ORIGIN, INJECTION, AND ACCELERATION OF CIR PARTICLES: THEORY

Report of Working Group 7

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Received: 8 March 1999; Accepted: 28 June 1999

Abstract. On the basis of the observational picture established in the report of Mason, von Steiger *et al.* (1999) the status of theoretical models on origin, injection, and acceleration of particles associated with Corotating Interaction Regions (CIRs) is reviewed. This includes diffusive or first-order Fermi acceleration at oblique shocks, adiabatic deceleration in the solar wind, stochastic acceleration in Alfvén waves and oblique propagating magnetosonic waves, and shock surfing as possible injection mechanism to discriminate pickup ions from solar wind ions.

1. Introduction

The basic features of energetic particles associated with CIRs have been known since the 1970s from Pioneer and Voyager observations (McDonald *et al.*, 1976; Barnes and Simpson, 1976) and are summarized in the accompanying paper by Mason and Sanderson (1999). Figure 1 shows, as a typical example, more recent measurements of various particle data and of the magnetic field magnitude during the passage of a CIR by Ulysses. From top to bottom are shown the electron and proton fluxes in the range 0.1–0.4 MeV and 0.8–1.0 MeV, respectively, and the solar wind velocity, temperature, density and magnetic field strength. The energetic protons exhibit two reasonably well-resolved peaks centered approximately on the forward and reverse shock. The intensity increase of the protons near the reverse

Space Science Reviews **89:** 369–399, 1999.

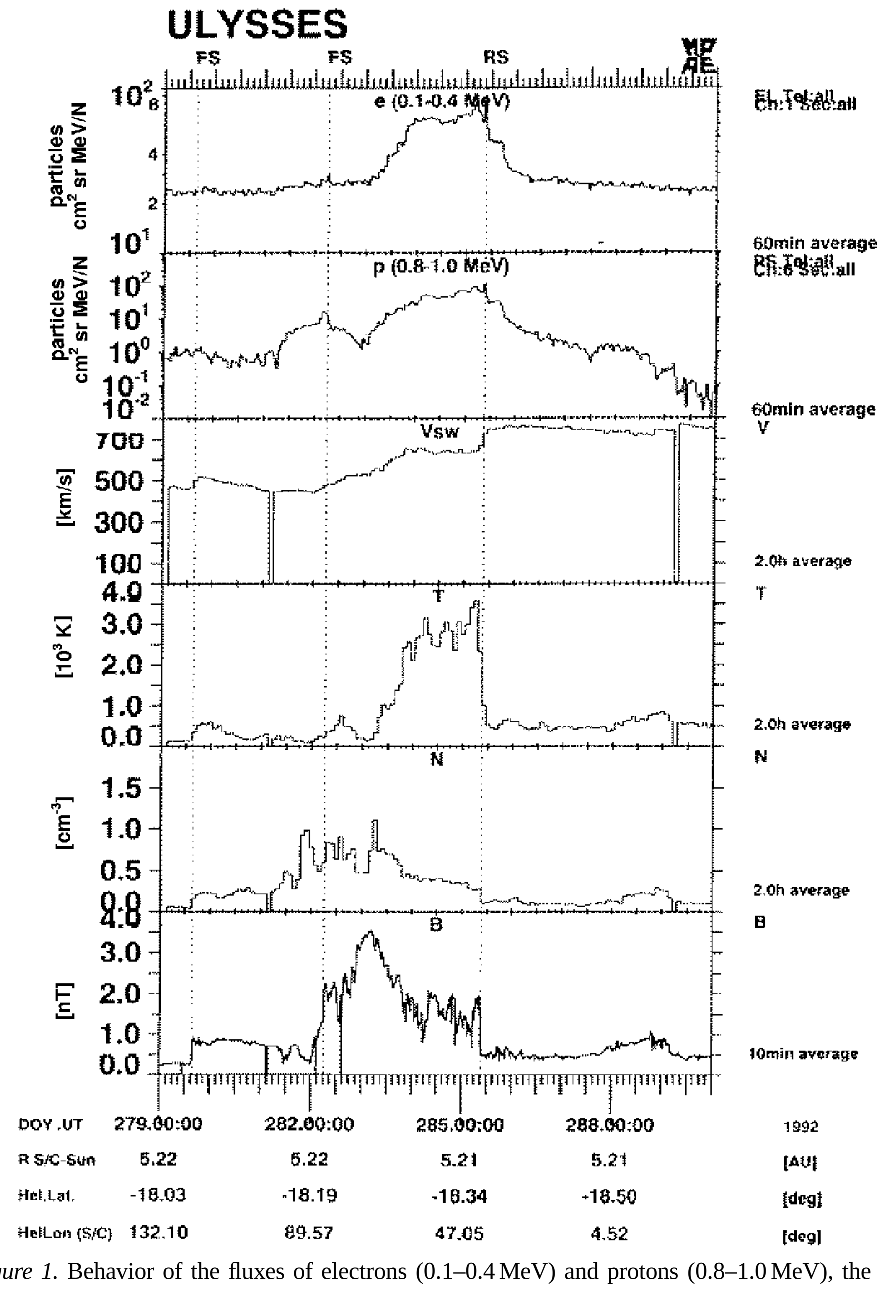

Figure 1. Behavior of the fluxes of electrons (0.1–0.4 MeV) and protons (0.8–1.0 MeV), the solar wind speed V_{SW}, the proton temperature T, the proton number density N and the magnitude of the magnetic field B during the crossing of CIR No. 5 by Ulysses.

shock is an order of magnitude larger than that near the forward shock, a feature which is well known from the Pioneer and Voyager measurements (Barnes and Simpson, 1976; Scholer *et al.*, 1980; Tsurutani *et al.*, 1982). In contrast to the two-peak ion structure the flux of energetic electrons is only enhanced at the reverse shock relative to the background by a factor of $\sim$3.

The spectral and compositional characteristics of CIR associated energetic ions have been summarized in the accompanying paper by Mason, von Steiger *et al.* (1999). In brief, the spectral form of CIR ions in the energy range above a few tens of keV is a power law with a steepening beyond $\sim$1 MeV and the composition is solar wind like. There are some differences, *e.g.*, the He/O and Ne/O ratios are observed to be a strong function of solar wind speed, while other ratios do not exhibit this variation. Strong heating and acceleration of pickup ions has been observed in association with CIRs (Gloeckler *et al.*, 1994). As stressed by Gloeckler (1999) in an accompanying paper, the intensity of pickup He^+ above twice the solar wind speed exceeds within CIRs at $\sim$5 AU that of suprathermal solar wind He^{++} even though solar wind He^{++} is at least a factor of 10^3 more abundant than pickup He^+. Furthermore, there is no difference in the spectral shapes of pickup He behind the forward and the reverse shock, while the solar wind heating efficiency is rather different. Gloeckler (1999) further points out that the inner source pickup ions with a C/O ratio of $\sim$1 could also contribute to CIR accelerated particles.

As can be seen from this brief overview and from the more detailed papers in this volume any successful theory of injection and acceleration of CIR particles has to explain a wide variety of phenomena. The present state of the theory is far from such a goal and what we have is merely a number of possible mechanisms. These mechanisms are not necessarily exclusive, but may well work at the same time for different species or at different locations. This chapter collects the mechanisms as favored by the different authors.

In the next section M. A. Lee and J. Kóta first introduce the general energetic particle transport equation and then evaluate the transport coefficients by quasi-linear theory. J.R. Jokipii then reviews the theory of diffusive or first order Fermi acceleration at oblique shocks. Subsequently M. A. Lee outlines in Sect. 4 a theory for diffusive acceleration at CIR shocks, which takes into account adiabatic deceleration in the expanding solar wind. This theory predicts the different spectral shapes at the forward and reverse CIR shocks, and has been successfully compared with spectral data in many cases. In Sect. 5 M. Scholer reviews work on particle injection at quasi-perpendicular shocks. Shock surfing may be an important injection and acceleration mechanism for pickup ions at the CIR shocks if the shock thickness is considerably less than the ion inertial length. The shock surfing theory is described by R. Kallenbach and M. A. Lee in Sect. 6. In this section there is also a discussion on theoretical constraints for a possible charge-per-mass (Q/A) dependence of the injection efficiency. There is observational evidence that stochastic acceleration is important for pickup ions well within the CIRs. L. A. Fisk reviews the theory of stochastic acceleration by Alfvén waves and by oblique propagating magnetosonic waves. He points out that transit time damping, *i.e.*, acceleration due to Landau resonance in magnetosonic waves is very efficient and can accelerate pickup ions well within the CIRs, where large fluctuations of the magnetic field strength have actually been observed. Only very little work has been done in the past on electron acceleration. In Sect. 8, G. Mann reviews what is known on electron acceleration and concentrates on shock drift acceleration of electrons. We should like to point out that in an accompanying paper Scholer (1999) presents a complementary review of injection and acceleration at CIR shocks. In the present chapter only the section on shock surfing addresses the problem of particle injection, and this process concerns basically the injection of pickup ions. Little theoretical/simulational work has been done on the problem of injection of solar wind ions at quasi-perpendicular shocks. A discussion of some of the ideas can be found in the article by Scholer (1999).

