae; this latter possibility is supported by the observation that the magnetic fields of white dwarfs appear to be higher for higher mass white dwarfs. This could be due to an overall evolutionary convergence of massive stars, which is indeed suggested by evolutionary calculations which, however, do not take strong magnetic fields or near critical rotation into account. This relationship suggests as an alternative the speculation that the mechanical energy of supernovae explosions is the potential energy from the transient phase when the interior of the exploding star collapses to an accretion disk, which loses its angular momentum due to torsional Alfvén waves. Assuming a core rotation of a star shortly before its explosion, which corresponds to critical rotation on the surface, this speculation leads with no further dominant parameter to an explosive energy of about $10^{51}$ erg (see paper CR I), corresponding very well to the typical mechanical energy of supernova explosions. Conceptually such a picture has some interesting similarities to protostar formation. Related ideas have been expressed and discussed by Kardashev (1970), Bisnovatyi-Kogan (1970), LeBlanc & Wilson (1970), Ostriker & Gunn (1971), Ammuel et al. (1972), Bisnovatyi-Kogan et al. (1976), and Kundt (1976), with Bisnovatyi-Kogan (1970) the closest to the argument here (see paper CR I). It remains as an interesting, but possibly quite difficult next step to calculate the evolution of massive stars allowing for the possibility of extreme rotation in the core and strong magnetic fields.

2.2.3. The Particle Spectrum above the Knee

Consider the derivation of the spectrum beyond the knee. Since the maximum energy particles can attain is a strong function of colatitude, the spectrum beyond the knee requires a discussion of the latitude distribution, which we have to derive first. The latitude distribution is established by the drift of particles which builds
up a gradient which in turn leads to diffusion down the gradient. Hence it is clear that drifts towards the equator lead to higher particle densities near the equator, and drifts towards the poles lead to higher particle densities there. Thus the equilibrium latitude distribution is given by the balancing of the $\theta$-diffusion and the $\theta$-drift.

The diffusion tensor component $\kappa_{\theta\theta}$ can be derived similar to our heuristic derivation of the radial diffusion term $\kappa_{rr}$, again by using the smallest dominant scales. The characteristic velocity of particles in $\theta$ is given by the erratic part of the drifting, corresponding to spatial elements of different magnetic field direction, and this is on average the value of the drift velocity $|V_{d,\theta}|$, possibly modified by the locally increased values of the magnetic field strength, and the characteristic distance is the distance to the symmetry axis $r \sin \theta$; this is the smallest dominant scale as soon as the thickness of the shocked layer is larger than the distance to the symmetry axis, i.e. $\sin \theta < U_2/U_1$. Thus we can write in this approximation, postulate 3,

$$\kappa_{\theta\theta,1} = \frac{1}{3} |V_{d,\theta}| r (1 - \mu^2)^{1/2}. \quad (63)$$

Here $\mu$ is again the cosine of the colatitude on the sphere we consider for the shock in the wind. Interestingly, this can also be written in the form

$$\frac{1}{3} r_g c (1 - \mu^2)^{1/2},$$

where $r_g$ is taken as positive; we also note that $c (1 - \mu^2)^{1/2}$ is the maximum drift speed at a given latitude, valid for the local maximum particle energy. This suggests that the latitude diffusion might be usefully thought of as diffusion with a length scale of the gyroradius, and the particle speed, to within the angular factor which just cancels out the latitude dependence of the magnetic field strength in the denominator of the gyroradius.

We assume then for the colatitude dependence a powerlaw $(1 - \mu^2)^{-a}$ and first match the latitude dependence of the diffusion term and the drift, and then use the numerical coefficients to determine the exponent in this law. The diffusion term and the drift term have the same colatitude dependence since the double derivative and the internal factor of $(1 - \mu^2)$ lead to a $(1 - \mu^2)^{-a - 1}$ for the diffusive term, while the drift term is just the simple derivative giving the same expression. For $(1 - \mu^2) \ll 1$ the condition then is $\frac{2}{3} a^2 = \pm a$.

It is important to remember the sign of these terms. The diffusive term is always positive, while the $\theta$-drift term is negative for $Z B_s$ negative. This means for positive particles and a magnetic field directed inwards the $\theta$-drift is towards the pole. In that case then the exponent $a$ is either zero or $a = 3/2$. Since the drift itself clearly produces a gradient, the case with $a = 0$ is of no interest here. It follows that for positive particles and an inwardly directed magnetic field the latitude distribution is strongly biased towards the poles, emphasizing in its integral the lower energies, and thus making the overall spectrum steeper beyond the knee energy. When the magnetic field is directed outwards and the particles are positive, the drift is towards the equator with then a positive gradient with $(1 - \mu^2)^{3/2}$, again in the limit $(1 - \mu^2) \ll 1$. We
note that this exponent 3/2 is reduced in the case, when the erratic part of the drift is increased over the steady net drift component.

In our model for the diffusion in $\theta$ we have used the drift velocity and the distance to the symmetry axis as natural scales in velocity and in length. When the $\theta$-drift reaches the maximum derived earlier, then the latitude drift changes character. This happens at a critical energy, which is reached at

$$\kappa_{\theta,1} = \kappa_{\theta,\text{max}}.$$  \hspace{1cm} (64)

which translates to

$$E_{\text{crit}} = \left(\frac{3}{4} \frac{U_1}{c}\right)^2 E_{\text{max}} = E_{\text{knee}}.$$  \hspace{1cm} (65)

We emphasize that two different basically geometric arguments lead to the same critical energy, $E_{\text{knee}}$. This means, postulate 4, that for

$$E > E_{\text{knee}}$$  \hspace{1cm} (66)

the drift energy gain is down by 2/3 to the level what the pure gradient and curvature drift yields, which results in

$$x = 11/6.$$  \hspace{1cm} (67)

The reduced drift energy gain reduces the value of $x$ below the limiting value derived earlier, of 9/4. This then leads to an overall spectrum of

$$E^{-29/11}$$  \hspace{1cm} (68)

before taking leakage into account, and

$$E^{-98/33} \approx E^{-3},$$  \hspace{1cm} (69)

with leakage accounted for. This is what we wanted to derive. The observational test with airshower data confirms (see paper CR IV and below), that the particle spectrum beyond the knee is close to the spectrum derived here; in fact it suggests a slightly steeper spectrum, of close to $E^{-3.07}$, about 0.1 steeper than derived here. This is clearly within the uncertainties of our analytical approach.

2.2.4. Observational Test: Radio emission

There are a number of consistency checks that are possible with the radio emission of Wolf Rayet stars, radiosupernovae and OB stars (Abbott et al. 1986, Bieging et al. 1989, Cassinelli 1983, 1991, Maheswaran & Cassinelli 1988, 1992; the references for radiosupernovae are given below; for details see paper CR II). Additional checks with X-ray and $\gamma$-ray data are possible as well as with novae.

First of all, it is possible to demonstrate that the classical dynamo process (Ruzmaikin et al. 1988) can produce high magnetic fields inside of massive stars, of order $10^6$ to $10^7$ Gauss, so that it appears quite plausible that already OB stars
which are radiative outside have appreciable strengths of their surface magnetic fields, of order $10^3$ to $10^4$ Gauss. Rotationally induced circulations can carry the fields on fairly short time scales to the surface for rapidly rotating stars. Wolf Rayet stars have as their surface former convective layers (Langer 1989) and so easily can have as high magnetic fields as suggested by the cosmic ray argument, also of order $10^3$ to $10^4$ Gauss. Secondly, it is possible to show that the magnetic fields can provide additional momentum to the wind of Wolf Rayet stars, and so may help to explain the strength of these winds (see paper CR II), acting as an amplifier to the radiative forces.

The radio luminosities of Wolf Rayet stars and of radiosupernovae can readily be understood with synchrotron emission and free-free absorption (for free-free absorption Altenhoff et al. 1960, for a proper integration of free-free emission Biermann et al. 1990, for radiosupernovae Weiler et al. 1986, 1989, 1990, 1991, Panagia et al. 1986, for the supernova 1987A, Turtle et al. 1987, Jauncey et al. 1988, Staveley-Smith et al. 1992, for the new supernova in M81 e.g. Strom et al. 1993). The comparison suggests that there is a critical shock velocity below which electron injection gets very weak. For the winds of massive stars this critical velocity is of the order of 5000 km/sec, interestingly, but probably misleadingly close to the shock velocity for which the downstream thermal electrons reach the speed of light.

2.2.5. Observational Test: Air Shower data

We wish to test the overall model proposed by asking whether it can successfully account for the spectrum and chemical composition at particle energies beyond $10^4$ GeV: the special difficulty in this endeavour is the fact that we know the chemical abundances near TeV particle energies already, and extrapolating these spectra does not readily give the bump and knee of the well established overall particle spectrum.

The difficulty in all such attempts originates in the relation between energy estimation with air showers and the mass of the primary nucleus. Generally showers initiated by heavy nuclei develop and are absorbed faster in the atmosphere. Heavy nuclei thus produce showers of smaller size at the observation levels of all existing experiments. The effect is stronger for inclined showers, which have to penetrate through larger atmospheric thickness. We model the shower size that the primary cosmic ray flux generates in the Akeno detector (Nagano et al. 1984), using both vertical showers and slanted showers; having performed such a fit, we then reconstruct the all particle spectrum.

Various cosmic ray experiments have given data about the chemical composition near TeV energies and somewhat beyond, up to 100 TeV, most notably from the Chicago Group (Grunsfeld et al. 1988, Müller 1989, Swordy et al. 1990, and Müller et al. 1991), the JACEE experiment (Parnell et al. 1989, Burnett et al. 1990, and Asakimori et al. 1991a, b), Simon et al. (1980) and Engelmann et al. (1990). The Oxygen data, where we have averaged existing data from all four experiments when independent measurements exist in the same energy range, clearly demonstrate a power law behaviour with a somewhat flatter spectrum than the low energy index of about 2.75. Oxygen has a spectrum of about 2.64, corresponding well to the argu-
ment made in CR I, that supernova explosions into existing stellar winds produce not only higher particle energies, but also spectra flatter than the lower energy particles from Sedov phase explosions into the interstellar medium. We use all these known chemical abundances as a given input, combining the elements into six groups, and the sources into three different sites as discussed above. The six element groups which we use are: a) Hydrogen, b) Helium, c) Carbon, Nitrogen and Oxygen, d) Neon to Sulphur, e) Chlorine to Manganese, and f) Iron. We argue that the galactic cosmic rays cut off at $10^8 Z\text{GeV}$ and that the extragalactic component is nearly all protons and Helium nuclei. In paper CR II (Biermann and Cassinelli 1993) we argue that this is in fact suggested by a number of stellar observations.

