Relativistic and Non-Relativistic Shock Acceleration in Various Objects

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Abstract

The acceleration of energetic charged particles at relativistic and non-relativistic shock waves plays an important role in powering esztablished and potential cosmic TeV-gamma-ray sources as supernova remnants, jets of active galactic nuclei and gamma-ray bursts. We therefore review some recent results on non-relativistic shock acceleretion and describe the current standing of the theory for relativistic shock acceleration. From the experience with non-relativistic shocks we emphasize that the microphysics of the interaction of plasma wave turbulence with the shock wave is important not only for the resulting momentum spectrum of accelerated particles but also for the shock wave structure itself.

1. Introduction

Three types of explosive astrophysical events are of interest to air Cherenkov telescopes: supernova remnants (SNRs), gamma-ray bursts (GRBs) and jets of active galactic nuclei (AGNi). Supernova remnants result from the massive explosions of stars at the end of their life, and are usually modelled as spherical outflows with large but non-relativistic outflow velocity $V_0 << c$. The leading collapsar (Woosley 1993, Paczynski 1998) and supranova (Vietri and Stella 1998) models for gamma-ray bursts involve a relativistic outflow that emanates from a compact source, where due to strong stellar envelope magnetic fields and/or explosions in a pre-existing pulsar wind bubble anisotropic relativistic outflow velocities result with initial Lorentz factors $\Gamma_0 = (1 - (V_0/c)^2)^{-1/2} \simeq 300$. AGN jets can be regarded as channelled collimated relativistic outflow components ($\Gamma_0 \simeq 30$) propagating in the dynamic jet medium which might be identical to the ambient interstellar medium (e. g. Pohl and Schlickeiser 2000). In all three sources it is of interest to unravel the basic physical processes that convert the huge kinetic outflow energy into observable nonthermal radiation.

According to current understanding there are two relevant dissipation processes of outflows in their interaction with the ambient cosmic plasmas:

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(1) two-stream plasma instabilities of the dense outflow plasma in the dilute interstellar medium,
(2) the formation of magnetized relativistic and non-relativistic shock waves in the dilute interstellar medium environment.

Because the surrounding interstellar medium from a plasma physics point of view is extremely collision-pure, a magnetohydrodynamic description of the interaction and the forming shock waves is not appropriate at all spatial scales of interest, so that both dissipation processes have to be investigated using a full plasma kinetic theory.

Both dissipation processes lead to efficient acceleration of energetic charged particles that then produce the observable nonthermal radiation. As a consequence of the two-stream plasma instability, electrostatic and magnetohydrodynamic plasma turbulence is generated that then accelerates charged particles by resonant wave-particle interactions. If indeed magnetized shock waves form, energetic particles can be accelerated by the diffusive shock wave acceleration mechanism. In the following I review some recent results on non-relativistic shock acceleretion and describe the theoretical difficulties occuring for relativistic shock acceleration.

2. Standard model of particle acceleration in outflow sources

In most existing radiation models of GRBs and AGN jets, it is assumed that adiabatic shock waves form from the interaction of the outflow with the surrounding interstellar medium, whose properties are determined by the magnetohydrodynamic shock relations, including sometimes nonlinear effects due to the back reaction of the accelerated particles. Energetic charged particles are efficiently accelerated by the diffusive shock wave acceleration mechanism, providing power law distribution functions in particle rigidity for relativistic hadrons and pairs. By simple equipartition arguments, it is assumed that a considerable fraction of the outflow energy is transformed into these power law distributions of energetic particles. Time-dependent modelling of the nonthermal synchrotron, synchrotron-self-Compton and external Compton cooling of energetic pairs in the evolving outflow source region are then performed to explain multiwavelength spectra and light curves in the optically thin case. Analogous calculations in an optically thick environment very often start from energetic hadrons that via cascades involving inelastic hadron-hadron-interactions and/or photomeson interactions provide copious amounts of secondary pairs.

Simple leptonic cooling models are remarkably successful in explaining the observed multiwavelength spectra of AGN jets (see Dermer and Schlickeiser 2002 and references therein) and GRB afterglows (e.g. Dermer, Böttcher and Chiang 2000); however, the measured rapid variability in the TeV light curves from Mrk 501 and Mrk 421 remains a challenge for these types of models. Especially, the long-time TeV monitoring of Mrk 501 (Gaidos et al. 1996; Quinn et al. 1999; Aharonian et al. 1999, 2001) indicates that during active phases (of order several months) many individual short time (of order days) intense TeV-flares occur. Within simple one-zone pair cooling models this implies that for every indivudual flare a new emission component has to be ejected from the central source that cools down with associated high-frequency TeV emission first, and delayed emission at sub-TeV frequencies. Within the shock acceleration and the hadron cascade scenarios it is unclear which processes provide the rapid short-term variability.