2. Energetic Particle Transport and the Diffusion Tensor

M. A. LEE and J. KÓTA

2.1. Introduction

The transport of energetic particles in the heliosphere is effectively described by the equation (Parker, 1965; Gleeson and Axford, 1967)

$$\frac{\partial f}{\partial t} + (\mathbf{V} + \mathbf{V}_D) \cdot \nabla f - \nabla \kappa \cdot \nabla f - \frac{1}{3} \nabla \cdot \mathbf{V}\, p \frac{\partial f}{\partial p} = Q \tag{1}$$

where $f(\mathbf{x}, p, t)$ is the particle omnidirectional distribution function, p is momentum magnitude, $\mathbf{V}(\mathbf{x}, t)$ is the solar wind velocity, and $Q(\mathbf{x}, p, t)$ is the source term (to describe for example ion injection at a shock). The magnetic field controls the spatial transport through the drift velocity, $\mathbf{V_D} = (pvc/3q)\nabla \times (\mathbf{B}/B^2)$, where v and q are the particle speed and charge, and $\mathbf{B}$ is the average magnetic field, and through the spatial diffusion tensor κ. Equation 1 is based on the assumption that $v \gg V$ and the spatial scalelength, L, is sufficiently larger than the scattering mean-free-path: $|\kappa| \ll vL/3$. The latter assumption is equivalent to the requirement that the particle distribution be nearly isotropic. If these conditions are not met then one should consider the more general Fokker-Planck equation (Skilling, 1971; Isenberg, 1997; Kóta and Jokipii, 1997) which remains valid for slow particle velocities and is equally applicable to either strong or weak scattering. Equation 1 may also be extended to include terms which describe stochastic acceleration (see Sect. 7), and viscous acceleration due to a shear in $\mathbf{V}$ (Earl *et al.*, 1988; Jokipii *et al.*, 1989).

The drift term describes transport due to curvature, gradient, and magnetization drifts in the "average" (usually viewed as an ensemble average) magnetic field. All other effects of the magnetic field are assumed to be diffusive and are lumped into the diffusion tensor κ. The diffusion tensor is often assumed to be axisymmetric about the unit vector $\mathbf{b} = \mathbf{B}/B$ (although it need not necessarily be), and therefore has the form

$$\kappa_{ij} = \kappa_{\perp}\delta_{ij} + (\kappa_{\parallel} - \kappa_{\perp})b_i b_j \tag{2}$$

where δ_{ij} is the Kronecker symbol. $\kappa_{\parallel}$ ($\kappa_{\perp}$) describes spatial diffusion parallel (perpendicular) to $\mathbf{b}$ due to magnetic field fluctuations. It is the efficiency of this spatial diffusive transport which insures near isotropy of the particle distribution. An additional antisymmetric component of the diffusion tensor, $\kappa_A \varepsilon_{ijk} b_k$, associated with the regular spiraling motion, is absorbed in the drift term in Eq. 1.

2.2. Quasilinear Derivation of $\kappa_{\parallel}$

The derivation of $\kappa_{\parallel}$ proceeds from the pitch-angle diffusion equation within a magnetic flux tube in the limit of strong scattering

$$v\mu \frac{\partial F}{\partial s} = \frac{\partial}{\partial \mu}\left[(1-\mu^2) D \frac{\partial F}{\partial \mu}\right] \tag{3}$$

where $F(s,t,\mu)$ is the particle phase-space distribution function, s is arclength along the flux tube, μ is the cosine of the particle pitch-angle, and D is the pitch-angle diffusion coefficient. Assuming that magnetic fluctuations vary only with s (slab model), the quasilinear theory yields (Jokipii, 1966, Jokipii, 1971b, Lee, 1971; 1982)

$$D = \frac{\pi q^2}{2m^2 c^2 |\mu| v} I\left(\frac{\omega}{\mu v}\right) \tag{4}$$

where m is particle mass, $\omega = qB/mc$ is the cyclotron frequency, and $I(k)$ is the wave intensity (or power) defined by

$$I(k) = \frac{1}{2\pi}\int_{-\infty}^{\infty} ds < \delta\mathbf{B}(s_0) \cdot \delta\mathbf{B}(s_0 + s) > e^{-iks} \tag{5}$$

(k is wavenumber). The cyclotron-resonance condition dictates that particles are scattered only if $kv\mu = \omega$. Equation 4 can be generalized to include wave fluctuations which also propagate oblique to the magnetic field.

With $F = f + g(\mu)$ and $|g| \ll f$ (nearly isotropic distribution), Equation 3 may be integrated to yield (Jokipii, 1966; Hasselmann and Wibberenz, 1970)

$$S = \frac{v}{2}\int_{-1}^{1} d\mu\, \mu g(\mu) = -\frac{v^2}{8}\int_{-1}^{1} d\mu \frac{1-\mu^2}{D}\frac{\partial f}{\partial s} = -\kappa_{\parallel}\frac{\partial f}{\partial s} \tag{6}$$

Thus, $\kappa_{\parallel}$ is a weighted integral over $D(\mu)^{-1}$. If $D(\mu)$ vanishes for a range of μ, $\kappa_{\parallel}$ diverges since particles cannot be scattered through this range of μ within the quasilinear theory. Higher-order corrections to $D(\mu)$ may be important for large-amplitude fluctuations and may lead to finite $\kappa_{\parallel}$ in this case.

2.3. THE PERPENDICULAR DIFFUSION COEFFICIENT $\kappa_\perp$

Stochastic transport normal to the average field is more complicated and controversial. The resonant scattering which causes the pitch-angle diffusion and parallel spatial diffusion also contributes to $\kappa_\perp$, since any scattering in pitch-angle causes the particle to shift the field line about which it gyrates. In addition, a given field line "random walks" about the average **B** causing perpendicular transport of a particle following the field line, the field line followed by the particle is ill-defined due to the finite gyroradius of the particle, and a particle drifts stochastically due to fluctuations on a scale larger than the gyroradius. All these effects, which contribute to perpendicular transport normal to **B**, are proportional to the intensity of the fluctuations (see *e.g.*, Forman *et al.*, 1974, Moussas *et al.*, 1982).

If fluctuations are small and gradients are not confined to the directions normal to **B**, then $\kappa_\perp \ll \kappa_\|$ and it is often appropriate to neglect $\kappa_\perp$. Alternatively, it is often assumed ad hoc that $\kappa_\perp = \eta\kappa_\|$, where η is constant and satisfies $\eta \ll 1$. For large amplitude fluctuations ($\delta B \sim B$) isotropic spatial diffusion is often appropriate ($\kappa_\perp \approx \kappa_\|$). A further idealized model is based on "hard sphere scattering" and yields

$$\kappa_\perp = \frac{v^2\tau}{3}\frac{1}{1+\omega^2\tau^2} \qquad \text{and} \qquad \kappa_A = \frac{v^2\tau}{3}\frac{\omega\tau}{1+\omega^2\tau^2} \tag{7}$$

where τ is the characteristic timescale for a large-angle (hard-sphere) scatter. Since $\kappa_\| = v^2\tau/3$, this formula gives the expected result $\kappa_\perp \approx \kappa_\|$ in the limit $\omega\tau \ll 1$. κ_A is the antisymmetric diffusion responsible for drift motion.

The contribution to $\kappa_\perp$ from the random walk of field lines is (Jokipii, 1971b)

$$\kappa_\perp = \frac{v}{2B^2} I(k=0) \tag{8}$$

where $I(k)$ is the magnetic field fluctuation power as given by Eq. 5 at zero wavenumber. Evaluation of the power at $k = 0$ follows from the fact that the long wavelength fluctuations dominate the random walk of field lines.

2.4. CHALLENGES

It is clear that Eqs. 7 and 8 for $\kappa_\perp$ are specialized and do not include all processes leading to stochastic transport across the average magnetic field. A more general approach follows from the Taylor-Green-Kubo formula (see Forman, 1977; Bieber and Matthaeus, 1997)

$$\kappa_{ij} = \int_0^\infty dt \langle v_j(t_0) v_i(t_0+t)\rangle \tag{9}$$

where the brackets $\langle\rangle$ describe an ensemble average over an appropriate ensemble of particle trajectories. Equation 9 also yields the drift transport (contained in the the antisymmetric terms) in the average magnetic field as described separately in

Eq. 1. In a sense Eq. 9 simply shifts the difficulty from evaluating κ_{ij} to evaluating the velocity correlation $\langle v_j(0)v_i(t)\rangle$. Bieber and Matthaeus (1997) postulate an exponential decay of $\langle v_j(0)v_i(t)\rangle$ with time to infer perpendicular diffusion and effective drift velocities, that are formally equivalent to Eq. 7 but τ includes the effects of both scattering and random walk of field lines.

If particles scatter back and forth in pitch angle, but remain strictly tied to field lines, then perpendicular diffusion results solely from the random walk and mixing of field lines. This idealized, but still physically valid, process is the so-called compound diffusion (Lingenfelter *et al.*, 1971) which is a non-Markovian motion, and which yields a slower diffusion than Brownian motion. In this case, the mean square displacement perpendicular to the mean field increases as $\langle \Delta x^2\rangle \propto t^{1/2}$, in contrast to the $\langle \Delta x^2\rangle \propto t$ dependence of standard diffusion. Compound diffusion may serve as a fair description when particle transport across the actual field lines is negligible. For the non-Markovian compound diffusion, $\langle v_j(0)v_i(t)\rangle$ has a long-time anticorrelation trend and may differ substantially from an exponential decay.

Observational investigation of an appropriate function for $\kappa_\perp$ has been limited, since $\kappa_\perp$ is usually dominated by $\kappa_\|$. However, Dwyer *et al.* (1997) have observed large anisotropy components normal to the average magnetic field for $\sim$100 keV/amu ions in the fast solar wind stream at 1 AU upstream of several large CIRs. The dependence of the anisotropy on the orientation of **B** allowed them to deduce that these normal anisotropy components are due to large perpendicular diffusion with $\kappa_\perp \sim \kappa_\|$ even though the scattering fluctuations are apparently not large amplitude. The reason for the large $\kappa_\perp$ is not known.

3. Theory of Shock Acceleration

J. R. JOKIPII

The mechanism for accelerating particles to 10^5–10^6 eV energies and higher, in the heliosphere, is generally thought to be diffusive shock acceleration. It is possible that other mechanisms have a role in the acceleration of the low-energy pickup ions and thermal ions to energies where shock acceleration takes over. However, shock acceleration may be responsible for accelerating these particles as well.

Diffusive shock acceleration has the virtue of naturally producing a power-law energy spectrum which is quite close to the spectrum observed. Diffusive shock acceleration is the natural consequence of the diffusive transport of fast charged particles at and near collisionless shock waves. Hence it can be derived readily from the general transport equation (1).