The main uncertain parameters are then the following:

1) The spectral difference between the Sedov phase explosion particles and the wind explosion particles.

2) The particle energy $\sim Z$ at the cutoff for the Sedov phase explosion acceleration.

3) The relative strength of the polar cap component.

4) The spectral difference between the energy range below and above the knee.

5) The particle energy $\sim Z$ of the knee.

6) The upper cutoff energy $\sim Z$ for the component beyond the knee.

The theory predicts specific numbers for all these parameters. The division of the abundances between the Sedov phase accelerated particles and the wind explosion accelerated particles can be made using existing data, especially from JACEE, and thus is no free parameter.

We use directly the shower size spectra to reconstruct the cosmic ray spectrum under the assumption of different chemical compositions. We do this simultaneously for two zenith angles, vertical ($\sec\theta = 1.0$) and for $\sec\theta = 1.2$, i.e. one moderate slant angle. At these angles the statistical errors are small. The shower size is expressed by $N_e$ - the total number of charged particles at observation level. The flux of showers with $N_e = 10^6$ at $\sec\theta = 1.2$ is smaller than the vertical one by a factor of 7. The shapes of the size spectra are also different.

The detailed modelling procedure is described in Stanev et al. (1993, paper CR IV). The comparison with the data demonstrates that the fits are well within the acceptable error ranges. The model has the following parameters:

The cutoff energy (exponential cutoff) for the Sedov phase accelerated particles is 120 TeV for protons; the spectral index of the wind component is 2.66 below the knee and 3.07 above the knee; this steepening occurs at 700 TeV for protons and is rigidity dependent. The ratio of the polar cap component to the steeper wind component is 1 for heavy nuclei at the bend. Although the polar cap component does not contribute to shower sizes significantly below or above the knee, it is essential for reproducing the sharp break at $N_e = 10^6$. The reproduction of this shower break, as well as the observed continuous flattening of the spectra of all heavy nuclei requires the introduction of a very flat ($E^{-2}$ at injection) acceleration component, i.e. the polar cap component. All four major parameters are close to their expected values (predicted in the earlier papers CR I, CR II, and CR III).
It is clear that within the framework of our picture the knee feature in the overall spectrum can only be due to the addition of the polar cap component. Without it the generated size spectrum cannot have a sharp bend in any model similar to ours, with rigidity dependent features and composed of different nuclear components.

We note that the shower size distribution is very sensitive to changes in the parameters. Relatively small changes (by more than 20%) of the cutoff and bend energy, as well as in the composition of the wind component, affect the calculated size spectra so strongly that they become inconsistent with data. Together with the direct measurements in the 1 TeV region the shower size distributions constrain the models to within a small parameter space. Fitting the shower size distribution at different zenith angles is critical despite its fairly large error bars.

Several properties stand out:

1) The correct overall spectrum is lower in the knee region and beyond than the conventionally derived spectrum; the factor is about 2. The knee itself is not as sharp but the change of the spectral index takes place at approximately the same energy.

2) The composition changes rapidly in the knee region, becomes increasingly heavier, and in our current interpretation is dominated by the Neon group above the knee.

3) The proton flux at energy above $10^8$ GeV represents the extragalactic component. Because the wind component has a maximum energy of $10^8$ GeV for protons ($2.6\times10^9$ GeV for Fe) the calculated shower size spectra would become too steep for sizes above $10^7$ without this extra component. This implies an extragalactic flux (presumably mostly protons) that is consistent with the independently derived estimate by Gaisser et al. (1993) using the Fly’s Eye data; a full comparison with an extragalactic model (Rachen & Biermann, 1992, 1993: paper UHE CR I) has been done in Rachen et al. (1993, paper UHE CR II).

4) The comparison demonstrates that our curves fit all existing data quite well, including the trend, observed in heavy nuclei, to exhibit a flattening of the spectrum at the approach to the knee. The seeming inconsistency with the highest energy direct data for iron is possibly due to the fact the the JACEE experiment (where these points come from) has presented its measurements for nuclei with $Z > 17$. It is difficult to judge what the absolute normalization of the cosmic ray flux beyond the knee is. The presentation of the spectrum in the form used here ($E^{-2.75}dN/dE$) tends to exaggerate differences in the absolute normalization. A relatively small error of the energy determination in air showers (typical error of 20%) changes the normalization of the absolute flux by a large amount (65%).

We conclude that the most stringent test yet of the model proposed in the earlier papers of the series on cosmic ray physics (CR I, CR II, CR III, UHE CR I, UHE CR II), which give predicted spectral shapes for all three sites of origin for cosmic rays is successful. Several detailed conclusions can be made:

1) The abundances of the cosmic rays from explosions into stellar winds do not require nor imply any admixture from the heavy elements newly produced in the star exploding. All the heavy elements already present in the stellar wind (enriched already
to some degree) prior to the explosion participate in the feeding of the accelerated particle population, and no other additional source is required.

2) The nuclei accelerated in the polar cap region are essential for a successful fit of the shower size spectra. Although the total energy carried by such nuclei is not more than $1/100$ of the total wind component, they contribute a major fraction in the knee region.

3) There is no requirement for other cosmic ray sources, either from spectral arguments or from abundance arguments; thus pulsars and compact X-ray binary systems may accelerate lots of particles, but they need not play a dominant role out in the typical interstellar medium.

4) The large fraction of heavy nuclei inferred from the Fly’s Eye data (Gaisser et al. 1993) requires the acceleration of heavy element nuclei out to at least $3 \times 10^9$ GeV. This implies (see CR I and II) that the winds of the stars that explode as supernovae have magnetic fields at least as strong as 3 Gauss at a fiducial distance from the star of $10^{14}$ cm. Since this limiting particle energy is already derived by using the spatial limit given by the condition that the Larmor radius of the particle fit into the space available, there is no other way than indeed having these high fields in the stellar winds of massive stars in the context of our theory (note that also Usov & Melrose 1992 argue for high magnetic fields, on different grounds). The magnetic field strengths required are of order $10^3$ to $10^4$ Gauss on the surface of the stars, with the exact value clearly depending on stellar radius and exact wind configuration. An important next step is to work out the consequences of these models at lower energies.

2.3. Starburst and normal Galaxies

Starburst galaxies clearly produce a large number of supernovae, accelerate cosmic rays to energy densities in their interstellar medium far above that of the interstellar medium in the solar neighborhood, and also produce break-outs from their gaseous disks to produce winds, which are visible in a) optical polarized light and line emission, in b) X-rays as well as in c) nonthermal radio emission (Kronberg et al. 1981, 1985, Schaaf et al. 1989). Normal galaxies are also thought to have galactic winds, albeit of less conspicuous power (Breitschwerdt et al. 1991, 1993).

We would like to obtain an estimate for the cosmic ray production. For this we need to discuss three points, a) the power output of galaxies in cosmic rays, b) their spectrum, and c) their low energy cutoff.

As our reference we take the well studied starburst galaxy M82 (e.g. Rieke et al. 1980, Kronberg et al. 1981, 1985, Schaaf et al. 1989). M82 is the first galaxy to yield a large number of observable radio supernova remnant candidates (Kronberg et al. 1985), the stronger ones being time variable; this allows a first estimate of the actual supernova rate. In Kronberg et al. (1985) we estimate a supernova rate of about 1 per 3 years, with an uncertainty of a factor of three. Assuming that the efficiency of producing cosmic rays is about 10% and the energy input per supernova is $10^{51}$ ergs, we thus arrive at a cosmic ray production luminosity of M82 of $10^{42}$ ergs/sec; we have to compare this with the farinfrared luminosity of (Rieke et al. 1980) $4 \times 10^{10} L_\odot$. This gives a ratio of cosmic ray luminosity to farinfrared
luminosity of 0.0069 with an uncertainty of a factor of three; we adopt as our reference ratio 0.01 and keep the uncertainty in mind. It is an important next step to work through the large amount of radio data assembled by Kronberg and his colleagues on M82 to check on all these detailed models.

The spectrum of cosmic rays outside the galaxy converts back to its source spectrum in the simple leaky box argument; although there are obvious problems with such arguments (see Biermann 1993), there is nothing better at present and so we will use this approach here as well. The acceleration of cosmic rays is discussed above. For the low energy cosmic rays this means that their spectrum is about $p^{-2.4}$ at injection, and again outside the galaxy.

This injection spectrum is for the transrelativistic regime. Hence the next question is the low momentum cutoff in this spectrum. Here we have information from two sources. First, when one repeats the argument of Spitzer & Tomasko (1968) for interstellar clouds (Jokipii and Biermann, 1991, as yet unpublished) then one finds that in order to explain the cosmic ray ionization rate in interstellar clouds required by the chemistry (Black & van Dishoeck 1987, Black et al. 1990, Black & van Dishoeck 1991) low momentum cutoffs corresponding to kinetic particle energies of $30 - 100$ MeV. Second, the spallation of cosmic rays produces Be and other light elements (Gilmore et al. 1992) also gives an indication of the low kinetic energy cutoff of the cosmic ray spectrum; again, the cutoff is in the same kinetic energy range. The high end of the particle spectrum is due to explosions into winds, and thus has a rather steep spectrum, but is dominated by heavy nuclei.

We thus find the following estimates: a cosmic ray luminosity of about 1% of the farinfrared luminosity, with a spectrum in momentum of $E^{-2.4}$, and a low kinetic energy cutoff of 30 MeV. Except for the spectrum, these estimates are fairly uncertain.

These cosmic ray particles are subject to adiabatic loss as they come out from a galaxy, and so their energy density as well as their low and high end particle energies are lowered by a factor which is likely to be of order at least 10. For various plausible models of galactic evolution this extragalactic cosmic ray population can ionize and heat the intergalactic medium (Nath & Biermann 1993), and help to let us understand the extreme degree of ionization of the intergalactic matter implied by the Gunn-Peterson test as well as the low Compton $y$ parameter implied by the COBE data (Mather et al. 1993). The temperature of the intergalactic medium in such models is about an order of magnitude higher than based on photoionization models, but still quite low compared to temperatures which would be required to give a substantial contribution to the X-ray background. A further desirable next step would be a calculation of models for Lyman $\alpha$-clouds taking low energy cosmic ray particles into account.