Another drawback of the standard model is that the diffusive acceleration mechanism even in the test particle limit is not fully understood both at non-relativistic and at relativistic shock speeds. For relativistic shocks the inferred universal particle's power law spectral index value $s = 2.23 \pm 0.01$ (Heavens and Drury 1988, Bednarz and Ostrowski 1998, Kirk et al. 2000) differs from what is needed for radiation modelling. Open issues for non-relativistic shock acceleration are the question of generating flat ($s < 2$) particle spectral indices and the resulting electron-proton ratio. As we shall argue, these discrepancies may have their origin in the non-kinetic treatment of the magnetized shock wave properties in these cosmic collision-pure environments. In any case, it is appropriate to inspect the microphysics of shock wave acceleration more closely.

3. Test-particle acceleration in relativistic and non-relativistic flows

Existing analytical work on the acceleration of test particles in relativistic and non-relativistic flows is based on two particle transport equations: the Fokker-Planck equation (1) in case of relativistic flows, and the diffusion-convection equation (3) in case of non-relativistic flows. Both have their origin in the quasilinear description of plasma kinetics. Due to the high conductivity of most cosmic plasmas large-scale steady electric fields are absent, so that the interest concentrates on magnetized plasma. The quasilinear approach to the interaction of charged particles with partially random electromagnetic fields ($\vec{B}_0 + \delta\vec{B},\ \delta\vec{E}$) is a first-order perturbation calculation in the ratio $q_L = (\delta B/B_0)^2$ and requires smallness of this ratio with respect to unity. In most cosmic plasmas this requirement is well satisfied as has been established by direct in-situ measurements in interplanetary plasmas, or due to saturation effects in the growth of fluctuating fields. The standard quasilinear approach also requires incoherent mode coupling of the fluctuating electromagnetic fields described as the superposition of individual plasma wave modes.

The stationary Fokker–Planck equation for the gyrotropic particle phase space density $f(z,p,\mu)$, where z denotes the spatial variable along the ordered uniform magnetic field $\vec{B_0} = B_0 e_z$ and $\mu = p_\parallel/p$, is formulated in the mixed comoving frame (i.e. p,μ are measured in the fluid frame while z is measured in the lab frame). The Fokker–Planck equation reads (Kirk et al. 1988)

$$\Gamma(U+v\mu)\Big[\frac{\partial f}{\partial z} - \frac{\partial U}{\partial z}\Gamma^2(m^2c^4+p^2c^2)^{1/2}(\mu\frac{\partial f}{\partial p} + \frac{1-\mu^2}{p}\frac{\partial f}{\partial \mu})\Big] = S(z,p,\nu)+\mathcal{E}(f) \quad (1)$$

where $S(z,p,\mu)$ denotes the source term, v the particle velocity, U the flow velocity and $\Gamma = (1-\frac{U^2}{c^2})^{-1/2}$.

$$\mathcal{E}(f) = \frac{\partial}{\partial \mu}\Big[D_{\mu\mu}\frac{\partial f}{\partial \mu} + D_{\mu p}\frac{\partial f}{\partial p}\Big] + \frac{1}{p^2}\frac{\partial}{\partial p}\Big(p^2[D_{\mu p}\frac{\partial f}{\partial \mu} + D_{pp}\frac{\partial f}{\partial p}]\Big) \quad (2)$$

is the Fokker-Planck wave-particle interaction term where the three Fokker–Planck coefficients $D_{\mu\mu}, D_{\mu p}, D_{pp}$ depend on the nature and statistical properties of the plasma wave turbulence.