Associated with the solution to the pitch-angle-averaged distribution function f is a streaming flux $\mathbf{S}$, which may be written in terms of f as

$$\mathbf{S} = -\kappa\nabla f - \frac{1}{3}\mathbf{V}p\frac{\partial f}{\partial p}. \tag{10}$$

with an associated anisotropy magnitude

$$\delta = 3|\mathbf{S}|/(vf). \tag{11}$$

The anisotropy magnitude must be small compared with unity for the diffusion approximation to be valid.

Consider the solution of the above transport equation in a standard, planar shock configuration. We work in the shock frame, with the shock at $x = 0$, and take the magnetic field to be weak enough that the flow is unaffected by it. The $x - z$ plane is chosen to contain **B**. The upstream and downstream quantities are given the subscripts 1 and 2, respectively. We denote the strength of the shock by the ratio of upstream to downstream velocity $r_{\mathrm{sh}} = V_1/V_2 = \rho_2/\rho_1$, where ρ is the density. The normal magnetic field is unchanged at the shock and the transverse field increases by r_{sh}, $B_{z2}/B_{z1} = r_{\mathrm{sh}}$.

The solution to the transport equation may be obtained by solving it in the upstream regions and relating the solutions across the shock by the jump condition (obtained by integrating across the shock) (Jokipii, 1987)

$$\left[\kappa_{xx}\frac{\partial f}{\partial x} + \frac{V_x}{3}\frac{\partial f}{\partial \ln p} - \frac{pcv}{3q}\frac{B_z}{B^2}\frac{\partial f}{\partial y}\right]_1^2 = Q_*. \tag{12}$$

where Q_* is that part of the source which is concentrated at the shock, and where κ_{xx}, the coefficient of diffusion normal to the shock face, may be written, in terms of the angle θ between the magnetic field and the x direction, as $\kappa_{xx} = \kappa_{\|}\cos^2(\theta) + \kappa_{\perp}\sin^2(\theta)$.

The first two terms on the left yield the standard jump condition, and the third term gives the effects of the gradient and curvature drifts at the shock front. The third, drift term vanishes in planar, one-dimensional systems, since $\partial f/\partial y$ is zero in this case. However, even though the drifts at the shock do not appear explicitly in the mathematics for the diffusive shock acceleration at a planar shock, they play an extremely important role in understanding the physics of particle acceleration in all but purely parallel shocks. This is because it may be demonstrated (see, *e.g.*, Jokipii, 1987) that a significant fraction of the energy gain comes from drifting along the shock face, in the $\mathbf{V}\times\mathbf{B}$ electric field. Due to the fact that the mathematics is the same, the energy spectrum produced at a near-planar perpendicular shock including drifts is the same power law in momentum that is produced in the absence of drifts

$$f(p) \propto p^{-3r_{\mathrm{sh}}/(r_{\mathrm{sh}}-1)}. \tag{13}$$

The only difference in the spectrum produced including drifts from simple shock theory is that the particles may be accelerated considerably faster at oblique or quasi-perpendicular shocks, where $\kappa_{\perp}$ contributes more to κ_{xx}, as the acceleration time is

$$\tau_{\mathrm{acc}} \approx 4\frac{\kappa_{xx}}{V_1^2}, \tag{14}$$

and usually $\kappa_\perp \ll \kappa_\|$. This extra energy gain comes from the drift in the electric field, as mentioned above. Also, the injection of low-energy particles, such as thermal particles or pickup ions, is more difficult at quasi-perpendicular shocks than at quasi-parallel shocks. One may also show that the e-folding fall-off distance of the accelerated particles upstream of the shock is the same, $L = \kappa_{xx}/V_1$.

The association of energetic nuclei with CIRs has been recognized for more than 20 years (Barnes and Simpson, 1976; McDonald *et al.*, 1976). Since the shocks are generally highly oblique, drifts must play a role in the acceleration. Currently, it is thought that the particles are accelerated to several MeV energies at the forward and reverse shock waves bounding the CIR's. Less clear is the initial source of these particles. Early observations near 1 AU suggested that the composition resembled solar particles (*e.g.* Gloeckler *et al.*, 1979). Recent observations show that interstellar pickup ions are the likely source of many of these particles (Gloeckler *et al.*, 1994, and Gloeckler, 1999, in this volume).

4. A Model for Diffusive Shock Acceleration at CIRs

M. A. LEE

As pointed out in the previous section, steady state theory of diffusive shock acceleration at a planar shock predicts a power law distribution function, which is not observed at the CIR shocks for the full particle energy range. Fisk and Lee (1980) have solved the diffusion equation (1) by taking into account adiabatic deceleration (last term on the left hand side). Under the assumptions that the ion distribution is stationary in the frame corotating with the Sun and the spatial diffusion tensor is dominated by diffusion parallel to the Archimedes spiral magnetic field, the transport equation (Eq. 1) can be written

$$V\frac{\partial f}{\partial r} - \frac{1}{r^2}\frac{\partial}{\partial r}\left[r^2\kappa\frac{\partial f}{\partial r}\right] - \frac{2V}{3r}v\frac{\partial f}{\partial v} = 0 \tag{15}$$

where κ is here and subsequently the radial spatial diffusion coefficient, and drift transport has been neglected. Equation 15 describes the ion transport within an Archimedes spiral flux tube upstream of either the forward shock in the slow stream, or the reverse shock in the fast stream. The terms on the left hand side of Eq. 15 describe advection of the ions with the solar wind, diffusion within the flux tube, and adiabatic deceleration in the expanding solar wind. The shock acceleration is introduced with the boundary conditions at the shock ($r = r_s$) that (1) f be continuous and (2) at speeds above ion injection, the component of the ion streaming **S** normal to the shock front be continuous. Neglecting diffusive transport within the CIR, where the large-amplitude turbulence suppresses diffusive transport along the average field, these conditions combine to yield the boundary condition at $r = r_s$:

$$-\kappa\frac{\partial f}{\partial r} = \frac{1}{3}Vv\frac{\partial f}{\partial v}(1 - R^{-1}) \tag{16}$$

where R is the shock compression ratio. Fisk and Lee (1980) made the reasonable choice that $\kappa = \kappa_0 vr$, consistent with scattering mean free paths which increase with r and which are independent of rigidity (Palmer, 1982). With this choice Fisk and Lee derived the asymptotic solution to Eqs. 15 and 16 for large v as

$$f \sim (\frac{r}{r_\mathrm{s}})^{2R/(R-1)+V/(\kappa_0 V)} v^{-3R/(R-1)} \exp(-6\kappa_0 vR/[V(R-1)^2]) \tag{17}$$

The middle factor of Eq. 17 is the standard power law characteristic of shock acceleration at a stationary planar shock. The exponential factor describes the competing effect of adiabatic deceleration. Interestingly the factors $(r/r_\mathrm{s})^{V/(\kappa_0 V)} v^{-3R/(R-1)}$ give the expected distribution for low speeds (above the injection speed), even though the solution is asymptotic in v: the first factor describes the convective/diffusive ramp and the second factor is the standard power law. Thus Eq. 17 would appear to be more generally valid. Eq. 17 indeed satisfies the boundary condition of Eq. 16 exactly. If Eq. 17 is substituted into Eq. 15 one term remains uncanceled. This term is smaller than the other terms if $\ln(r_\mathrm{s}/r) < 3/2$ (for $\kappa_0 v < V$) or if $\ln(r_\mathrm{s}/r) < (3/2)\kappa_0^2(v/V)^2$ (for $\kappa_0 v > V$). Thus the spatial realm of validity of Eq. 17 increases with increasing v, as expected. At the lower speeds ($\kappa_0 v < V$) the solution requires at least $r_\mathrm{s}/r \leq 4.5$ with improving accuracy closer to the shock.

Equation 17 accounts for observed features of many corotating ion events: (1) the spectrum is exponential in v at higher energies with an e-folding speed independent of species; (2) at lower energies at the shock the spectrum is a power law; (3) the gradient, $f^{-1}\partial f/\partial r \propto r^{-1}$, is larger in the inner heliosphere; (4) if κ_0 and R are similar at the forward and reverse shocks, then the e-folding speed at higher energies is larger at the reverse shock since V is larger in the fast stream, in general agreement with higher ion intensities observed at the reverse shock. Furthermore, the general decrease in ion intensity within the CIR simply arises from adiabatic deceleration of the ions, which are trapped there by the large-amplitude turbulence. In addition, adjacent to the stream interface the ion intensity is expected to exhibit a dip since these field lines do not intersect the shocks (Palmer and Gosling, 1978).

Mason *et al.* (1997) present WIND observations of the large corotating ion event of DOY 340–343, 1994. At energies less than $\sim$1 MeV/amu in the fast stream they measure a power-law dependence of differential intensity ($\propto v^2 f$) on energy with an index of $\sim$2.2. According to Eq. 17 that implies a compression ratio $R \sim$1.9, which is reasonable for a strong reverse shock. For the same event at energies greater than $\sim$1 MeV/amu they measure an exponential spectrum with an e-folding speed of $\sim 8.5 \times 10^{-2}$ (MeV/amu)$^{1/2}$ = 1200 km/s. With an observed solar wind speed of $\sim$700 km/s, Equation 17 implies $\kappa_0 \sim 4.1 \times 10^{-2}$. With $\kappa_0 vr = \lambda_r v/3$, we obtain a radial scattering mean free path $\lambda_r/r \sim 0.12$, which again is reasonable. However, it should be noted that Mason *et al.* (1997) do not observe the reduced intensity at low energies expected from the first factor in Eq. 17; the low energy ions appear to be more mobile than expected. Reames *et al.* (1997) compare Eq. 17 with spectra observed by WIND during the event of May 30–June 9, 1995, and find very good agreement at three different times during the event. The adjusted parameters are very reasonable with the exception of the implied shock compression ratio late in the event, which is too large. Desai *et al.* (1999) fit the last two factors of Eq. 17 to the spectra measured by Ulysses during the hour centered on shock passage for all forward and reverse shocks encountered from Day 183, 1992 to Day 91, 1993. The inferred power-law spectral index was generally much smaller than the predicted value, $3R/(R–1)$; assuming an exponential in energy rather than speed appeared to provide a better fit and increased the inferred power-law spectral index to a value closer to the predicted value. These discrepancies appear to emphasize the importance upstream of the shock of the first factor in Eq. 17, which would harden the spectrum and reduce the inferred spectral index, and of a sheath of enhanced turbulence adjacent to the shock, which would modify the speed dependence of the exponential factor in Eq. 17.