3. Radio Galaxy Hot Spots: Highest Energy Cosmic Rays

Any model of the origin of the highest energy cosmic rays has to fulfill the following conditions: (i) the predicted sources must be able to accelerate particles to
the observed energies up to at least 100 EeV, (ii) the energy content in relativistic particles and the number density of the sources must be sufficient to provide the observed UHE-CR flux, (iii) the observed UHE-CR spectrum must be fitted in the energy region where the considered contribution dominates over all others, and (iv) the chemical abundances must match the new Fly’s Eye results (Gaisser et al. 1993).

The investigation of FR II radio galaxies constitutes a fairly good basis for this purpose: FR II galaxies are very bright radio objects, their distribution and evolution is well known up to rather high redshifts (Peacock 1985, Barthel 1989), and their content of relativistic particles can be well estimated from their synchrotron emission (Rawlings & Saunders 1991). The only problem is to obtain a reliable estimate of the proton to electron ratio in the energy density of relativistic particles; this uncertainty introduces a “fudge factor” in our model, which can only be tested on plausibility but not calculated at present.

3.1. Prediction

Clearly, the hot spots in FR II galaxies are not expected to be the only extragalactic objects that can accelerate particles to ultra-high energies. Most active galactic nuclei might provide much better conditions for that, and shock acceleration can also take place in the jets of less luminous FR I galaxies (e.g. M87). The reason for the restriction to the rare class of FR II galaxies is, that highly energetic charged particles produced deep inside galactic structures will suffer substantial adiabatic losses on their way out to the extragalactic medium. This problem applies to all possible cosmic ray sources except the FR II hot spots, since they are located at the edge of the extended radio lobe of the galaxy and the particles can readily enter the extragalactic space. However, protons can escape from galactic cores with very high energies, since they can undergo isospin flips in $pp$ or $\gamma p$ reactions, leaving the galaxy as an UHE neutron. At present it is not clear whether the resulting CR spectrum even extends to the highest energies considered here (Protheroe & Szabo 1992, Mannheim 1993a, b); however, it is clear that airshower data require heavy nuclei to dominate beyond the knee in contradiction to the model by Protheroe & Szabo (1992).

Our model is based on the present knowledge about diffusive shock acceleration, radio galaxy properties, galaxy evolution and a simple model for intergalactic cosmic ray propagation.

The hot spots in strong extended radio galaxies, denoted as FR II galaxies following the classification of Fanaroff & Riley (1974), are identified as the endpoints of powerful jets ejected by the active nuclei of the galaxies deep into the extragalactic medium (see, e.g., Meisenheimer & Röser 1989). It has been shown that the radio-to-optical spectra emitted by those hot spots can readily be explained as synchrotron radiation from particles accelerated at a strong shock wave by the first order Fermi mechanism (e.g. Biermann & Strittmatter 1987, Meisenheimer et al. 1989).

The basic argument (Biermann & Strittmatter 1987) runs as follows: Protons and other heavier nuclei are accelerated as explained above (see, e.g. Drury 1983) for parallel shocks; these energetic nuclei have fairly low losses and thus attain large
particle energies, and are thus the first "messenger" to the upstream flow of the coming shock transition. The nuclei cause plasma turbulence to develop which initiates a cascade in wavenumber space. This cascade is assumed to be, on the same basis as argued above (Matthaeus & Zhou 1989, and by analogy with the solar wind where the turbulence can be measured by spacecraft), of a Kolmogorov character. Thus the mean free path for the scattering of the electrons is fixed by the energetic nuclei. This mean free path for the scattering of electrons in turn leads to a maximum particle energy for electrons, which translates to a maximum synchrotron emission frequency, which is independent of the strength of the magnetic field and thus ought to be similar in both compact sources as well as in extended hot spot regions; this is exactly the behaviour found in many red (nonthermal) quasar emission spectra, in BL Lac objects, in the M87 jet, and in radio hot spots. Thus, the ubiquitous cutoff (many papers starting with Rieke et al. 1976) in the electron synchrotron emission finds a ready explanation. There is some variance in the exact cutoff frequency due to the possibilities (all neglected except the last point by Biermann & Strittmatter 1987) due to a) a slight obliqueness of the shock, b) different transport coefficients on both sides of the shock, c) weak relativistic boosting between shock frame and observers frame, d) time dependence and e) the possible influence of a radiation field, which can modify the observed cutoff frequency. There is a possibility to check the assumed Kolmogorov law with the data of Meisenheimer et al. (1989): since particles accelerated at the shock continuously lose energy as they move downstream with the flow, the integration of the emission over the downstream region produces a bend in the emission spectrum at some frequency which is inversely related to the length of the effective downstream region. With this effective length available from observation together with the bend and the cutoff frequency one can limit the possible range of turbulence spectra of the plasma cascade, and the result suggests support for the initial assumption of a Kolmogorov spectrum.

Given this agreement, we can then use the theory to derive maximum particle energies for protons and other nuclei. For the protons a high energy cutoff of the order of $10^{21}$ eV is found if their energy is mainly limited by synchrotron losses and p-γ-collisions rather than by the finite range of the diffusion region or by interactions with ambient photons.

The simple canonical theory of diffusive shock acceleration explains the main hot spot features and is therefore a very attractive model for the acceleration of cosmic ray particles. Clearly, slightly flatter spectra as predicted by enhanced acceleration models (Krülls 1992) are consistent with most of the results within the error estimates, even required by the observations of Pic A west.

FR II galaxies are extended, steep spectrum radio galaxies, that are uniquely classified by their double structure due to their luminous hot spots. This morphology occurs for radio luminosities above $2 \cdot 10^{26}$ WHz$^{-1}$ at 178 MHz, which makes it easy to distinguish FR II galaxies from other radio galaxies in the radio luminosity function. The present state of knowledge is that roughly 70% of all steep spectrum radio sources in this power range show FR II structure (Perley 1989), but on the other hand there may also be some compact flat spectrum sources that can be identified with FR II
galaxies having their jets oriented along the line of sight (Padovani & Urry 1992). At high redshifts (above z=1), however, little is known about the morphology of the sources and it is not clear whether the Fanaroff-Riley relation still holds; but we will see that this uncertainty affects the cosmic ray spectrum only below 1 EeV.

The synchrotron radiation emitted by relativistic electrons is the only (or at least the main) information we get about the content of relativistic particles in jets. However, the total synchrotron luminosity $L_s$ of a source allows us to give minimum estimates about their mean energy density in relativistic particles $U_p$ and $U_e$, protons and electrons respectively, and in the magnetic field $U_B$. The minimum energy condition is given by, e.g., Pacholczyk (1970). We introduce the quantity $k_p = U_p/U_e$. The emission of photons in the X and γ regime from UHE protons in hot spots by synchrotron radiation and the proton-induced electromagnetic cascades can be expected (Mannheim & Biermann 1989, Mannheim et al. 1991). Recent GRO detections of blazars, which may be considered as compact Doppler-boosted emission from jets support this expectation (Mannheim & Biermann 1992, Mannheim 1993a, b) and give evidence for the presence of energetic protons as employed here. In the future, if good observations of X-ray and hard γ emission from hot spots can be obtained, it may even be possible to derive definite estimates on $k_p$. However, any such estimate has to overcome the difficulty that we do not know how to estimate the efficiency of converting jet-power into ultimate energy density in relativistic nuclei in the hot spots.

For considering FR II galaxies as cosmic ray sources, we are only interested in the “proton luminosity” of a source. From the synchrotron luminosity one can derive the total jet power, assuming some value for $k_p$. Rawlings & Saunders (1991) gave the jet power for 39 FR II galaxies, assuming no protons in the jet ($k_p = 0$).

$f$ is the fudge factor in our normalization containing all the uncertainties briefly mentioned above. We derive $f$ from fitting our calculated spectra to the data and therefore obtain information on the relative proton content $k_p$ required.

To calculate the extragalactic source function of cosmic ray protons, we have to apply our knowledge about jet powers and proton luminosities to the epoch-dependent radio luminosity function (RLF) of FR II galaxies. The RLF $n_G(P_\nu, z)$ gives the number density of radio galaxies with a specific radio luminosity $P_\nu$ at a given frequency $\nu$ at an epoch $z$ respective to the comoving unit volume element. RLF’s can be derived from radio source counts, using the available information about redshifts and assuming some modelling parameters (Peacock & Gull 1981, Windhorst 1984, Peacock 1985).

To connect the proton source function to the RLF, the given relation between total synchrotron luminosity and kinetic jet energy is not very useful; we rather have to find a relation between jet energy and radio luminosity at the frequency for which the RLF is given. In this work, we use the various luminosity functions given by Peacock (1985) for different modelling parameters, based on source counts at 2.7 GHz for steep spectrum radio sources.

Protons with Lorentz factors $\gamma_p \geq 10^{10}$ can undergo high energy reactions even with the very low energy photons of the universal microwave background radiation.
(MBR), as pointed out first by Greisen (1966) and independently by Zatsepin & Kuzmin (1966). The attenuation length reduces to merely 10 Mpc for $\gamma_p > 10^{11}$ (Stecker 1968), what causes a cutoff in the particle spectra of distant extragalactic sources, usually denoted as the Greisen-cutoff. Hillas (1968) demonstrated the influence of the cosmological evolution on this effect; for very distant sources the Greisen-cutoff appears at much lower energy, because of the higher density and temperature of the MBR at earlier epochs. The most extensive investigations of the transport of UHE protons in the evolving MBR were done by Hill & Schramm (1985) and by Berezinsky & Grigor’eva (1988), who both considered the modification of the spectrum due to the cosmological evolution of the proton sources. We will use here the method of Berezinsky and Grigor’eva.

In all the discussion so far we only discussed the evolution of proton spectra in time; the spatial propagation was always assumed to be rectilinear. We apply our method to the calculation of proton spectra from single sources with a clearly defined spatial distance; if there is any appreciable magnetic field in extragalactic space, the particles propagate by a combination of drift, convection and diffusion, so that our results may no longer be valid. We use the approximation of straight line paths for the energetic particles. The limitations and merits of such an approach are discussed at some length in paper UHE CR 1.

The results derived from this modelling lead to spectra with an $E^{-2}$ behaviour at energies well below EeV, then a section with roughly $E^{-2.75}$ and a bump, and finally a sharp cutoff near 100 EeV; the chemical composition is mostly Hydrogen.