Linear stability calculations show that magnetized plasmas contain low-frequency magnetohydrodynamic turbulence such as shear Alfven waves and fast and slow magnetosonic waves. Because for these plasma waves the magnetic part of the Lorentz force is much larger than the electric part of the Lorentz force, the time scale for rapid pitch angle scattering of energetic charged particles is much shorter than the time scale for energy changes. As a consequence, the cosmic ray particles' gyrotropic distribution function adjusts rapidly to quasi-equilibrium, which is close to the isotropic distribution function, in the fluid frame. By the standard diffusion approximation the diffusion–convection transport equation for the isotropic part of the phase space density $F(z,p)$ can then be derived from the quasilinear Fokker–Planck equation (1). For non-relativistic bulk speeds $U << c$ the diffusion-convection equation reads (Kirk et al. 1988)

$$-S_0 = \frac{\partial}{\partial z}\Big[\kappa\frac{\partial F}{\partial z}\Big] - V\frac{\partial F}{\partial z} + \frac{p}{3}\frac{\partial V}{\partial z}\frac{\partial F}{\partial p} + \frac{1}{p^2}\frac{\partial}{\partial p}\Big[p^2 A\frac{\partial F}{\partial p}F\Big], \quad (3)$$

where the spatial diffusion coefficient κ, the cosmic ray bulk speed V and the momentum diffusion coefficient A are determined by pitch-angle averages of three Fokker–Planck coefficients

$$\kappa = \frac{v^2}{8}\int_{-1}^{1} d\mu \frac{(1-\mu^2)^2}{D_{\mu\mu}(\mu)},$$

$$V = U + \frac{1}{3p^2}\frac{\partial}{\partial p}(p^3 D), \quad D = \frac{3v}{4p}\int_{-1}^{1} d\mu (1-\mu^2)\frac{D_{\mu p}(\mu)}{D_{\mu\mu}(\mu)}$$

$$A = \frac{1}{2}\int_{-1}^{1} d\mu [D_{pp}(\mu) - \frac{D^2_{\mu p}(\mu)}{D_{\mu\mu}(\mu)}]. \tag{4}$$

The three Fokker–Planck coefficients entering the averaging in equations (4) are calculated (Hall and Sturrock 1967, Krommes 1984, Achatz et al. 1991) from ensemble-averaged first-order corrections to the particle orbit,

$$D_{\mu\mu} = \Re \int_0^\infty d\tau \, \langle \dot{\mu}(t)\dot{\mu}^*(t+\tau)\rangle \,, \quad D_{\mu p} = \Re \int_0^\infty d\tau \, \langle \dot{\mu}(t)\dot{p}^*(t+\tau)\rangle \,,$$

$$D_{pp} = \Re \int_0^\infty d\tau \, \langle \dot{p}(t)\dot{p}^*(t+\tau)\rangle \,, \tag{5}$$

with $\vec{p} = (p\sqrt{1-\mu^2}\cos\phi, p\sqrt{1-\mu^2}\sin\phi, p\mu)$.

In its general form the diffusion–convection transport equation contains spatial diffusion and spatial convection terms as well as momentum diffusion and momentum convection terms. Since the pioneering work of Fermi (1949, 1954) it has become customary to refer to the latter two as Fermi acceleration of second and first order, respectively. Note, however, that the momentum convection term only leads to acceleration for converging bulk flow (i.e., $dV/dz < 0$) but to deceleration for expanding flows (i.e., $dV/dz > 0$). The converging bulk flow condition $dV/dz < 0$ is fulfilled at cosmic shock waves and leads to diffusive shock acceleration. We also note that in general the cosmic ray bulk speed V differs from the gas speed U by an extra term that depends on the cross helicity (i.e. the relative intensities of forward and backward moving waves) in the fluid rest frame.

The values of the three quasilinear transport parameters (4) depend on the nature and the statistical properties of the electromagnetic turbulence and the turbulence-carrying background medium. Idealized physical situations can be constructed where some of the three transport parameters (4) do not occur, e.g., in the magnetostatic approximation of the turbulence the parameters $A = D = 0$, so that the transport equation (3) in particular would contain no momentum diffusion term. Despite its frequent use, such a truncated transport equation is unrealistic and its applicability therefore rather limited.

Most cosmic plasmas have a small value of the plasma beta $\beta = c_S^2/V_A^2$, whose square root defined by the ratio of the ion sound to Alfven speed, and thus indicates the ratio of thermal to magnetic pressure. For low-beta plasmas the two relevant magnetohydrodynamic wave modes are the
(1) *shear Alfven waves* with dispersion relation

$$\omega_R^2 = V_A^2 k_\parallel^2 \tag{6}$$

at parallel wavenumbers $|k_\parallel| \ll \Omega_p/V_A$, where Ω_p denotes the non-relativistic proton gyrofrequency, and

(2) the *fast magnetosonic waves* or *fast mode waves* with dispersion relation

$$\omega_R^2 = V_A^2 k^2, \quad k^2 = k_\parallel^2 + k_\perp^2 \tag{7}$$

for wavenumbers $|k| \ll \Omega_p/V_A$.