The theory of Fisk and Lee (1980) makes assumptions which may have to be relaxed to obtain better agreement with observations. Their neglect of diffusive transport perpendicular to the average magnetic field is at odds with recent observations in the fast stream for several events of large diffusive fluxes normal to the magnetic field (Dwyer *et al.*, 1997). Since CIR shocks tend to be quasi-perpendicular, even small perpendicular diffusion coefficients can facilitate ion transport from the shock into the upstream solar wind. The assumption that $\kappa \propto r$ does not allow for a sheath of enhanced turbulence adjacent to the shock, which may be excited by the accelerated ions. Finally, the theory of diffusive shock acceleration cannot address the mechanism of ion injection at the shock. At quasi-perpendicular shocks the energy threshold for injection into the process of diffusive shock acceleration is quite large. Based on Ulysses observations at $\sim$5 AU, Gloeckler *et al.* (1994) find that interstellar pickup ions are preferentially accelerated by the CIR shocks. This may point to shock surfing as an important injection mechanism.

5. Numerical Models for Ion Injection at Shocks

M. SCHOLER

As has been outlined in the previous sections, the model of diffusive shock acceleration at CIR shocks in the expanding solar wind explains many observations associated with corotating energetic particle events. However, diffusive shock theory does not deal with the so-called injection problem, *i.e.* why and how a certain part of the thermal ions is extracted from the solar wind and becomes a suprathermal population, which is further accelerated by diffusive shock acceleration. Diffusive shock acceleration theory is thus not able to make statements about elemental abundances in corotating energetic particle events. In this and in the following section we will discuss analytical and numerical attempts which deal with the problem of injecting solar wind and pickup ions into a diffusive shock acceleration process.

Since CIR shocks tend to be quasi-perpendicular we have to deal at first sight only with the injection problem at quasi-perpendicular shocks. The only analytic attempt which is applicable to the problem of solar wind ion injection deals with quasi-parallel shocks: Malkov and Völk (1995) have extended the theory of diffusive particle acceleration to low energies, where the difference between the upstream and the downstream fluid frame is essential and the particle distribution is highly anisotropic at the shock front. In their model wave excitation and pitch angle scattering are treated self-consistently by assuming that pitch angle scattering is due to self-excited MHD waves propagating along the ambient magnetic field. These waves are excited in cyclotron resonance due to the pitch angle anisotropy of the backstreaming ions, *i.e.*, by an electromagnetic ion/ion beam instability. However, since it is assumed that the source of the injected particles are those downstream heated particles with an upstream velocity exceeding the shock velocity, the formalism does not really treat the injection problem, but rather treats acceleration in the thermal and suprathermal energy regime.

Shock surfing, which belongs to the category of shock drift acceleration processes at perpendicular shocks, has been modeled analytically by Lee *et al.* (1996) and will be discussed in detail in the next section to offer an explanation for the extremely preferential injection of pickup ions into CIR shock acceleration.

Numerical models for injection and acceleration at collisionless shocks fall into two categories:

- In Monte Carlo simulations it is assumed that thermal particles collide with scattering centers (magnetic irregularities) in the same way as the accelerated population. The process of acceleration during the sampling of the velocity difference across the shock can be modeled by the Monte Carlo technique (Ellison *et al.*, 1990). The scattering mean free path is usually assumed to be a power law function of momentum down to thermal energies. This approach does not really solve, but rather circumvents, the injection problem.
- In kinetic and hybrid simulations the interaction of particles with the shock is modeled self-consistently. Hybrid simulations of collisionless shocks treating the electrons as a neutralizing fluid not only successfully explain the shock micro-structure, but also show how diffuse suprathermal particles are directly injected out of the incident thermal plasma at quasi-parallel shocks (Quest, 1988; Giacalone *et al.*, 1992; Scholer, 1990).

5.1. MONTE CARLO SIMULATIONS

In the Monte Carlo simulations the mean free paths assumed for the thermal and suprathermal population are justified a posteriori from the mean free paths obtained by the hybrid simulations. These mean free paths are calculated from spatial intensity profiles of suprathermal particles upstream of the shock over one or more e-folding distances. Extrapolating the diffusion coefficient down to thermal energies in particle scattering across the shock does not necessarily describe correctly the process of thermal particle injection at quasi-parallel shocks. In particular, the assumption is made implicitly that the particles conserve their adiabatic moment during the shock interaction, which assumes that the magnetic field varies smoothly on spatial scales larger than the thermal gyroradius.

As mentioned, CIR shocks tend to be quasi-perpendicular shocks so that the Monte Carlo simulations and the one-dimensional and two-dimensional hybrid simulations are only of limited use. In the Monte Carlo model injection is basically due to that part of the downstream thermal distribution that has an upstream-directed velocity greater than the downstream bulk flow speed parallel to the magnetic field. Scattering is assumed to be isotropic in the de Hoffmann-Teller frame, where the upstream drift electric field is transformed to zero. In this model the efficiency of injection decreases rapidly with increasing shock obliquity and is effectively shut off when $\Theta_{Bn} > 30°$ (Baring *et al.*, 1994), since for large obliquity the thermal particles are rapidly swept downstream by the flow. Injection at quasi-perpendicular shocks can only be achieved by including cross-field diffusion. Baring *et al.* (1995) have extended the Monte Carlo technique by including cross-field diffusion: they found that cross-field diffusion effectively traps the thermal particles in the shock environs when the ratio of scattering mean free path to the particle gyroradius λ/r_g is of the order of one. Baring *et al.* (1995) were able to fit the proton spectrum for a quasi-perpendicular interplanetary traveling shock ($\Theta_{Bn} = 77°$) by assuming that $\lambda/r_g = 4$. Kinetic simulations should, in principle, include the cross-field diffusion process self-consistently. However, simulations in one or two dimensions have serious limitations. First, a reduction of dimensions implies a reduction in the allowable wave vector space. Second, Jokipii *et al.* (1993) have presented a general theorem according to which they show that charged particles in fields with at least one ignorable spatial coordinate are effectively forever tied to the same magnetic lines of force, except for motion along the ignorable coordinate. This theorem was derived by Jokipii *et al.* (1993) in a heuristic manner and has recently been derived rigorously by Jones *et al.* (1998). Since long-time 3-D simulations of shocks are not feasible at present, other approaches are necessary.

5.2. Hybrid Simulations

Giacalone *et al.* (1994) took refuge to a similar approach as did Baring *et al.* (1995). In 1-D hybrid simulations of perpendicular shocks they imposed an assumption on the ion motion so that diffusion across the magnetic field is possible. The cross-field scattering efficiency is expressed in terms of a scattering time τ measured in units of the inverse ion gyrofrequency, during which the gyrophase is randomised. Pickup ions were included self-consistently. Assuming a scattering time τ_{sc} of the order of $20\Omega^{-1}$ (Ω is the gyrofrequency) they obtained fast injection and acceleration of pickup protons, He^+, and heavier ions at perpendicular shocks. However, solar wind particles could not be injected into the acceleration process unless the rather unphysical value of $\Omega\tau = 1$ was chosen. As will be outlined in the next section, another injection and acceleration mechanism for pickup ions may be shock surfing.

As stated earlier, energetic ion ($>$50 keV) intensities at CIR reverse shocks typically are greater and broader in extent than at CIR forward shocks. Giacalone and Jokipii (1997) discussed two specific mechanisms for producing this effect. First, they pointed out that the model discussed by Fisk and Lee (1980) produces somewhat flatter spectra at the reverse shock, and this would, in general, produce higher intensities at the reverse shock, at least at the higher energies. Second, they introduced another possible effect which would produce an overall enhancement at the reverse shock at all energies, if the energetic particles were pickup ions: The pickup ions have a higher energy in the fast wind into which the reverse shock is propagating than in the slow wind into which the forward shock is propagating, because their energy (relative to the local plasma) is proportional to the square of the flow speed. They pointed out that this effect could be used to help establish the role of pickup ions in CIR-associated energetic particles.

As proposed by Scholer (1999) there might be another solution to the injection problem at CIR shocks without invoking extremely small mean free paths or small scattering times for thermal solar wind ions. Hybrid simulations show that solar wind ions are injected at quasi-parallel shocks. Since these ions are accelerated during their first shock encounter to energies considerably higher than that corresponding to specular reflection, backstreaming ions occur for shock angles $\Theta_{Bn} > 45^\circ$ (Scholer, 1998). Furthermore, the characteristic time required to produce backstreaming ions is short and only of the order of several gyroperiods. The interplanetary magnetic field (IMF) is fluctuating on all time and length scales. Large amplitude long wavelength fluctuations of the IMF may well lead locally to sporadic quasi-parallel situations, where ion injection can occur, although, on average, the CIR shocks are quasi-perpendicular. The observations at CIR shocks show that compared to pickup He^+ solar wind He^{2+} ions are accelerated less effectively although they are present in number densities that exceed those of pickup ions by orders of magnitude. The scenario with an occasional injection of ions during a more quasi-parallel situation would also explain the preferential injection of pickup He^+ ions relative to solar wind He^{2+}: Scholer and Kucharek (1999) have recently demonstrated that quasi-parallel shocks have a reflection efficiency for pickup ions exceeding that for solar wind ions by one to two orders of magnitude.