### 3.2 Tests with Airshower data

Here we compare the prediction introduced above with observational data from both the Fly’s Eye and the Akeno airshower detectors. These experimental data are now available from the analysis of the chemical composition near EeV energies (Gaisser et al. [1993] for the Fly’s Eye experiment in an accurate analysis, and from Stanev et al. [1993] for the Akeno experiment as a consistency check) and demonstrate that below 1 EeV protons show a flatter spectrum than the overall spectrum and thus their relative proportion increases with energy. The comparison of the spectral data and the flux shows that the ultra high energy component of the cosmic rays can indeed be understood as arising from the hot spots of powerful radiogalaxies, in flux, spectrum and chemical composition: This result puts stringent limits on the propagation of high energy cosmic rays and thus on the properties of the intergalactic magnetic field.

We can use the world data set on the cosmic rays to obtain an averaged spectrum of all cosmic rays beyond 0.1 EeV as a basis for the absolute normalization of the new component.

The measurement of the flux of ultra high energy cosmic rays is difficult because of their small flux (less than 0.05 particles per m$^2$ ster.year above $E = 10^8$ GeV). Still relatively large statistics have been collected during the last 20 years by four experiments, three of which are still active (Akeno (Nagano et al. 1984; 1992), Yakutsk (Efimov et al. 1991) and the Fly’s Eye (Cassiday et al. 1990; Loh et al. 1991)).
The array at Haverah Park (Cunningham et al. 1980) does no longer exist. A recent presentation and comparison of the results of the four arrays is made by Sokolsky et al. (1992).

Apart from the differences in normalization all data sets show the same slope ($\gamma = 3.0 - 3.1$) in the region $3 \times 10^{17} - 3 \times 10^{18}$ eV. We use the data in this range to normalize experiments to each other. First an average spectrum is calculated using the original experimental errors. Then we determine for each experiment the energy shift $\Delta E$ that brings the data set in best agreement with the average spectrum. The resulting shifts are smaller than the suspected systematic errors – we obtain $\Delta E/E = -16\%$ for Yakutsk, $-13\%$ for Akeno, $-3\%$ for Haverah Park (which has the smallest statistical errors) and $+6\%$ for the Fly’s Eye. We apply these energy shifts to the full data sets and obtain a world data set where the scatter is smaller than the individual error bars. Finally we bin in logarithmically equal energy bins, combine individual data points within the same bin and produce the average flux. This is the flux that we shall use to estimate the strength of different chemical components further down.

This procedure is not unique and does not eliminate the systematic errors in the energy derivation. We claim, however, that it introduces a standard systematic error, which is nearly the same for all data sets. The energy shifts required for the averaging are remarkably small. A better way to achieve a standard energy derivation from different experiments is a detailed study of the experimental algorithms (M. Lawrence et al. 1991) and a comparison and possible improvement of their theoretical bases.

We have produced a series of model calculations for an input spectrum of $E^{-2}$ with an intrinsic exponential cutoff at different energies in the range about 100 EeV, and a Hubble constant of 75 km sec$^{-1}$ Mpc$^{-1}$. Such calculations involve an estimate for the conversion of radio luminosity to jet power (Rawlings and Saunders 1991), from jet power to energy in energetic particle populations, and, most importantly, the ratio between the energy density of the energetic electrons, and the energetic nuclei – after all, the observed synchrotron emission traces only the relativistic electrons. The combination of these factors is the “fudge factor” introduced above, and it is probably dominated by the unknown nuclei to electron ratio. Low energy fits to the Fly’s Eye data for protons (see below) imply this factor to be near 3, and so suggest that the proton to electron ratio is of order 20 or less (see UHE CR I and above), a rather uncertain limit.

We have used two independent methods to estimate the extragalactic flux from different air shower experiments. First we fit the shower size distributions published by the Akeno experiment (Nagano et al. 1984) for different zenith angles. Then we use the results of a recent analysis (Gaisser et al. 1993) of the cosmic ray composition around $10^{18}$ eV from the measurements of the Fly’s Eye detector (Cassiday et al. 1990).

The Akeno shower array is a traditional detector that measures the number of charged particles in the shower ($N_c$) that reach the observation level. We use data sets for vertical ($\sec \theta = 1.0$, atm. depth = 920 g/cm$^2$) and slightly inclined
(\text{secan} \theta = 1.2, \text{atm. depth} = 1104g/cm^2) \text{showers. The fitting procedure is described in detail in paper CR IV and consists of stepping through the energy spectrum of each chemical component with a small logarithmic step (10^{0.01}) and calculating the size of a number of showers at that energy and primary mass with the parametrization of Gaisser (1979). The resulting sizes at both depths are binned with the appropriate weight, and the sum for all components is then compared with the experimental data.}

In the spectrum and composition model of paper CR IV all galactic nuclear components have the same spectral index of 3.07 at energies above $4 \times 10^{16}$ eV. The comparison with the experimental $N_e$ distribution demonstrates that galactic cosmic rays fit well the shower size distribution up to $N_e = (1. - 3.) \times 10^7$, but are insufficient to maintain the calculated spectrum, especially for inclined showers, in agreement with experimental data for bigger $N_e$. A better fit requires the introduction of a flatter cosmic ray component at energy above $10^{16}$ eV. The agreement is achieved by introducing a component, consisting of pure Hydrogen, with an energy spectrum of $(1.4 \pm 0.3)10^{-7} \times E^{-2} \text{ (cm}^2 \text{ster} \text{s GeV)}^{-1}$ at energies between $10^{16}$ and $5 \times 10^{17}$ eV and $(0.47 \pm 0.09) \times E^{-2.75} \text{ (cm}^2 \text{ster} \text{s GeV)}^{-1}$ at higher energy, where $E$ is measured in GeV. Because the spectral shape of this component is entirely different from those of the galactic cosmic rays it is very likely to represent an emerging extragalactic cosmic ray flux.

For the second estimate we use the recent results from the analysis of the Fly’s Eye measurements of the depth of maximum ($X_{\text{max}}$) distribution in terms of cosmic ray composition. The Fly’s Eye is a different type of detector, which observes directly the longitudinal development of air showers through the detection of the fluorescent light from the atmospheric Nitrogen atoms, induced by the shower charged particles. The amount of fluorescent light is proportional to the number of charged particles after an account is made for light scattering and absorption. The data analysis fits individual data points (taken at various atmospheric depths) to a shower profile and derives the depth ($X_{\text{max}}$) and size ($N_{\text{max}}$) of the shower maximum. $N_{\text{max}}$ is proportional to the energy of the primary nucleus and $X_{\text{max}}$ depends on the energy and mass of the primary nucleus. The sensitivity to the primary mass comes from the rate of energy dissipation, which is faster in showers initiated by heavy nuclei. $X_{\text{max}}$ of Fe generated showers is about 100 g/cm$^2$ shallower than that of proton generated shower of the same energy.

The basic idea of the shower analysis is to simulate a large number of air showers and compare the results to data. The error bars include both the statistical errors from the experimental statistics and the fitting procedure and the systematic errors from the particle physics input in the simulation, which also make the errors very asymmetric.

The best curve has an intrinsic cutoff of 100 EeV, and a spectrum which is nearly $E^{-2.75}$ between 1 EeV and 30 EeV. We selected this curve because it fits best the experimental data simultaneously in the low and the high energy range. The fit is quite good, considering the error bars on the data, fitting to about $2\sigma$ or better.

Most of the existing theoretical proposals to explain the origin of the cosmic ray particles at energies beyond the knee can be excluded on the basis of the presented
data:

1. Beyond the knee the particles are dominantly heavy nuclei and not all protons as follows from the proposals by Protheroe & Szabo (1992), and by Salamon & Stecker (1992).

2. The galactic wind model by Jokipii & Morfill (1987) would give chemical abundances much closer to solar abundances for these cosmic ray particles than the models presented here, as would the reacceleration model (Ip & Axford 1992) except for the highest particle energies. Again, the observed abundances clearly disagree with these models. The essential reason for the difference in the predicted chemical abundances is, that reacceleration models derive from the low energy cosmic rays and thus have those abundances, albeit modified, while the model proposed here derives its nuclei near to and beyond the knee from the strongly enriched winds of evolved massive stars.

The difference curve ought to correspond to the galactic contribution of heavy nuclei; this is first of all confirmed by the Fly’s Eye analysis, which does suggest, that in this particle energy range there are very few nuclei of Carbon, Nitrogen, and Oxygen, and many heavier nuclei. Second, over the particle energy range which is well measured, i.e. below a particle energy of about 10 EeV, the curve is well described by a powerlaw with an exponential cutoff near 5 EeV, just as suggested by the earlier cosmic ray arguments (papers CR I, CR II, and CR IV) for the galactic contribution of cosmic rays. At energies above 30 EeV the subtraction is clearly unreliable, since the difference is between two large numbers of similar numerical value.

We conclude first that a detailed check of the existing prediction with the new data now available is successful. The extragalactic contribution can be readily modeled in a) particle energy, b) flux and spectrum, and c) chemical composition.

There is a second important conclusion: Strong shocks in extragalactic jets and their associated hot spots do produce energetic nuclei. This is of interest, since in all arguments about Gamma-ray sources like the quasar 3C279, Mkn 421 and the like (e.g. Hartman et al. 1992a, b, c, Punch et al. 1992), there is always the question whether active galactic nuclei accelerate nuclei at all in their jets. The successful fit made here suggests strongly, that this is the case, and that protons dominate the process. This gives strong support to hadronic interaction models to explain the gamma-ray emission from such quasars such as the model we have proposed (Mannheim & Biermann 1992; Mannheim 1993a, b; also see following section).

4. Radioloud Quasars

This section has been written mostly by K. Mannheim.

Extragalactic radio sources are the largest known dissipative structures (non-thermal objects) known in the universe and as such highly interesting candidate sources of cosmic rays (see references in Berezinsky et al. 1990 and above). There are two ways for them to inject cosmic ray baryons into the intergalactic medium; (i) by escape from hot spots often a few hundred kiloparsecs away from active nucleus of the host galaxy (Biermann 1991, Rachen & Biermann 1992, and papers UHE CR I and
UHE CR II) and (ii) by neutron escape from the sub–parsec scale of the jet or the nucleus (Protheroe & Szabo 1992). Both these mechanisms avoid catastrophic adiabatic losses preventing such escape otherwise. However, do we have any indications that baryons are actually present in radio jets?