In the limiting case (commonly referred to as slab model) of parallel (to $\vec{B}_0$) propagation ($\theta = k_\perp = 0$) the shear Alfven waves become the left-handed circularly polarised Alfven-ion-cyclotron waves, whereas the fast magnetosonic waves become the right-handed circularly polarised Alfven-Whistler-electron-cyclotron waves.

4. Diffusive acceleration at relativistic shocks

Since the pioneering work of Kirk and Schneider (1997) solutions of the Fokker-Planck equation (1) are derived for a Fokker-Planck interaction term

$$\mathcal{E}(f) \simeq \frac{\partial}{\partial \mu}\Big[D_{\mu\mu}\frac{\partial f}{\partial \mu}\Big] \tag{8}$$

in a step function shock profile

$$U = \begin{cases} U_+ & \text{in } z > 0 \\ U_- & \text{in } z < 0 \end{cases} \tag{9}$$

so that $\partial U/\partial z = 0 \ \forall z \neq 0$. The simplified Fokker-Planck operator (8) is appropriate for the low-frequency MHD modes (6) and (7) for which $D_{\mu p}$ and D_{pp} are of order $(V_A/v) << 1$ and $(V_A/v)^2 << 1$, respectively, smaller than $D_{\mu\mu}$. The steady-state Fokker-Planck transport equation (1) then reduces to

$$\Gamma_\pm(U_\pm + v\mu)\frac{\partial f_\pm}{\partial \mu} = \frac{\partial}{\partial \mu}\Big[D^\pm_{\mu\mu}\frac{\partial f_\pm}{\partial \mu}\Big] \tag{10}$$

on both sides of shock. Adopting shock properties (relation between U_+ and U_-) from relativistic MHD shock equations the solutions of Eq. (10) are obtained (for a recent review see Kirk and Duffy 1999) by expanding the solutions in terms of the eigenfunctions of the Sturm-Liouville-type pitch-angle scattering operator, and matching the solutions at $z = 0$, i.e. $f_+(z = 0, p, \mu) = f_-(z = 0, p, \mu)$.

Of course, the derived solutions depend strongly on the adopted form of the pitch angle Fokker-Planck coefficient $D^\pm_{\mu\mu}$ that contains all the microphysics of the particle-wave interaction. All existing analysis assume the same Fokker-Planck coefficient

$$D^+_{\mu\mu} = D^-_{\mu\mu} = D_{\mu\mu}(\mu, p) \tag{11}$$

on both sides of the shock, chosing a simple separable form

$$D_{\mu\mu}(\mu, p) = D_1(p) D_2(\mu).$$

According to our experience from non-relativistic shock acceleration, where the interaction of Alfven waves with the shock wave has been studied in greater detail, assumption (11) is very problematic and most probably not applicable.

5. Diffusive acceleration at nonrelativistic shocks

The transport equation (3) is used to investigate the test-particle diffusive acceleration of cosmic ray particles in quasi-parallel shock waves. Here finite plasma wave speed effects are important: this concerns in particular the cosmic ray bulk speed V in Eq. (4) that can be different from the gas speed U for nonzero values of $D \neq 0$.

Following earlier work by McKenzie and Westphal (1969) and Scholer and Belcher (1971), Vainio and Schlickeiser (1998, 1999, 2001) calculated anew the transmission of small-amplitude parallel-moving Alfven waves through a parallel super-Alfvenic shock. In their investigation Vainio and Schlickeiser combined the equations for
(1) the continuity of the transverse momentum

$$\left[\rho U_n \vec{U}_t - \frac{B_n \vec{B}_t}{4\pi}\right] = 0, \tag{12}$$

where the shock bracket $[X] \equiv X_1 - X_2$ denotes the difference of the upstream (index 1) and downstream (index 2) value of the physical quantity X, U_n (B_n) and $\vec{U}_t$ ($\vec{B}_t$) are the normal (to the shock) and tangential gas flow velocity (magnetic field) components, respectively;
(2) the continuity of the normal magnetic field

$$B_{n,1} = B_{n,2} = B_0; \tag{13}$$

(3) the continuity of the tangential electric field

$$\left[U_n \vec{B}_t - B_n \vec{U}_t\right] = 0; \tag{14}$$

(4) and the continuity of the mass flux

$$\left[\rho U_n\right] = 0 \tag{15}$$

with the different relation of velocity and magnetic field fluctuations for forward (f) and backward (b) moving Alfven waves, i. e.