In concluding this section it should be noted that there exists no difficulty in injecting pickup ions into a shock acceleration process. As will be discussed in Sect. 7 pickup ions are easily accelerated by transit time damping in the downstream region. Transit time damping is acceleration in obliquely propagating fast magnetosonic waves by Landau ($n = 0$) resonance. Since the minimum energy in order to accelerate ions by the $n = 0$ resonance in magnetosonic waves is given by the Alfvén speed, thermal solar wind ions are not accelerated by transit time damping. Pickup ions are injected and accelerated at perpendicular shocks when a reasonable amount of cross field scattering occurs, whereas for injecting solar wind ions the scattering time has to be of the order of the inverse ion gyrofrequency. The kinetic simulations indicate that solar wind ions are easily injected into a diffusive acceleration process for more quasi-parallel shock configurations. This process also favors the injection of pickup ions by one to two orders of magnitude. Finally, shock drift acceleration also favors injection of pickup ions.

6. Shock Surfing and Shock Drift

R. KALLENBACH and M. A. LEE

A possible injection mechanism which strongly favors the injection of pickup over solar wind ions is "shock surfing" (Sagdeev, 1966; Lee *et al.*, 1996; Zank *et al.*, 1996; Zilbersher and Gedalin, 1997; Lipatov *et al.*, 1998). As illustrated in Fig. 2, a fraction of the gyrating pickup ions approach the shock under a peculiar angle with $|v_x| \ll u$, where u is the solar wind bulk velocity in the shock frame. These ions find themselves trapped between the electrostatic shock potential and the upstream Lorentz force. With each reflection at the shock potential they gyrate parallel to the motional electric field, picking up energy and surfing along the shock surface. Eventually their energy in the shock normal direction exceeds the shock potential, or the Lorentz force exceeds the electric field given by the gradient of the potential, and the ions gyrate downstream with a substantial energy gain. These ions can attain the threshold for diffusive shock acceleration at a perpendicular shock. In the simple case of a perpendicular shock with x in the direction of the upstream shock normal and $\mathbf{B} = B_z\mathbf{e}_z$, the ion equations of motion are

$$\frac{dv_x}{dt} = -\frac{q}{m}\frac{d\Phi}{dx} + v_y\omega, \quad \frac{dv_y}{dt} = -(u+v_x)\omega, \quad \frac{dv_z}{dt} = 0 \tag{18}$$

where $\omega = qB_z/m$, $-u$, $(u > 0)$ is the upstream flow speed relative to the shock, and Φ is the shock potential. On the right hand side of the equation for v_x are the two forces which can trap the ion and cause v_x to oscillate. Averaging the equation for v_y over this oscillation period yields $v_y = v_{y0} - u\omega t$. Taking a step function is in many cases a reasonable approximation for the potential of quasi-perpendicular shocks because their characteristic length L is usually much less than u/ω, the characteristic gyroradius of a pickup ion. For a discontinuous potential the ion is reflected (at time $t = 0$) in the upstream direction without change in v_{y0} and v_{z0} but changing the sign of v_{x0} if $mv_x^2/2 < q\Phi_0$ and $v_{y0} < 0$. The time between two bounces τ is then approximately given by $x = (v_{x0} + \omega v_{y0}\tau/2 - u\omega^2\tau^2/6)\tau = 0$; from this it also follows that there is a velocity change $\Delta v_x = -u\omega^2\tau^2/6$ compared to the case of no acceleration by the motional electric field. As $\omega\tau \approx -2v_{x0}/v_{y0}$ one can estimate (Lee *et al.*, 1996) that after the first bounce or two in the adiabatic limit, $\tau {v_y}^{-1} dv_y/dt \ll 1$, the relations

$$|v_x| = |v_{x0}|\left(\frac{|v_y|}{|v_{y0}|}\right)^{1/3}, \quad \varepsilon_x = \varepsilon_{x0}\left(\frac{\varepsilon_y}{\varepsilon_{y0}}\right)^{1/3}, \tag{19}$$

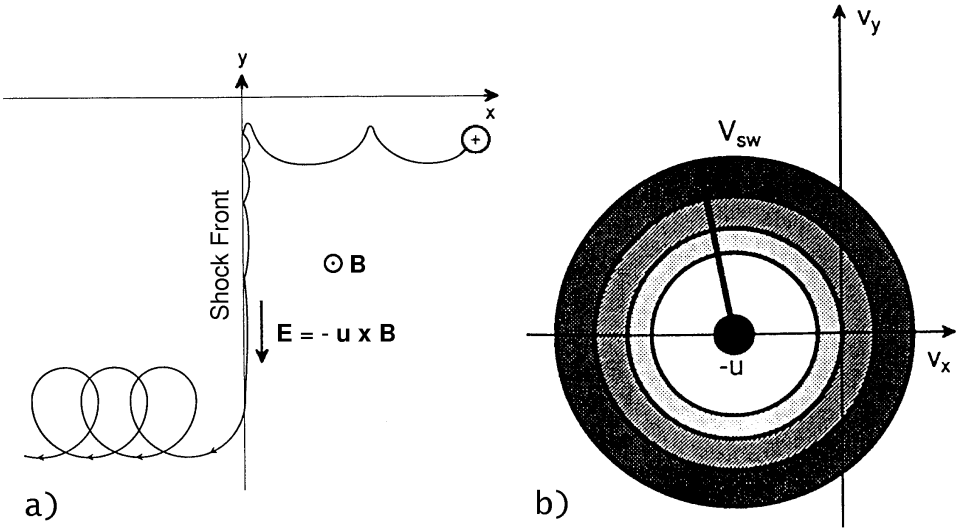

Figure 2. a) Schematic diagram of an ion interacting with a perpendicular shock at $x = 0$ when the ion first encounters the shock with $v_x \ll u$ and $v_y < 0$. **E** denotes the motional electric field, $\mathbf{E} = -uB_1\mathbf{e}_y$ ($\mathbf{u} = -u\mathbf{e}_x$). b) Pickup ion and solar wind velocity-space distributions in the plane ($v_z = 0$). The distributions are in the shock frame in which the solar wind (dark circle) is normally incident with a speed u. The pickup ion distributions are spherically symmetric in the solar wind frame with speeds of the order of the solar wind bulk velocity $V_{\rm SW}$ (adapted from Lee *et al.*, 1996).

are valid, where $\varepsilon_{x,y}$ are the particle energies in the x and y directions. When $mv_x^2/2 > q\Phi_0$ the ion is transmitted downstream, which can yield very large injection energies $mv^2/2 \gg mu^2/2$ because the ion gains much more energy in the y direction than in the x direction. For a continuous potential the electric field $E_x = \Phi_0/L$ across the shock is finite and in addition has to balance the Lorentz force on the order of qv_yB_z due to the high ion velocity along y, so that the ion is transmitted if $q\Phi_0 < \varepsilon_{x0}(\varepsilon_y/\varepsilon_{y0})^{1/3} + m\omega L|v_y|$ where the variation of B_z inside the shock layer has been neglected. Including a possible y component B_y of the magnetic field in the shock layer, the condition for transmission becomes $q\Phi_0 \leq \varepsilon_{x0}(\varepsilon_y/\varepsilon_{y0})^{1/3} + m\omega L|v_y| + qL|v_z||B_y|$. It has been shown by Lee *et al.* (1996) that for adiabatic particle motion ($\tau {v_y}^{-1} dv_y/dt \ll 1$) the relation between the energy gain in the x and y directions represented by the first right-hand term is conserved inside and outside the shock. In case of small ε_{x0} and small B_y the maximum injection energy is determined by $q\Phi_0/L \approx m\omega|v_y|$, which implies an injection energy of $mv^2/2 \approx mu^2/2\ [(\frac{m_p}{m})^2 \frac{q^2u^2}{e^2L^2\omega^2}]$. Here we used $\Phi_0 \approx \frac{m_p}{2e}u^2$ as the upper limit of the magnitude of the potential jump at the shock, which applies for high Mach numbers; in reality this jump is typically still about a few tenths of this value. If $L \ll u/\omega$ the mechanism yields large injection energies. The size of L and its implications on the efficiency of shock surfing remains to be discussed later in this section after the presentation of the case of non-adiabatic motion.

In the case of non-adiabatic motion inside the shock layer the acceleration of particles depends on the detailed field configuration there. Usually the equations of motion are non-linear and non-integrable and have to be treated numerically (Lembege *et al.*, 1983). However, here we present the example of a magnetosonic shock wave, where the solutions to the equations of motion can be approximated analytically. We refer to Tidman and Krall (1971) who described the potential step of a shock wave by the potential slope ahead of a solitary magnetosonic wave in a two-fluid plasma of protons and electrons. Electric and magnetic field **E** and **B** of the solitary wave are given by

$$\mathbf{E} = -B(x)V_{ye}(x)\mathbf{e}_x - uB_1\mathbf{e}_y,\ \mathbf{B} = B(x)\mathbf{e}_z,\ B(x) = B_1 + \frac{4\Delta B \exp\left(\frac{|x|}{L}\right)}{\left[1+\exp\left(\frac{|x|}{L}\right)\right]^2},$$

$$V_{ye}(x) = \frac{2u^2 - V_A^2(x) + V_{A,1}^2}{2u^2\mu_0 N_1 e}\frac{dB}{dx},\ V_A^2(x) = \frac{B^2(x)}{\mu_0 N(x) m_p},\ V_{A,1}^2 = \frac{B_1^2}{\mu_0 N_1 m_p}. \quad (20)$$

where B_1, N_1, and $V_{A,1}$ are the upstream magnetic field, plasma density, and Alfvén velocity, respectively, and $N(x) \approx N_1$; $B_2 = B_1 + \Delta B$ is the downstream magnetic field. The smaller the shock (soliton) width L, the larger is the electron fluid velocity $V_{ye}(x)$ and therefore the trapping field within the shock wave. The equation of motion for a test particle $(q,m) = (Qe,Am_p)$ with cyclotron frequency $\omega(x) = qB(x)/m$ including its energy gain can be described as:

$$\frac{dv_x}{dt} = \omega(x)v_y - \omega(x)V_{ye}(x),\quad \frac{dv_y}{dt} = -\omega(x)v_x - \omega_1 u,$$
$$\frac{1}{m}\frac{d\varepsilon}{dt} = -\omega_1 v_y u - \omega(x)V_{ye}(x)v_x. \quad (21)$$