The leptonic component ($e^\pm$) of nonthermal particles in jets are the well–studied origin of synchrotron radiation from radio to X-ray frequencies (e.g., see the conference proceedings edited by Maraschi et al. 1989 and Zensus et al. 1987, Bregman 1990). Subtracting thermal emission components like dust infrared radiation (see, e.g., Chini et al. 1989b) and the big blue bump in quasars one is left with a continuum spectrum of highly varying flux and polarization. Flat–spectrum radio sources are characterized by jets oriented at small angles to the line of sight, so that the radiation from the base of the jet is Doppler–boosted towards the observer. Some of these sources, the blazars, show very active behaviour of the most compact regions of the jet, especially with respect to polarization. Probably all flat–spectrum radio sources appear as blazars during active episodes. Since a) particle acceleration takes place in radiojets as is obvious from synchrotron observations, and b) the direct observations of high energy cosmic rays (see Gaisser et al. 1993, and UHE CR II) it is a small step to argue that this acceleration mechanism not only concerns electrons (and positrons), but protons (and nuclei) as well.

4.1. The Hadronic Cascade

Now, recent gamma ray detections and the radio/X-ray correlation of extragalactic flat-spectrum radio sources together with the direct detection of the extragalactic component of cosmic rays (see above) make the existence of a ultra–relativistic proton population in jets very probable. The protons with maximum Lorentz factors in the range $10^9 - 10^{11}$ generate hard photons with energies from keV to TeV via pion and pair photoproduction and subsequent synchrotron cascade re-processing. The target of the energetic baryons is self–consistently provided by the synchrotron photon soup from the primary (accelerated) electrons, while for the secondary pairs cooling always dominates over acceleration. To simplify things the compact jet with relativistic electrons and protons giving rise to the radioloud quasar broadband emission shall be coined the “proton blazar”.

The simplest approach towards considering protons in addition to electrons as radiative constituents of radiojets is to make use of the Blandford & Königl (1979) model, where the jet is assumed to contain a tangled magnetic field dominating (in pressure over radiation) and relativistic particles conically streaming away from the central engine. Within this prescription the physical conditions inside the radiojet plasma can be inferred from continuum observations at radio through optical frequencies. It is then straightforward to calculate secondary emissivities from proton interactions taking into account that the protons will reach much greater energies than the electrons, just because they suffer weaker energy losses at the same energy as electrons.

The pions and pairs created by interactions on the soft target field produce gamma rays to which the jet is optically thick. The reason is that the soft tar-
get photons (from the primary accelerated electrons) are distributed with an inverse powerlaw, so that the higher the gamma ray energy, the more soft photons satisfy the threshold condition for two photon pair production. Thus, a new generation of pairs is produced which then in turn produces gamma rays again (due to synchrotron losses). The crucial parameter is the target radiation compactness \( l \propto L/r \), where \( L \) denotes the soft luminosity and \( r \) the characteristic size of the emitting region, which can also be calculated from the Blandford and Königl model yielding typical values \( l \approx 10^{-4} - 10^{-5} \) for radioloud quasars (extended hot spots have even lower compactnesses). The energy above which the jet becomes optically thick with respect to pair creation is then given by

\[
E^*_\gamma \approx 10Djl^{-1} \text{ MeV}
\]

hence \( E^*_\gamma \approx 1 - 10 \text{ TeV} \) for typical Doppler factors of \( D_j \approx 10 \). This is in accordance with recent GRO and Whipple (Punch et al. 1992) observations. The fact that \( l \ll 1 \) means that the electromagnetic (synchrotron) cascade is unsaturated, i.e. photons at MeV are optically thin. This results in a spectrum which can be described as the superposition of only three to four cascade generations \( \text{pair} \rightarrow \gamma \text{pair} \rightarrow \gamma \ldots \) and the emerging shape of the spectrum can be quite different depending mainly upon the ratio \( \xi = E_{p,\text{max}}/E^*_\gamma \). When \( \xi \) is very great the spectrum has flux indices \( \alpha_X \approx 0.5 \) up to the MeV range and then steepens towards \( \alpha_\gamma \approx 1 \) until it turns over steeply in the TeV range (merged cascade generations at gamma ray energies and the surviving final cascade generation at X-ray energies, cf. 3C279). As \( \xi \) becomes smaller, the X-/gamma ray spectrum becomes decomposed into two bumps, one peaking in the MeV range with indices \( \alpha_X \approx 0.5 \) and \( \alpha_\gamma \approx 1.5 \) and the other peaking in the TeV range with \( \alpha_\gamma \approx 0.8 \) until it turns over steeply. Additional damping of gamma rays may result from strong extended infrared photons fields as in 3C273.

### 4.2. Critical Test: Neutrinos

Hadronic interactions lead to considerable neutrino production of roughly the same luminosity as in the electromagnetic channels. Using gamma ray observations to normalize neutrino fluxes one can readily calculate expected count rates in neutrino detectors. It is crucial to consider that while the gamma ray spectrum reflects cascading, the neutrino spectrum reflects the initial injection. Due to the inverse powerlaw nature of the target agent this means that the neutrino flux density spectrum from \( p\gamma \) interactions is flat. Since the energy loss time scale for \( pp \) collisions is constant, in contrast to the \( p\gamma \) losses described above, the emissivity of secondaries from \( pp \) collisions is much steeper, i.e. reflects the power law of the protons. Therefore the neutrinos from \( pp \) collisions are important at lower neutrinos energies than the neutrinos from \( p\gamma \) collisions, i.e. below \( \approx 10^{-2}E_{\nu,\text{max}} \).

The possibility of detecting neutrinos from flat–spectrum radio quasars with the Fly’s Eye experiment is investigated in Mannheim et al. (1992). There is a significant chance that the HiRes (Cassiday et al. 1989) might be able to detect proton blazars. The neutrino flux crosses the atmospheric background flux at roughly
\( E_\nu \approx 1 \text{ PeV} \) and has a spectrum \( F_\nu \propto E_\nu^{-1} \) below approximately the same energy and is flat above up to roughly \( 10^9 \) GeV.

Neutrino detection at lower energies is feasible with experimental designs currently under construction like AMANDA, DUMAND II or BAIKAL (Stenger et al. 1991, Halzen 1991) in the TeV range and by advanced analysis of horizontal atmospheric showers up to the PeV range (Halzen & Zas 1992). Although extragalactic neutrino beams would be a unique tool for elementary particle physics to study weak interactions in a kinematical regime far beyond planned laboratory designs, it remains to be shown whether it is possible to detect the neutrinos from radioloud quasars below PeV because of the flatness of the flux.

4.3. Further Consequences

Hadronic interactions exhibit still another fingerprint: In every second inelastic proton-photon event the leading nucleon emerging from the interaction fireball is a neutron. Since neutrons are not magnetically confined anymore, they can escape the acceleration region in the jet (Biermann & Strittmatter 1987). The optical depth for neutrons transverse to the jet is less than unity, so that the ultra–relativistic neutrons can leave the source without turning back into confined protons (provided that the distance of the jet from the site of the thermal UV nucleus is great enough to prevent damping by these photons). Since the frame in which the scattering centers isotropize the accelerated protons rests in the bulk flow of the jet, the emitted neutrons as seen from an observer stationary with respect to the host galaxy are streaming along a cone with the blazar in its apex. As a corollary it follows that *Baryonic blazar beams contain energetic photons, neutrinos and neutrons with comparable luminosities.*

It remains to be shown, whether the neutrons have observable consequences (cf. Kirk & Mastichiadis 1989).

While it has been suggested by Sikora et al. (1989) and Begelman et al. (1991) (cf. MacDonald et al. [1991] for further consequences of energetic neutrino production.) that neutrons from a hypothetical accretion disk in an AGN can explain observed gas outflows, the mechanism proposed works best for moderately energetic neutrons. To convert the power of the relativistic neutrons into kinetic power of thermal plasma authors assume effective coupling of the two media via excitation of Alfvén–waves by \( \beta–\) decay protons.

However, it is by no means clear how the microphysics shall interplay to lock the particles isotropically to the plasma when their energy is very high (e.g., Berezinsky et al. 1990). It must be remembered that the luminosity of neutrons from the jet peaks at the highest neutron energy (\( \approx 10^9 \) GeV). Without such isotropic locking there are no adiabatic losses. Moreover, the distance neutrons travel before suffering \( \beta–\) decay is given by

\[
R_n = \gamma_n c \Delta \tau_n \approx \left( \frac{\gamma_n}{10^8} \right) \text{kpc}
\]

Proton blazars generating most neutron luminosity at \( \gamma_n \geq 10^9 \) are therefore clearly injectors of cosmic rays, because the neutrons decay at \( R_n \geq 10 \) kpc outside
the main central galaxy. Only for $\gamma_n < 10^9$ neutrons decay well within the host galaxy. The greatest number of these neutrons come from $pp$ collisions in the jet. Their luminosity is much less than that of neutrons from $p\gamma$ collisions. It must also be remembered that the great luminosity of the proton blazar appears only in the direction of the jet. After isotropization the true luminosity available for conversion into the kinetic power of a wind is reduced by the factor $D_j^{-4} \approx 10^{-4}$.

Other models based on $pp$-interaction in a thick accretion disk (Bednarek 1993), on electron/positron beams (Henri et al. 1993), and on synchrotron and inverse Compton processes (Maraschi et al. 1992, Dermer & Schlickeiser 1992) have been presented. Hadronic interactions distinguish themselves most strongly through the neutrino emission from the more classical electron/positron arguments; also, the observed high energy cosmic ray particles strongly suggest that protons and heavier nuclei get abundantly accelerated in the cosmos (see earlier sections in this review). The basic hadronic interaction processes are being redesccussed in some detail by Berezinsky & Gazizov (1993a, b).

To summarize our model, what can be learned from adding relativistic protons to the usual relativistic electrons observed in compact radio jets?

First of all, if the protons manage to be accelerated up to the high energies where their energy loss rate is equal to the electron energy loss rate, they generically induce high energy emission from X–rays to gamma rays. The spectrum is complex, since it results from the superposition of several cascade generations. These various cascade generations allow a variety of source spectra (e.g. 3C279 vs. 3C273). High redshift sources suffer from interaction with a possible infrared background (Stecker et al. 1992), and so their observable photon energy range is limited.