$$\delta \vec{U}^{\mathrm{f}} = -\frac{\delta \vec{B}^{\mathrm{f}}}{(4\pi\rho)^{1/2}}, \quad \delta \vec{U}^{\mathrm{b}} = \frac{\delta \vec{B}^{\mathrm{b}}}{(4\pi\rho)^{1/2}}. \tag{16}$$

To arrive at a complete set of equations for the downstream values, i.e., to be able to determine the gas compression ratio $r = \rho_2/\rho_1 = U_{n,1}/U_{n,2}$ of the shock, Eqs. (12) – (16) must be completed by yet two equations (e.g., Boyd and Sanderson 1969) describing the continuity of the normal momentum

$$\left[\rho U_n^2 + P + \frac{B_t^2}{8\pi}\right] = 0 \tag{17}$$

and energy flux (for an adiabatic equation of state, $P\rho^{-\gamma_g} = \text{const.}$)

$$\left[\frac{1}{2}\rho U_n (U_n^2 + U_t^2) + \frac{\gamma_g P U_n}{\gamma_g - 1} + \frac{U_n B_t^2}{4\pi} - \frac{B_n(\vec{U}_t \cdot \vec{B}_t)}{4\pi}\right] = 0, \tag{18}$$

respectively.

One then arrives (Vainio and Schlickeiser 1999) at a cubic equation for the shock compression r that includes the influence of finite Alfven wave pressure and that can be solved in a parametric form. For $\gamma_g = 5/3$ THe Alfvenic Mach number $M = U_{n,1}/V_{A,1}$ obeys

$$M^2 = (1+y)r(y) \tag{19}$$

$$r(y) = \frac{8y^2(y+1) - 6\beta y^2 - q_{L,1}(y+1)(5y-3)}{2y^2(y+1) + q_{L,1}(y+1)(y+3)}, \tag{20}$$

where $q_{L,1} = (\delta B_1/B_0)^2$, and the parameter y runs between

$$\frac{\beta - 1 + q_{L,1} + \sqrt{(\beta + 1 + q_{L,1})^2 - 4\beta}}{2} < y < \infty. \tag{21}$$

Moreover, with this approach *the downstream electromagnetic field properties can be calculated from the specified upstream electromagnetic field.* In particular, specifying the upstream Alfven wave cross helicity state $H_{c,1}$, that indicates the relative fraction of forward and backward moving Alfven waves, so that the upstream cosmic ray bulk speed is $V_1 = U_1 + H_{c,1}V_{A,1}$, Vainio and Schlickeiser calculated the resulting downstream cosmic ray bulk speed $V_2 = U_2 + H_{c,2}V_{A,2}$. This immediately yields the scattering center compression ratio

$$r_k \equiv \frac{V_1}{V_2} = r\frac{M + H_{c,1}}{M + r^{1/2}H_{c,2}}, \tag{22}$$

which in general is different from the gas compression ratio r. Because the shock wave is collisionless, *it is this scattering center compression ratio, and not the gas compression ratio, that determines the spectral index of the power law momentum spectrum of the accelerated cosmic ray particles.* When downstream momentum