In addition to the cyclotron terms, the equation for dv_x/dt contains a trapping term due to the shock potential, and the equation for dv_y/dt contains a term describing the acceleration in the motional electric field that determines the kinetic energy gain with time. The energy gain is positive for $-v_y > v_x V_{ye}(x)\omega(x)/(u\omega_1)$ which excludes practically all solar wind ions with typically $v_x \approx -u$ and $v_y \approx 0$ (note that $V_{ye}(x){<}0$) from injection but not the pickup ions with $|v_x| \ll u$ and negative v_y as initial conditions. The second term on the right-hand side can be integrated analytically. The first term corresponds to the energy gain due to the acceleration in the motional electric field and can be estimated for the usual case $u \ll |V_{ye}|$ so that $v_x^2 \ll v_y^2 \approx 2E/m \approx \omega_1^2 u^2 t^2$. However, this is only valid until $|v_y|$ approaches the maximum electron drift velocity $|V_{ye,m}|$ (see also Lembege *et al.*, 1983) at $x \approx 1.3L$. Then, the Lorentz force overcomes the trapping force due to the shock potential and bends the ion motion into the negative x direction with $|v_x|$ comparable to $|v_y|$ to leave the shock layer. This leads to a term $|V_{ye,m}|^2/2$ (Ohsawa, 1985) independent of Q/A, which turns out to be dominant in most cases for the energy gain of the particle transmitted through the shock:

$$\frac{\varepsilon}{m} \approx \frac{\alpha}{2} V_{A,1}^2 \frac{m_p}{m_e} \left(\frac{\Delta B}{B_1}\right)^{3/2} \left[1 - \beta \left(\frac{\Delta B}{B_1}\right) \frac{V_{A,1}^2}{u^2} - \gamma \left(\frac{\Delta B}{B_1}\right)^2 \frac{V_{A,1}^2}{u^2}\right]^2$$
$$- \frac{Q}{A} \frac{V_{A,1}^2}{2} \left[\left(1 + \frac{V_{A,1}^2}{2u^2}\right) \left(\frac{B_2^2}{B_1^2} - 1\right) - \frac{V_{A,1}^2}{16u^2} \left(\frac{B_2^4}{B_1^4} - 1\right)\right]. \quad (22)$$

The numerical constants α, β, and γ are typically on the order of 3, 0.7, and 0.2, respectively, depending on the individual ion trajectory. The electron drift velocity $|V_{ye,m}|$ and hence the energy gain of the particle crucially depends on the shock width L; it is approximately proportional to $1/L^2$. The theoretical value for L is close to the electron inertia length (Tidman and Krall, 1971), $L = c(B_1/\Delta B)^{1/4}/\omega_{pe}$. In that case, maximum particle energies of 100 keV/amu can be reached in CIR shocks, assuming an Alfvén velocity of about 50 km/s and jumps in the magnetic field by a factor $\sim$3. Data of the SOHO/CELIAS proton monitor at 1 AU suggest that the change in proton speed typically occurs on scales that are on the order of the proton inertia length or even longer (see *e.g.*, Ipavich *et al.*, 1998). However, the proton fluid is not expected to react within the electron inertia scale length. The proton scale length rather determines the decay of the "overshoot" visible in the magnetic field data presented in Fig. 4 of Livesey *et al.* (1982). The overshoot thickness of shocks observed by ISEE 1 and 2 is a few proton inertia lengths. The leading ramps of the shocks, in fact, have a thickness on the order of the electron inertia length, which matches theory quite closely. The many "wiggles" in the downstream turbulence also have similar short leading magnetic field ramps.

With increasing shock strength the terms containing Q/A become more important. At moderately strong shocks such as interplanetary shocks, particles with low Q/A are favored for injection and acceleration. In fact, particle abundances observed in the anomalous component of cosmic rays suggest that pickup ions with small Q/A are accelerated or transported more efficiently. In case of larger Alfvén velocities such as for shocks close to the Sun the accelerated particles may undergo even stronger Q/A-fractionation. This would match observations in solar energetic particles (Breneman and Stone, 1985): Elemental particle abundances can be ordered by Q/A where individual events favor large or small Q/A, possibly depending on the shock parameters according to Eq. 22. Observed abundances of the anomalous component of the cosmic rays also show Q/A as an ordering parameter that is related to the acceleration and/or the transport efficiency of pickup ions in the heliosphere.

After the shock surfing the particles are injected into the diffusive acceleration process or they are possibly further accelerated by shock drift, which is a process similar to shock surfing except that particle is transmitted across the shock twice during each gyration.

Shock surfing and shock drift describe in an idealized way part of the motion during the injection and acceleration process. As pointed out in Sect. 3 where a general theory of shock acceleration is given (see also Jokipii, 1987), a significant fraction of the energy gain comes from drifting along the shock face, in the motional electric field. However, for pre-acceleration and injection the particle interaction with magnetosonic waves is not only important at the main ramp of the shock wave, but in a more statistical form in the waves and turbulence of the solar wind downstream from the shocks as outlined in the next section.

7. Statistical Acceleration in the Solar Wind

L. A. FISK

The solar wind contains extensive waves and turbulence. The statistical interaction of energetic particles by these waves and turbulence is another mechanism for particle acceleration. Indeed, in regions of enhanced turbulence such as Corotating Interaction Regions (CIRs) it is reasonable to expect that some of the acceleration of energetic particles, particularly the relatively low-energy particles, must be statistical acceleration. It is possible also, given that turbulence is present throughout the solar wind, that there is a residual statistical acceleration always present, whose effects accumulate during the transit of the solar wind outward through the heliosphere and result in ever-present accelerated particles in the outer heliosphere.

Statistical acceleration is most readily described as a diffusion in momentum space. In other words, Equation 1 for the time variation of the isotropic particle distribution function, f, needs to include a term,

$$\left(\frac{df}{dt}\right)_{TT} = \frac{1}{p^2}\frac{\partial}{\partial p^2}\left(p^2 D_{pp}\frac{\partial f}{\partial p}\right) \tag{23}$$

to describe statistical acceleration. Here, D_{pp} is the rms change in momentum per unit time, averaged over particle direction. The change in momentum resulting from particles interacting with waves and turbulence is equally likely to be positive or negative; it is a true diffusion in momentum. However, since the magnitude of the particle momentum, or equivalently the energy, is only positive, the diffusive process moves the particles in the direction of increasing momentum or energy, and is a net acceleration.

To illustrate the change in the mean energy that results it is possible to describe the diffusion term in Eq. 23 in terms of the differential number density U, per unit interval of kinetic energy T, or as:

$$\left(\frac{df}{dt}\right)_{TT} = \frac{\partial}{\partial T}\left(D_{TT}\frac{\partial U}{\partial T}\right) - \frac{\partial}{\partial T}\left(\frac{D_{TT}U}{2T}\right) \tag{24}$$

Here, $D_{TT} = v^2 D_{pp}$, where v is particle speed. The statistical acceleration results in an rms change in energy, which is described by the coefficient D_{TT}, and a mean change in energy which is given by the term $D_{TT}/2T$.

There are two principal forms for propagating variations in the magnetic field in the solar wind which can give rise to statistical acceleration. (It is important to note that the waves must be propagating; static structures will not accelerate since an electric field is required.) Alfvénic fluctuations are the most common, and do not involve significant variations in the magnitude of the magnetic field strength. Magnetosonic waves are also present, in smaller amplitudes, and involve significant magnitude variations. There is, however, considerable difference in the acceleration capability of these two wave forms.

In the case of Alfvén waves the particles are pitch-angle scattered by the moving Alfvén wave. They gain or lose speed equal to the Alfvén speed, and, as has been shown by, *e.g.*, Jokipii (1971a) and Wibberenz and Beuermann (1972), the diffusion coefficient in energy space is given approximately by:

$$D_{TT} \approx V^2 T^2 / \kappa_{||} \tag{25}$$

where $\kappa_{||}$ is the spatial diffusion coefficient for pitch angle scattering parallel to the mean magnetic field direction. Thus, the rate of statistical acceleration depends inversely on the scattering mean free path. Said another way, pitch-angle scattering and statistical acceleration by Alfvén waves both depend on resonant scattering, where the interaction is at the first harmonic of the cyclotron frequency.

Indeed, it can be shown that the required mean free path for any interesting statistical acceleration from Alfvén waves is extremely small (Fisk, 1976a). For example, if the Alfvén speed is $\sim$50 km/s and the mean free path is $\sim$0.01 AU, then the characteristic acceleration time for 100 keV/amu particles is still $\sim$20 days. Mean free paths are generally considered to be much longer than this in the solar wind, although perhaps in very turbulent regions they could be this small. In any event, only for time scales characteristic of the transit of the solar wind over large distances is acceleration by Alfvén waves likely to be important, *i.e.* it is unlikely to be important in the acceleration of particles in CIRs in the inner heliosphere, but perhaps, although probably not likely, it could yield accelerated particles during the transit of the solar wind into the outer heliosphere.

In the case of magnetosonic waves the acceleration is much more efficient, and, indeed, this efficiency is the expected reason for the low amplitudes of such waves in the solar wind. Here, the acceleration occurs not due to a resonant interaction at the cyclotron frequency, but rather at the Landau resonance, where the phase speed of the wave parallel to the mean field direction is equal to the parallel speed of the particle; this mechanism is known as transit-time damping. In this case acceleration can occur in the absence of significant pitch angle scattering.