At present it is unknown, whether there is also a generic proton/electron ratio and hence a generic ratio $L(> X)/L(< UV)$. The EGRET detections of flat–spectrum sources seem to indicate that either there is significant scatter in the distribution of the proton/electron ratio or that the proton acceleration (taking longer than the electron acceleration) is sometimes interrupted before reaching the maximum proton energies.

Secondly, the presence of protons leads to neutron production which has remarkable astrophysical consequences. At ultra–high energies the neutrons escape from the host galaxy without adiabatic losses injecting cosmic ray protons. At lower energies they can accelerate gas surrounding the radio jet within the host galaxy (and its halo).

Finally, with protons brought into the game, the flat–spectrum radio sources generate a diffuse background of high energy neutrinos which is detectable with planned experiments.

Important next steps are here first to try to understand the rapid time variability observed in such sources (Qian et al. 1991, Quirrenbach et al. 1991, 1992), and second to calculate the induced $\gamma$-ray background, both due to the interaction of cosmic ray particles interacting with the microwave background, as well as the
superposition of all flat radio spectrum quasars.

5. The Galactic Center and Radioweak Quasars

5.1. The Galactic Center

The recent radio observation of the Galactic center source Sgr A* with mm-VLBI by Krichbaum et al. (1993) for the first time resolved the source spatially; these observations suggest an elongated one-sided morphology quite similar to the jet observed in radio quasars. Fitting the flux density and the radio spectrum of this central source with the simple classical model for extragalactic radio sources (Blandford & Königl 1979), we can deduce limits for the important ratio of jet power to disk accretion luminosity (Falcke et al. 1992, 1993). It turns out that the data lead to a lower limit for this ratio of order 0.1. This suggests that even for such a weak source akin to an AGN, the jet power is of similar order of magnitude as the observed accretion disk luminosity. As an important next step it will be interesting to see whether this conclusion holds up for radioweak quasars in general; below we will speculatively assume that it does.

The feeding of the central accretion disk close to the black hole from larger scales further out, out to about 100 pc can be thought of as an overall nonsteady accretion flow (cf. Weizsäcker 1943, 1951, Lüst 1952) with a substantial effective viscosity in the interstellar gas (Linden et al. 1993, Biermann et al. 1993). The accretion rate for the central disk is very low (Falcke et al. 1993), while on the larger scale the cloud motions can be traced out very well with such an accretion disk model, which has a very much higher accretion rate, more in correspondence with the star formation rate in that area.

If there is a powerful jet, possibly emanating from the inner edge of the accretion disk, then nuclei might be accelerated to high particle energies in shocks at the boundary or inside. The effect of energetic particles on the accretion disk again, and also on molecular clouds will be considered below.

5.2. Radioweak Quasars

5.2.1. The Hadronic Cascade in the Accretion Disk

The active nuclei of galaxies (AGN), ranging in luminosity from Seyfert galaxies to quasars, are the most powerful individual sources of radiation in the Universe. To explain the power emitted by such objects, one generally assumes the existence of a central engine in which the gravitational energy of matter falling into a supermassive black hole gets converted into radiation. Even though so far only electromagnetic radiation has been observed, it is generally assumed that other particles are accelerated and emitted as well to explain the tight relationship between the non-thermal and thermal components in the IR, UV and X-ray spectra (Lightman & White 1988, Pounds et al. 1989, Fabian et al. 1990, Clavel et al. 1992, Chini et al. 1989a). Of special interest are neutrinos, since they can travel cosmological distances without
losing the information on the direction they originated from. Large underwater detectors (Alatin et al. 1991, Stenger 1991) or detectors in the antarctic ice cap (Barwick et al. 1991), are used as neutrino telescopes. Recent calculations have shown that the flux of neutrinos originating in an AGN could be detected by such experiments (Stecker et al. 1991, 1992a, 1992b, Szabo and Protheroe 1992, Mannheim et al. 1992, Biermann 1992, Mannheim 1992). Data from proton decay experiments (Meyer 1991) and airshower arrays Halzen & Zas (1992) is already sensitive enough to constrain such models significantly.

In the following we are interested in the production of neutrinos in radio-weak AGN. In the “standard” model for AGN one assumes that the central black hole is surrounded by an accretion disk of infalling matter. Besides that, one expects to find bipolar outflow of gas and plasma perpendicular to the disk (jet). Jets can be seen in different objects with disk accretion and in many AGN (Bridle & Perley 1984, Meisenheimer & Röser 1991). One expects shocks in the plasma of the jets (Mannheim 1992) which could accelerate protons through first order Fermi acceleration to energies ranging from $10^6$ GeV to $10^9$ GeV, depending on the distance to the black hole, with a powerlaw spectrum of $E^{-2}$ (Biermann & Strittmatter 1987). The observation of $\gamma$-rays with the same kind of spectrum (Punch et al. 1992) indicates that indeed there is shock-acceleration outside the core of the AGN, so that the photons can escape.

Niemeyer (1991) showed that the far infrared (FIR) emission of AGN can be explained by assuming that accelerated particles diffuse back from a central location above the accretion disk and heat dust clouds beyond the outer region of the accretion disk (see the next subsection). In the inner part of the disk, protons hitting the disk will initiate hadronic cascades through interactions with the gas in the disk. This could feed into an electromagnetic cascade and thereby generate the observed X-ray and gamma emission (Biermann 1992). Clearly, this is a speculative picture at present, but it is suggested to have merit on the basis of the analogous situation in the center of our Galaxy, discussed briefly above.

To describe the accretion disk, we use the model by Shakura & Sunyaev (1973). For our purposes, we will concentrate on the inner region, which is the radiation dominated part of the disk (neglecting the thin, innermost ring around the black hole).

Due to the observed short timescales of the variability of the X-ray and UV components of the AGN spectrum, this part of the electromagnetic radiation must predominantly originate in a small, central region around the central black hole (Stecker et al. 1991, 1992a, b) with a typical mass of $10^8 M_\odot$ for luminous AGN. We expect neutrino production to take place in the same region, so we need to know the incoming particle flux in this region. From Niemeyer & Biermann (1993 and below) we know that the FIR spectra are very well fitted by assuming that the protons originate from a source in the jet extending outward from $z_0 = O(3 \cdot 10^{15} \text{cm}) \approx 100R_\odot$.

For a source that close to the center of the disk, the discussion of the maximum proton energy in (Biermann & Strittmatter 1987) has to be extended to take the UV photons from the disk into account. The photon spectrum has a pronounced bump in the UV region (Stecker et al. 1991, 1992a, b, Clavel et al. 1992). Therefore,
the reaction $p\gamma \rightarrow N\pi$ is only effective for protons with energies above the threshold $E_{th} \approx 8 \cdot 10^6$ GeV for the interaction with UV photons. This interaction channel limits the maximum proton energy.

If we take $\dot{M}/M_{edd} \approx 0.1$ and $\alpha \approx 0.1$, the column density perpendicular through the disk is high enough compared to the mean free path for $pp$- and $np$-collisions of $O(50 \, \text{g/cm}^2)$ for a hadronic cascade to develop. The magnetic field confines the protons to the disk, which leads to a further increase of the effective column density seen by the protons. Even though neutrons which are produced in the cascade are not confined by the magnetic field, the amount of gas present is sufficient to prevent them from escaping; they interact before they can decay.

The hadronic cascade in the disk is much simpler than cascades in the earth’s atmosphere (Gaisser 1990, Gaisser & Yodh 1980, Gaisser et al. 1978), since the density of the accretion disk is so low that all unstable particles decay rather than interact. Therefore the cascade consists of a nucleonic part which feeds into the mesonic and electro-magnetic channels; no pion-nucleon reactions occur.

The electromagnetic output of the hadronic cascade is reprocessed in the inner disk. Pair cascades, inverse Compton scattering, and reflection on cold material change the shape of the initial $E^{-2}$-spectrum and provide a steep turnover around $E_\gamma \approx m_e \approx 511$ keV (Zdziarski & Coppi 1991, Zdziarski et al. 1990, Ghisellini 1987, Svensson 1987, Zdziarski et al. 1991, Done et al. 1990). This may produce the observed X-ray and gamma emission of the AGN which we will calculate in detail elsewhere.

The neutrino spectrum from a single source mirrors the $E^{-2}$-spectrum of the protons, only that it is shifted down by a factor of $\approx 0.05$. The upper cutoff of the neutrino spectrum depends on the details of the shock acceleration process in the jet, since that determines the maximum energy reached by the protons (Biermann 1992).

To be able to use the relation between electromagnetic output and neutrino production to predict the diffuse neutrino background, we have to estimate how much of the observed, diffuse X-ray and $\gamma$-ray background results from hadronic cascades in AGN.

There is growing evidence from the ROSAT all sky survey that an appreciable fraction of the background is due to active galactic nuclei, possibly a dominant proportion (Hasinger et al. 1991).

To be conservative, we allow for a factor of three maximum between the total hard X-ray emission from an AGN galaxy (starburst and active nucleus together) and the contribution strictly from the hadronic cascade.

The observations of the hard X-ray background at those photon energies minimally influenced by reprocessing (Fabian et al. 1990), i.e., at energies where the original $E^{-2}$ powerlaw is still visible, give a possible range for the unprocessed energy total of $1.0 \cdot 10^5$ to $1.4 \cdot 10^5 \, \text{eV cm}^{-2} \, \text{s}^{-1} \, \text{sr}^{-1}$. Using the above mentioned conservative estimate, we arrive at a lower limit of $3 \cdot 10^4 \, \text{eV cm}^{-2} \, \text{s}^{-1} \, \text{sr}^{-1}$ for the contribution of hadronic cascades to the X-ray and $\gamma$-ray background.

Since the cosmological redshift is the same for neutrinos and photons, we can scale the background neutrino luminosity using the fraction of the X-ray and $\gamma$-ray
background derived above. We get

$$N(E_{\nu}) = 1.7 \cdot 10^{-12} \left( \frac{E_{\nu}}{\text{TeV}} \right)^{-2} \text{cm}^{-2} \text{s}^{-1} \text{sr}^{-1} \text{GeV}^{-1}$$

(71)

as a conservative limit for the sum of all neutrino species. About 2/3 are muon neutrinos, the remaining 1/3 are electron neutrinos. This prediction is a factor of 2.5 lower than the experimental limit set by the Frejus experiment (Meyer 1991). For each family, the number splits evenly into neutrinos and anti-neutrinos. We expect the background spectrum to have a softer cutoff at high energies than a single source, since the maximum neutrino energy and the redshift vary between different AGN. Therefore, the cutoff in the superposition of all spectra will be smoothed out.

The main point in our model is the production of neutrinos as the result of a hadronic cascade initiated by $pp$-interactions. Compared to the $p\gamma$-channel, which has a threshold of $E_{\text{th}} \approx 8 \cdot 10^6 \text{GeV}$ for photoproduction on UV photons, all protons above a few hundred MeV contribute in $pp$-interactions. As a consequence, we expect neutrino emission even from sources with a low cutoff in the primary proton spectrum. This way, we include contributions from a larger class of sources both to the neutrino background and to the diffuse X-ray background. Similar to the result of other authors (Stecker et al. 1992a, b, Szabo & Protheroe 1992), our model displays the tendency to produce neutrinos strikingly close to existing limits (Meyer 1991, Halzen & Zas 1992).

Modifications of this $pp$-model can lead to a reduced neutrino flux while re-processing of the electromagnetic component ensures an unchanged X-ray emission. The most drastic modification — reducing the maximum proton energy and correspondingly the maximum neutrino energy — can ultimately decrease the observable extragalactic neutrino flux by moving the cutoff below the cross-over with the steeper spectrum of atmospheric neutrinos. More realistically, steepening of the AGN proton distribution at a sufficiently low break energy is an alternative to reduce the predicted event rate in a neutrino detector. In both cases, neutrino production via $p\gamma$-reactions becomes ineffective.

On the other hand, a proton spectrum flatter than $E^{-2}$ leads to the dominance of $p\gamma$-reactions at high energies. This, again keeping the X-ray background unchanged, enhances the neutrino flux far above the TeV range while reducing the flux below. Details depend on the modelling of the target photon field (i.e., contributions from the disk, corona, jet, etc.).

A likely consequence is the effect of the energetic particles (see, e.g. also, Ferland & Mushotzky 1984, Kazanas & Ellison 1986) and photons on stars in their neighborhood of the central engine, i.e. the rotating central black hole, the inner accretion disk and the inner jet. The stellar distribution is modified by disruption of those stars that get too close to the central black hole, by stellar collisions, and possibly also by agglomeration and interaction with the disk. Stellar atmospheres exposed to a strong flux of energetic particles may expand and thus start a stellar wind, which in turn gets ionized by the ultraviolet radiation field, a process which can be described as an HII-region turned inside out (Scoville & Norman 1988, Kazanas...
a similar effect can occur when energetic neutrinos penetrate the stars and change their structure from the inside, bloating the stars and thus moving lower mass main sequence stars into the red giant region (MacDonald et al. 1991). An additional process which may confuse the observations is the likelihood of star formation in the dense molecular clouds which are observed to exist rather close to the central engine, e.g. by absorption (Lawrence & Elvis 1982, Mushotzky 1982, Antonucci and Miller 1985, Barthel 1989) or by maser emission (e.g. Claussen & Lo 1986)). The broad emission line region may thus be the ensemble of slow stellar winds embedded in a strong ultraviolet radiation field.

An important point in our speculative model is that the acceleration of the protons and other nuclei takes place above the disk. This way, the acceleration takes place in an environment more suitable than an accretion disk.

5.2.2. The farinfrared emission of radioweak quasars

The FIR emission from active galactic nuclei was believed for a long time to be due to nonthermal processes (e.g. Edelson 1986). The observations of a few nearby Seyfert galaxies (Hildebrandt et al. 1977, Telesco & Harper 1980, Rieke & Low 1975, Roche et al. 1984, Neugebauer et al. 1979, Miley et al. 1984) gave the first hint, that dust emission might be important as an additional source. Then, the very sensitive mm-observations at the 30m IRAM telescope (Chini et al. 1989a, b) combined with the IRAS quasar survey (Neugebauer et al. 1986) turned the trend; the dust emission from radioweak quasars is now well accepted to be generally due to warm dust (Sanders et al. 1989, A. Lawrence et al. 1991, Barvainis 1990).

There are theoretical possibilities to reproduce many, but not all of the observed spectra with synchrotron emission due to two populations of relativistic electrons via self-absorption and emission (de Kool et al. 1989, Schlickeiser et al. 1991). However, the observations of very large amounts of molecular gas in radioweak quasars (Barvainis et al. 1989, Alloin et al. 1992) demonstrated that the radioweak quasars contain certainly sufficient gas and dust mass to produce all the emission observed in the FIR. Hence, our understanding of radioweak quasars now includes an extended region of dust clouds, presumably in a disk configuration. It is less clear what the energy source of this dust is. The difficulty lies in the fact that the FIR emission of radioweak quasars is the strongest of any of the observed wavelength bands usually, and therefore any primary source must be even stronger, allowing for some waste and inefficiency. Another difficulty is, that we can infer from Plancks law that the radial scale from which the farinfrared emission is arising, ranges over a fairly large radial scale out to distances of order 100 pc.

A natural hypothesis, which is actually supported by the discovery of large amounts of molecular gas, is that a circumnuclear starburst also provides the FIR emission seen. It appears as a very reasonable hypothesis that stars are being formed in these molecular clouds surrounding the central engine; a number of nearby Seyfert galaxies and galaxies with extreme FIR luminosities clearly demonstrate the activity of such starbursts (a now well established example is the Seyfert galaxy NGC1068, Wilson et al. 1992). However, starbursts have a very clear correlation between
their FIR and radio emission, and also have typical radio spectra: Normal galaxies have nonthermal radio emission spectra with flux density \( S_\nu \sim \nu^{-0.8\pm0.1} \), sometimes modified by some thermal free-free emission, or by free-free absorption in a clumpy interstellar medium, as in the famous starburst galaxy M82 (Kronberg et al. 1981, 1985). The ratio of FIR to radio emission is much larger in many cases for radioweak quasars than for starburst galaxies, clearly requiring an additional energy source (Chini et al. 1989a, 1989b). The task, however, remains and is an important next step, whether extremely intense regions of star formation also suppress the radio emission relative to the farinfrared emission; resolution of this question may come from infrared line emission observations. Finally, in some cases the spatial structure of the radioemission has been determined, and is found to be analogous to radio jets of low power (e.g., Lacy et al. 1992), again inconsistent with starbursts as the source of the dominant radioemission. Other possibilities are discussed in Niemeyer & Biermann (1993), like the X-ray heating model of Sanders et al. (1989); in such a model rather extreme warping of the layer of dust clouds is required to geometrically allow irradiation of all the dense gas. Here we introduce a hypothesis that circumvents geometric difficulties, but uses a putative jet and its particle acceleration:

We propose that the dusty clouds are heated by relativistic particles. These particles are thought to originate from a source on the axis of rotational symmetry, possibly a jet or knots in a jet and its interaction regions with the environment. The difference between radioloud and radioweak quasars in our picture is not that the jets in radioweak quasars have a low kinetic luminosity (Lacy et al. 1992), but that their radio emission is very much weaker, despite their large power. This is analogous to the weak nonthermal radio emission of Wolf Rayet stars as compared to the large nonthermal radioemission of radiosupernovae, which differs probably only in the velocity of the shockwave which causes the particle injection and acceleration (see Biermann & Cassinelli 1993). Thus, in our speculative picture radioweak and radioloud quasars both have jets of strong power, but in one case the jet power is dissipated fairly deep inside the host galaxy, and in the other case the dissipation is very far outside the host galaxy. The source is assumed to be located above the disk and to transform a fair fraction of the entire source accretion power into relativistic particles, analogous to the argument made above for the Galactic center source. Such a concept is consistent with the model for strong jets discussed in the literature. The particles easily scatter in the magnetic plasma that permeates the halo above the disk all the way from the central engine. And since dust clouds of the column densities typical for central regions of galaxies readily provide sinks for energetic particles, they can be heated.

We calculated (Niemeyer & Biermann 1993) the FIR spectra that can be produced in this model, and compare them to observations. The basic model parameter is the radial behaviour of the diffusion coefficient that governs the transport of the energetic particles from the putative central source out to the dust clouds in the disk. The essential test is then whether any reasonable model for this diffusion gives a physical explanation for the observed spectra.

Independent of any particular model for the diffusion coefficient, we assume
here for the diffusion coefficient a general powerlaw dependence on spherical radius \( r \). It can be shown that the spectrum itself is independent of the numerical value of the diffusion coefficient, since we have assumed a stationary state (see Niemeyer & Biermann 1993). However, for these basic assumptions to be justified, that lead to a stationary state, the diffusion coefficient has to be quite large. We consider also the possibility that the source is extended along the symmetry axis \( z \); this case can be considered as a limiting addition of many point sources. The physical concept is given by many knots in a jet.

Dust particles are of mixed chemical composition with different grain size distribution. Mathis \textit{et al.} (1977) developed a dust model for the composition, absorption and emission of interstellar dust. They succeeded to approximate the dust grain size distribution with a powerlaw. This powerlaw can be interpreted as a quasistationary fragmentation (Biermann & Harwit 1980).

We do not discuss in detail here how the energy is transformed from an impinging flux of energetic particles into an effective gas and dust heating. There are two channels which are clearly important.

a) The \textit{pp} or nucleus - nucleus collisions which transform particle energy into pions, which in turn decay and finally deposit their energy in electrons and positrons as well as \( \gamma \)-photons of a range of energies. These electrons/positrons and \( \gamma \)s in turn thermalize their energy by further encounters with atoms and their shells (see, \textit{e.g.}, Spitzer & Tomasko 1968). In this picture about 1/3 of the primary proton and nuclei energy density is dissipated, the rest goes into neutrinos and energetic photons.

b) The second process is the excitation of Alfvén waves by a gradient of the particle distribution, which is necessarily formed when many particles are absorbed in \textit{pp} and corresponding collisions (Skilling 1975a, 1975b, 1975c, Skilling & Strong 1976). This wave excitation in turn can lead to a very strong wavefield, which dissipates readily in a neutral-ion plasma like inside dense dust clouds; it would dissipate considerably less in a purely ionized medium, such as the innermost accretion disk. In molecular clouds of a low degree of ionization this diminishes the resulting \( \gamma \)-emission due to an appreciable optical thickness in the hadronic encounters. An important next step is to work out the consequences for the molecular clouds in the Galactic center, where there are unresolved questions regarding the flux of cosmic rays and their confinement in the galactic magnetic field (\textit{e.g.} Chi & Wolfendale 1993), the heating of clouds and the amount of molecular gas. In this case, a large proportion of the primary energy density of the protons and nuclei can be dissipated.