In Fisk (1976b) a detailed derivation is provided for the diffusion coefficient in momentum space, D_{pp}, for the transit-time damping of, or equivalently the statistical acceleration by, magnetosonic waves. The derivation is based on standard quasi-linear theory, which may not be too bad in this case, since the wave amplitudes are expected to be small. As expected, the diffusion coefficient depends on the amplitude of the magnitude of the field variations, on the phase speed of the waves, and it is an integral over the power of all the waves which satisfy the Landau resonance. In simple terms, the particles interact with moving magnetic gradients in the magnetic field, which, statistically, increase their parallel speeds.

The acceleration rates in Fisk (1976b) are considerable. That is, if there are magnitude variations in the magnetic field, there should be noticeable statistical acceleration. In the case of CIRs, such magnitude variations can be generated locally through the interaction of high and low speed flows, and the resulting shock waves. This local generation would be balanced, then, by the transit-time damping and subsequent acceleration of energetic particles. In the outer heliosphere, large-scale magnitude variations in the magnetic field are observed (Burlaga *et al.*, 1987) and could contribute to a general acceleration in the outer heliosphere. However, such variations may be more static than propagating.

The clearest evidence for statistical acceleration in the solar wind occurs for pick-up ions in CIRs. Schwadron *et al.* (1996) examined the spectrum of interstellar pickup ions in the solar wind as observed by Ulysses in 1992 at about 5.5 AU from the Sun, near the equatorial plane. During the period studied, some 22 forward and reverse shocks were observed, surrounding CIRs. The spectra of the pickup ions showed clear evidence of acceleration. Pickup ions are injected into the solar wind with a speed less than or equal to twice the solar wind speed, in the frame of the spacecraft; pickup ions observed in excess of this speed result from local acceleration. Clear evidence was seen for tails on the distribution function in excess of twice the solar wind speed. However, these tails clearly did not occur in coincidence with the shocks surrounding the CIRs, but rather were more likely, although not exclusively, to occur in coincidence with magnitude fluctuations in the magnetic field in the region between the shocks. Indeed, Schwadron *et al.* (1996) were able to fit quite well the observed spectra of the pickup ions by assuming the transit-time damping rates of Fisk (1976b).

8. Electron Acceleration

G. MANN

So far, we have only discussed injection and acceleration of ions at CIR shocks. CIR related shock waves are not only able to produce energetic ions but also energetic electrons, as can be seen from Fig. 1. While during this event the flux of energetic protons increases by a few orders of magnitude at both the forward and reverse shock, the flux of energetic electrons is predominantly enhanced only at the reverse shock. The evolution of energetic electrons and ions has been discussed by Simnett and Roelof (1995) and Roelof *et al.* (1996) for the period of the southbound journey of Ulysses. At low heliospheric latitudes both energetic electron and ion flux increases were seen with the appearance of the forward (FS) and reverse shocks (RS), with the larger flux increases usually observed at the reverse shocks (Keppler *et al.*, 1996). The last FS-RS pair was observed during CIR No. 15 at a distance of 4.58 AU and a latitude of -33.7°, although the RSs were observed by Ulysses for several more solar rotations (Gosling *et al.*, 1993a). After that, the peaks of energetic electron fluxes were delayed with respect to those of ions by a few days (Simnett and Roelof, 1995). Simnett and Roelof (1995) suggested the following explanation for this delay: Electrons accelerated at CIR shocks have high speeds (for 0.4 MeV $v = 0.8c$; c is the speed of light) and move essentially scatter-free. They are thus able to propagate upstream into the inner heliosphere. Since the IMF increases towards the Sun, these electrons are reflected due to the conservation of their magnetic moment and return to the RS for repeated acceleration.

Thus, a population of energetic electrons can be established due to the inter-relation of a global mirroring in the inner heliosphere and a local acceleration at CIR related shocks. After the spacecraft encounters the local corotating shock it will be magnetically connected to the shock beyond 5 AU. Since the RS is stronger at larger radial distances out to about 10 AU (Pizzo, 1994), it is able to accelerate electrons more efficiently. Thus, higher fluxes of energetic electrons are expected when the spacecraft is further upstream on field lines connected to the stronger shock. On the other hand, energetic ions have a velocity of $0.046c$ (for 1 MeV). Therefore they cannot penetrate deep into the inner heliosphere like the electrons. Furthermore their mean free path is smaller and they do not travel scatter free. Their maximum flux is expected close to the shock and falls off rapidly with distance along the field line in the upstream direction (Simnett and Roelof, 1995; Roelof *et al.*, 1996).

The efficiency of electron acceleration at CIR-related shocks is investigated by comparing the electron fluxes at the shock crossing with the shock parameters presented in Mason, von Steiger *et al.* (1999). Here, the electron fluxes are used as measured by the HISCALE instrument (Lanzerotti *et al.*, 1992) in the range 30–50 keV. The electron fluxes j_{shock} at the shock crossing are compared with the unperturbed fluxes j_0 determined during quiet periods before and after the CIR. The values for $\log(j_{\mathrm{shock}}/j_0)$ vary between 0.207 and 3.259 for 32 shock crossings at the CIRs 1–18. The reverse shocks of CIRs 4–7, 9, and 13 as well as the forward shock of CIR 6 accelerate electrons very efficiently, *i. e.* $j_{\mathrm{shock}}/j_0 > 54$. In 18 of these 32 shocks the flux ratios $\log(j_{\mathrm{shock}}/j_0)$ could be related to the jumps of the magnetic field B_2/B_1 and the Alfvén-Mach numbers M_{A} of the associated shocks as depicted in Fig. 3. Thus, the efficiency of electron acceleration expressed by $\log(j_{\mathrm{shock}}/j_0)$ increases for shocks with a stronger magnetic field jump B_2/B_1 and a higher Alfvén-Mach number M_{A} as can be seen from Fig. 3.

Since CIR-related shocks become stronger at larger heliospheric distances (*e.g.*, Pizzo, 1994) the intensity of energetic electrons produced by CIR-related shocks is larger at several AU than at 1 AU as indeed observed. Mason, von Steiger *et al.* (1999) reported that electrons in the energy range 50 to a few 100 keV are mainly produced at CIR related shocks at several AU, while only small intensities of energetic electrons have been measured at 1 AU. Note that there are also fewer shocks at 1 AU.

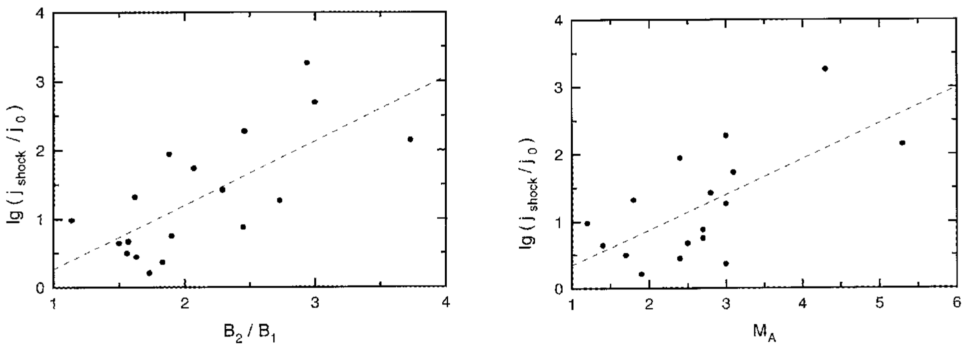

Figure 3. Correlation between the ratio $\log(j_{\rm shock}/j_0)$ and B_2/B_1 of the magnetic field (left) and the Alfvèn-Mach number $M_{\rm A}$ (right) for 18 CIR related shock crossings. The dashed lines represent the linear regression.

Shock waves can accelerate electrons by (1) shock drift acceleration, (2) a first-order Fermi process, or (3) by stochastic acceleration. Stochastic acceleration of electrons has been considered as an acceleration mechanism in the corona in connection with flares (Miller, 1997). However, in interplanetary space energetic electrons are almost scatter-free, so that stochastic acceleration here is unlikely. As outlined in Sect. 3, in the diffusive limit first-order Fermi acceleration incorporates shock drift acceleration at oblique shocks. However, since electrons behave essentially scatter-free the diffusion-convection equation (1) cannot be immediately applied to electron acceleration at CIR shocks. Furthermore, electrons behave quite differently in comparison to ions because of their smaller gyroradius, and they respond to rather different frequency regimes of the shock-induced plasma waves because of their opposite gyromotion and higher velocities. In the following, we will discuss shock drift acceleration of electrons in some detail.