Here we simply assume that a fixed fraction of the energy of the impinging particles is deposited as heat. We determine the dust temperature by assuming that emission and absorption balance.

Since the incoming flux of energy depends on radius in the disk geometry considered, the dust temperature is also dependent on radius \( r \). We calculate the FIR spectrum of a finite disk in the wavelength range 10 und 1300 \( \mu \)m. The most important characteristics to be interpreted are the rapid rise from the mm-wavelengths to the FIR \( L_\nu \sim \nu^\alpha \) with \( \alpha \geq 2.5 \), and, on the other hand, the slow decrease with frequency beyond the farinfared with a spectral index \( \alpha \simeq -1 \).
All cases considered of the radial behaviour of the diffusion coefficient show the steep increase from the mm-range towards the FIR, as required by the observations. The emission through the IR towards the NIR decreases very slowly. The general trend of the calculated spectra show a good representation of what is seen in the observations. In the mm to FIR wavelength range the local spectral index $\alpha$ is always $\geq 2.5$, independent of the specific dependence of the diffusion coefficient on radius. On the other hand, the spectra in the IR to NIR range are indeed strongly dependent on the source distribution function and on the specific law for the radial dependence of the diffusion coefficient chosen. The decline through the IR gets stronger with an increasing radial dependence of the diffusion coefficient.

It is important to note that there is no minimum distance the source has to have. This is because with decreasing minimum distance more particles are deposited in the inner part of the disk, but the radial dependence of the temperature remains the same in the outer part of the disk. The maximal IR–luminosity decreases with decreasing minimal distance $z_{\text{min}}$. The relation can be described by a logarithm scale factor: If the minimal distance $z_{\text{min}}$ decreases by a factor $X$, the IR–luminosity decreases by a factor $Y = \ln X$ for a given source power.

The calculations show that for sufficiently high values of the column density the optical depth at high infrared frequencies is significantly larger than unity. Hence the limiting case of high optical depth describes the spectrum, and the numerical value of the optical depth itself drops out. Correspondingly we treat the case of optical depth much smaller than unity at low frequencies for all values of column densities considered. In this case the actual value of the optical depth does matter, since the emission is directly proportional to it. Specifically, the dependence of the absorption cross section on frequency can be approximated in the FIR/mm-range by a power law. For extremely low values of the column density the optical depth becomes smaller than unity at all frequencies, and the local maxima and minima of the absorption cross section then determine the shape of the emission spectrum directly and produce the various maxima and minima that we described. Should some observed spectra show such features, then we could derive estimates for the column density.

Nearly all observed spectra can be fitted with a line source with an intensity $z^{-1}$. Therefore a basic model is a line source, with an intensity $z^{-1}$ starting at some indeterminded distance above the the center of symmetry - as long as this distance is not larger than of order 0.1 pc - and a diffusion coefficient which is $D \sim r$ or $\sim r^2$.

The source energy goes about equally into the central region, where the disk is too hot for dust, and the outer region, where infrared emission is caused. Of the fraction of energy which goes into the outer region, we have used $1/3$ to heat the dust and dense gas specifically, but, depending on the microphysics assumed (see the discussion above), this fraction could approach unity.

A line source with source strength proportional to $z^{-1}$ can be understood as due to the particle population energy density $\sim z^{-2}$, a lateral gradient scale $\sim z$, a transport coefficient $\sim z$ also, and a circumference $\sim z$; again we use here a scaling argument for the transport coefficient.

To summarize we find that for reasonable parameters we can indeed reproduce
the observed farinfrared spectra of radioweak quasars. We find that for a diffusion coefficient in the region above the disk, which scales with radius \( r \), and a line source strength \( \sim \nu^{-1} \), with a total source luminosity of 6 to 10 times the observed infrared luminosity, we can reproduce and interpret the spectra of radioweak quasars from the mm to the near infrared region.

There are obvious consequences of our model for the X-ray, gamma- and neutrino emission from the inner disk, as well as likely consequences for the heating of interstellar clouds close to a central engine (by wave excitation and dissipation), which remain to be worked out. The analogy to the Galactic center also suggests that the accretion disk luminosity crudely scales with the jet power, and with the FIR emission, but maybe larger by a factor in the range 1 to 10. Such correlations will provide a severe test for the proposed speculative model.

6. Conclusions and Outlook

Energetic particles can have an observable effect in many astronomical objects, from heating and ionization of interstellar material to TeV \( \gamma \)-rays from Active Galactic Nuclei (AGN). Key to all such arguments and an important test for our physical insight is the acceleration of the observed cosmic ray particles, their energy range, spectrum and chemical composition. There has been recent progress in this area and in this review we have described some of it.

6.1. Galactic Cosmic Rays

We propose a fully developed physical concept to account for the cosmic ray particle energies, spectrum and chemical abundances. We suggest, that the cosmic rays observed near earth arise from three sites:

1) Supernova explosions into the nearly homogeneous interstellar medium; these cosmic rays reach particle energies of about \( 100Z \) TeV, and have an injection spectrum of near \( E^{-2.4} \), where \( Z \) is the charge of the nucleus considered.

2) Supernova explosions into a former stellar wind produce a slightly flatter spectrum of near \( E^{-2.3} \) at injection below a bend near \( 700Z \) TeV, beyond which the spectrum steepens to near \( E^{-2.7} \) up to a maximum energy of about \( 100Z \) PeV. This means that heavy nuclei dominate at large particle energies. Since here the acceleration is in a shock that runs through the wind of an evolved star, the heavy elements are enriched. The heavy element contribution may even be important at low energies. Air shower data have been successfully fitted with this model in the energy range from 10 TeV to near EeV.

3) The hot spots of powerful radio galaxies provide the third site of origin. The predictions from this model prediction have also been fitted successfully to airshower data from 0.1 EeV to 100 EeV.

We emphasize that the proposal rests on a plausible but nevertheless speculative assumption about the nature of the transport of energetic particles in perpendicular shock waves, namely that there are large fast convective motions across the shock interface; this notion is, however, supported by radio polarization observations.
The model has predictive power. The predictions for both the galactic and the extragalactic cosmic rays underwent their most severe tests with the airshower data quite successfully, with all the parameters fitted rather close to their predicted values, and with little "wiggle room".

The leakage of low energy cosmic rays from normal and starburst galaxies at high redshift may ionize and heat the intergalactic medium, meeting all known constraints.

6.2. Extragalactic Cosmic Rays

The third site of cosmic ray production is the hot spots of powerful radio galaxies. The full tests against the new airshower data from Fly's Eye have been made, demonstrating that this concept successfully explains particle energy, spectrum and chemical composition in the particle energy range from 0.1 EeV to 100 EeV. One consequence is that the mean free path for high energy cosmic rays in the intergalactic medium has to be of order the bubble size inferred from galaxy distribution studies, and thus of the same order of magnitude as the distances between powerful radiogalaxies.

6.3. Radioloud quasars

The new GRO observations of such flat spectrum radio quasars demonstrate, that they have higher luminosities in the $\gamma$-ray range than at any other wavelength, with one object even showing emission near TeV photon energies.

We review the hadronic interaction picture to explain these observations, tracing the $\gamma$-ray emission to the consequences of hadronic interactions of nuclei accelerated in shocks in relativistic jets. Our model is capable of explaining a variety of hard X-ray to $\gamma$-ray spectra such as observed by Comptel and Egret. The spectra deduced have some interesting analogies to the spectra of $\gamma$-ray bursts, reviewed in chapter 3 of this book.

6.4. The Galactic Center and Radioweak quasars

The Galactic center has been demonstrated by recent mm-VLBI observations to contain an elongated morphology in its compact radio emission suggestively similar to jets in radioloud quasars; in simple models this jet can be shown to carry nearly as much power as is in the accretion disk luminosity. For radioweak quasars it has only recently become clear (Chini et al. 1989a) that their farinfrared emission is due to dust, and is thus thermal radiation. The origin of this energy is not clear; several possibilities exist, and one of these possibilities is the hypothesis, that the active nucleus provides a large proportion of its power in the form of energetic particles - arising from shocks in a putative jet analogous to our interpretation of the Galactic center - which in turn heat the molecular clouds which contain the dust. Various consequences of this hypothesis have been described, including some consequences for the emission from the inner accretion disk (X-rays and neutrinos) as well as the farinfrared emission spectra of quasars.
6.5. The 12 next steps

The comprehensive physical concept proposed here for AGN and Galactic sites of origin of cosmic rays stands and its essential parameters are fixed. There remain many important steps to be done and so the proposal can be understood as a program; we will list here those steps that seem feasible within the near future:

1. The secondary to primary ratio of the nuclei in the cosmic rays and their relation to the structure and evolution of molecular clouds.
2. A simulation of the selection effects governing our observations of normal Sedov-type supernova remnants.
3. A comparative study of different acceleration sites for energetic electrons, like supernova remnants, radiosupernovae, novae, stars in order to induce a physical concept for the critical velocity of electron injection. Obviously, the injection of protons and heavier nuclei ought to be studied in parallel.
4. A check on why the knee feature of the cosmic ray spectrum is so sharp. This sharpness suggests either a) a complete convergence of the properties of massive stars, including rotation and magnetic fields, or b) a speculative picture like the one suggested in paper CR I that traces the origin of the mechanical energy of supernovae to the potential energy at that depth where the innermost section of the star forms an accretion disk for a few seconds, before magnetic fields transport its angular momentum away and most of the disk collapses further.
5. The abundances of the chemical elements in the cosmic rays at low energy.
6. A study of the supernova remnant candidates in the starburst galaxy M82.
7. The effect of low energy cosmic rays from normal and starburst galaxies at high redshift on Lyman $\alpha$-clouds.
8. A calculation of the $\gamma$-ray background due to cosmic ray interaction.
9. The rapid variability of blazars, as inferred from the same shock acceleration and hadronic interaction concept, which successfully lets us understand the $\gamma$-ray emission, its variability and spectrum.
10. The connection between radioweak quasars and our Galactic center in structure and energetic photon emission.
11. The contribution to the farinfrared emission of radioweak quasars from the circumnuclear starburst.
12. The heating of giant molecular clouds in our Galactic center, their $\gamma$-ray emission and the cosmic ray flux near to the Galactic center.

Obviously, the most important task of all is to check the notion of the fast convective parallel transport at shocks, both observationally and computationally.

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