A fast magnetosonic shock represents a moving magnetic mirror at which charged particles can be reflected and accelerated. Following the nonrelativistic approach given by Holman and Pesses (1983) and Schwartz *et al.* (1983) the calculations are conveniently done in the de Hoffmann-Teller frame, in which the shock is at rest and the motional upstream electric field is removed. In this frame the particles are reflected by conserving their magnetic moment and kinetic energy. After the transformation into the inertial frame the velocity component $V_{r,||}$ parallel to the upstream magnetic field after the shock encounter is related to the initial one $V_{i,||}$ via

$$V_{r,||} = 2v_{\rm s} \sec\theta_{Bn} - V_{i,||} \tag{26}$$

Here, θ_{Bn} is the angle between the shock normal and the upstream magnetic field and $v_{\rm s}$ denotes the shock speed. This mechanism is called shock drift acceleration. Assuming a Maxwellian distribution for the electrons with a thermal velocity $v_{\rm th} = (k_B T/m_e)^{1/2}$ ($k_{\rm B}$, Boltzmann's constant; T, temperature; $m_{\rm e}$, electron mass) in the upstream region, a *shifted loss-cone* distribution

$$f(V_{r,\|},V_{r,\perp}) \;\sim\; \Theta(V_{r,\|}-V_{\rm s})\Theta(V_{r,\perp}-[V_{r,\|}-V_{\rm s}]\tan\alpha_{\rm lc}) \times \exp\{\frac{-[(-V_{r,\|}+2V_{\rm s})^2+V_{r,\perp}^2]}{2v_{\rm th}^2}\} \tag{27}$$

(where $V_{\rm s} = v_{\rm s}\sec\theta_{Bn}$) results for the reflected particles (Leroy and Mangeney, 1984; Wu, 1984). Here, Θ denotes the step function and $V_{r,\perp}$ is the velocity component perpendicular to the upstream magnetic field for the reflected electrons. The loss-cone angle $\alpha_{\rm lc}$ is given by the jump of the magnetic field B_2/B_1 across the shock, $\alpha_{\rm lc} = \arcsin[(B_2/B_1)^{-1/2}]$. The ratio between the velocity gain (Eq. 26) and the thermal velocity is given by $\Delta V/v_{\rm th} = (2m_{\rm e}/m_{\rm p})^{1/2}\beta^{-1/2}2M_{\rm A}\sec(\theta_{Bn})$ (β, plasma beta; $m_{\rm p}$, proton mass). Inspecting Eqs. 26 and 27 the energy gain and the number of accelerated electrons increase with increasing Alfvén-Mach number $M_{\rm A}$ and magnetic field jump B_2/B_1 (or decreasing $\alpha_{\rm lc}$), respectively. Both properties are seen in the observations as demonstrated in Fig. 3. Adopting $M_{\rm A} = 3.0$, $\theta_{Bn} = 70°$, and $\beta = 1$ as typical parameters of reverse shocks (*cf.* Table I of Mason, von Steiger *et al.*, 1999) the velocity gain $\Delta V/v_{\rm th} \approx 0.6$ is small for a single shock encounter. But the mirroring of the accelerated electrons in the inner heliosphere as proposed by Simnett and Roelof (1995) leads to multiple encounters with the shock and in turn to higher electron energies. It is an open question how solar wind electrons are injected into such a shock drift acceleration process. We note in this respect that backstreaming suprathermal solar wind electrons are commonly observed upstream of corotating forward and reverse shocks beyond $\sim$2 AU (Gosling *et al.*, 1993b). The shocks are evidently able to accelerate electrons directly out of the solar wind.

9. Summary

9.1. Acceleration

Ions. There is general agreement that diffusive or first-order Fermi acceleration at the corotating shocks is responsible for many corotating energetic ion events. Diffusive shock acceleration explains the intensity peaks at well developed forward and reverse CIR shocks between about 3–5 AU. Diffusive shock acceleration theory, including adiabatic deceleration in the expanding solar wind, explains reasonably well the spectral shapes of many events, as well as the differences in intensities between forward and reverse shocks.

Energetic ion increases have also been observed at trailing edges of compression regions even when no reverse shock was detected. Currently, it cannot be decided whether acceleration in these cases occurs at the trailing boundaries of these compression regions in a similar fashion as at shocks, or whether acceleration occurs at localized shocks and the particles then undergo cross-field diffusion onto field lines not connected with the shock. The theory of shock acceleration is rather well developed, the missing ingredients as far as CIR shock acceleration is concerned are details of the transport coefficients, *i.e.*, the radial dependence, the dependence on energy, and on mass/charge of the spatial diffusion tensor. These transport coefficients are important for predicting spectral shapes and abundance variations away from the source regions, *i.e.*, at different heliographic longitudes or in the inner solar system. Another open question is the importance of the large-scale structure and spatial extent of the CIR shocks on the acceleration process.

Electrons. Electron acceleration at CIR shocks has not received much theoretical attention. Shock drift acceleration plays a more important role than for ions. One important question is whether electrons can get accelerated out of the solar wind or whether an energetic background population is further accelerated by shock drift acceleration. There is experimental evidence that electrons can be accelerated to suprathermal energies at CIR shocks (Gosling *et al.*, 1993b). Since the electron mean free path is large, such a drift cannot be described in terms of a diffusive acceleration mechanism. Electrons can more easily probe large distances along the magnetic field; the large-scale IMF structure is therefore important for an understanding of the electron time-intensity profiles.

9.2. Origin and Injection

The abundances of different energetic particle species is rather given by pre-acceleration and injection mechanisms which discriminate particles by their velocity distribution and therefore also by their origin. A number of possibilities have been discussed in the literature for the injection and/or acceleration of pickup ions at or downstream of CIR shocks. These mechanisms or scenarios either favor pickup ions relative to thermal solar wind ions or do not work for thermal ions at all.

Pickup Ions. Strong heating and acceleration of pickup ions has been observed in association with CIRs. Above twice the solar wind speed the intensity of pickup He^+ exceeds within CIRs that of solar wind He^{++}. Pickup ions have speeds between zero and twice the solar wind speed and thus part of the pickup ion distribution constitutes a suprathermal particle population. This eases the injection problem for pickup ions as compared to thermal solar wind ions. For efficient injection and diffusive acceleration at a perpendicular shock it is sufficient that there is cross-field scattering with a scattering time of about 20 times the inverse proton gyrofrequency. Pickup ions are then injected and accelerated at a perpendicular shock, but solar wind ions are not. However, the observational evidence in data from Ulysses at $\sim$5 AU that energetic pickup ions are enhanced in their abundance over the solar wind ions by several orders of magnitude is notable and requests a further quantitative theoretical investigation. Some analytical models have been proposed for preferential injection and acceleration of pickup ions:

– Shock surfing works for pickup ions but not for thermal solar wind ions. This is simply due to the fact that part of the pickup ion distribution has almost zero velocity with respect to the shock. However, shock surfing is only an efficient acceleration mechanism up to energies of $\sim$100 keV/amu if the length scale of the leading ramp of the cross shock potential is on the order of the electron inertial length; there is observational evidence for such strong field gradients in magnetosonic shock waves at Earth's bow shock. At larger shock length scales, part of the pickup ions are, like solar wind ions, just specularly reflected at a quasi-perpendicular shock and are convected downstream.
– Transit time damping is another mechanism which favors pickup ions relative to solar wind ions. It is a statistical way of ion acceleration by obliquely propagating magnetosonic waves. It has been proposed as an acceleration mechanism for pickup ions in the region downstream of the CIR shocks, and there is observational evidence for a correlation between suprathermal pickup ion fluxes and fluctuating magnetic field strengths in CIRs. The accelerated pickup ions can travel upstream and could then participate in a diffusive shock acceleration mechanism which boosts them to higher energies.

Solar Wind Ions. One of the main open questions in models is how solar wind ions are injected into a diffusive acceleration process at CIR shocks; there exists so far no reasonable model for the injection of thermal ions into a first-order Fermi acceleration process at quasi-perpendicular shocks although there has already been experimental evidence that this happens at the bow shock and at shocks driven by coronal mass ejections at 1 AU:

– One proposal made is that since the IMF fluctuates on all time and length scales the IMF direction is variable and, although CIR shocks are on average quasi-perpendicular, they are sufficiently often quasi-parallel. Injection of thermal ions at quasi-parallel shocks occurs on time scales of a few ion gyroperiods; in addition, injection of pickup ions at quasi-parallel shocks is by orders of magnitude more efficient so that this scenario would also favor pickup ions.
– Solar wind ions could participate in a Fermi type process at the shock if they scatter very efficiently, almost at the Bohm diffusion rate, across field lines. This seems rather unlikely, since with such a small scattering time (of the order of the inverse gyrofrequency) the whole magnetic field profile across the shock would change considerably. Such drastic changes have at least not been seen at Earth's bow shock. One of the outstanding questions is a self-consistent determination of the cross-field diffusion coefficient down to thermal energies.

Finally, in Table I we summarize the models discussed in this working group chapter and the related key observational features which are reported in the accompanying papers by Mason, von Steiger *et al.* (1999), Gloeckler (1999), and Mason and Sanderson (1999) in this volume. For a discussion of models on the energetic particle transport from CIRs to high latitudes we refer to the introductory article by Fisk and Jokipii (1999) and the working group report by Kunow, Lee *et al.* (1999).

TABLE I

Summary of models on ion injection and acceleration in CIRs and related observations

Process	**Related Observations**
Diffusive shock acceleration	- Intensity peaks of accelerated particles at well developed forward and reverse CIR shocks between about 3–5 AU. - Power law energy spectra in the range ∼10–1000 keV/amu. - Same time-intensity profiles from He through Fe (no strong Q/A dependence above injection threshold). - Different energetic ion intensities at forward/reverse shocks.
Adiabatic deceleration	- Steepening of energy spectra above ∼1 MeV/amu.
Diffusive transport	- At 1 AU generally sunward flow away from the location of peak intensities at ∼3–5 AU). - No change in spectral forms out to tens of AU (energetic particle transport from ∼3–5 AU). - Poor inward transport of He^+ at tens of keV/amu between 1 and 5 AU. - Strong inhibition of particle transport across stream interface (small cross-field diffusion).
Injection processes sensitive to the seed particle velocity distributions (transit time damping and/or shock surfing)	- Suprathermal tails in the velocity distributions of different species have identical spectral shapes (pre-acceleration mechanisms are velocity dependent). - Strong enhancement of pickup over solar wind ions. - He/O and Ne/O ratios increase with solar wind speed. - C abundance enhanced by a factor of 2–3 over O (predominantly inner source pickup C). - Non-shock pre-acceleration takes place in the turbulent inecliptic solar wind; tens of keV ions at 1 AU even in absence of shocks. - More pre-accelerated He^+ than He^{++} is observed downstream of shocks where turbulence is stronger. - Spectral shapes of tails are complicated (not simple power laws). - CIR pre-acceleration is limited to about 10 keV/amu.
Shock drift acceleration of electrons	- 50 keV to 100 keV electrons at several AU; small intensities at 1 AU.

Acknowledgements

Two of the co-authors (SC and RK) contributed to this article in the frame of the INTAS cooperative project “The Heliosphere in the Local Interstellar Cloud”.

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