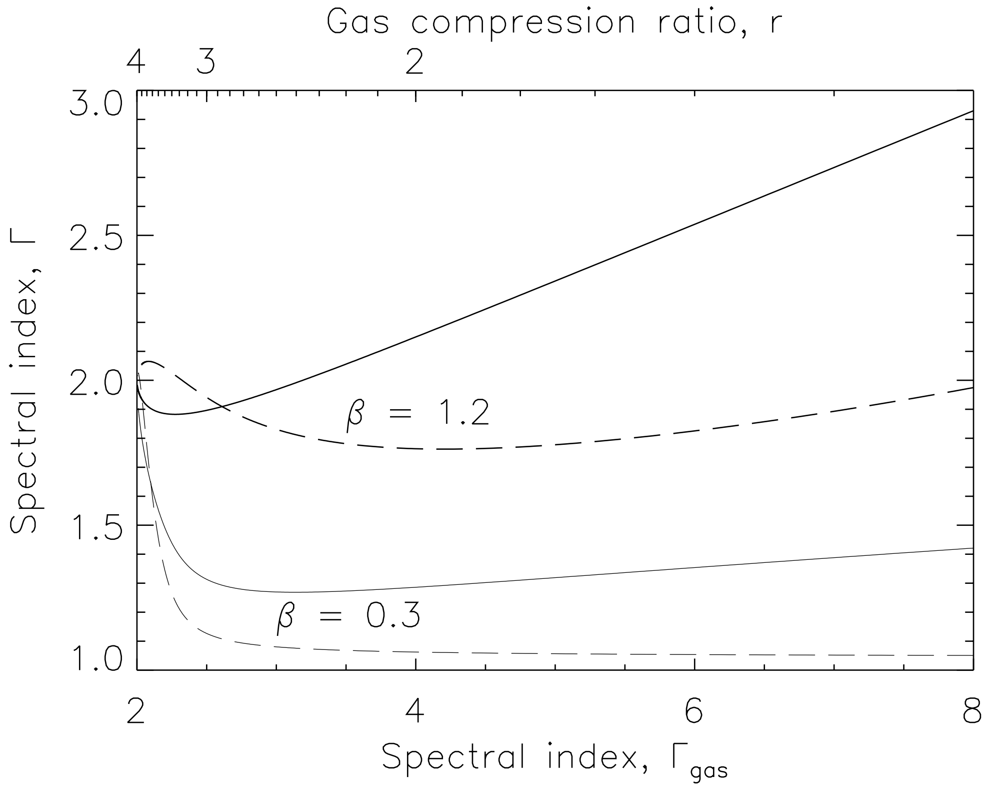

Fig. 1. Cosmic-ray spectral index produced by an adiabatic shock with a constant upstream plasma beta, neglecting stochastic acceleration in the downstream region. Dashed and solid lines give the results for $H_{c,1} = +1$ and -1, respectively. The magnetic amplitude of the upstream waves is $b = 0.1$ and their spectral index is $q = 1.5$. The Alfven wave normal momentum and energy flux are includeded in deriving the shock's gas compression ratio. From Vainio and Schlickeiser (1999).

diffusion is neglected, the particle differential energy spectrum at the shock is — up to a cut off determined by losses, particle escape, finite geometrical shock extent and/or finite acceleration time — a power law in momentum

$$dJ/dE \propto p^{-\Gamma}, \quad \Gamma = \frac{r_k + 2}{r_k - 1}, \tag{23}$$

whose spectral index Γ is solely determined by r_k. In Fig. 1 we show the calculated cosmic ray spectral index values as a function of the spectral index of the conventional theory, $\Gamma_{\rm gas} = (r+2)/(r-1)$ and as a function of the gas compression ratio for four different specified upstream states and an adiabatic index $\gamma_g = 5/3$.

In all cases, the scattering center compression ratio r_k differs significantly from the gas compression ratio r. Practically never does the scattering center compression ratio r_K agree with the gas compression ratio r. In particular, for low upstream plasma beta small spectral index values $\Gamma \leq 2$ are possible, whereas the gas compression ratio value is limited to $r \leq 4$ for $M \rightarrow \infty$. Thus, the model, being able to generate particle power law spectra harder than the originally limiting value $\Gamma = 2$, avoids the discrepancy noted by, e.g., Lerche (1980), Drury (1983) and Dröge et al. (1987) that the original shock wave acceleration theory in its simplest test-particle form is not in accord with the observed flat particle spectra in shell-type supernova remnants and bright spiral galaxies. And it again is keeping track of different propagation speeds of forward and backward moving waves that leads to this significant result.

The principal difference between the gas compression ratio and the scattering center compression ratio, being equivalent to the difference between the effective plasma wave velocity and the gas velocity, and the possible consequences for the spectral index of the differential momentum spectrum of accelerated particles, has been already noted by Bell (1978), see his Eqs. (11) and (12); although he did no quantitative calculations of this effect. By calculating the correct transmission coefficients of Alfven waves through the shock from the Rankine-Hugoniot continuity equations Vainio and Schlickeiser (1988, 1999) demonstrated that precisely this effect can account for the generation of particle spectral indices flatter than $\Gamma = 2$.

6. Summary and conclusions

The acceleration of energetic charged particles at relativistic and non-relativistic shock waves plays an important role in powering established and potential cosmic TeV-gamma-ray sources as supernova remnants, jets of active galactic nuclei and gamma-ray bursts. I have reviewed some recent results on non-relativistic shock acceleretion and described the current standing of the theory for relativistic shock acceleration.

There are three important clues from our insight into non-relativistic shock acceleration:

1) The *microphysics* of the plasma wave - shock interaction *is important* not only for the resulting momentum spectrum of accelerated particles (difference between gas compression ratio and scattering center compression ratio) but also for the shock wave structure (influence of finite wave pressure on the shock's cubic (20)).

2) *Cosmic ray transport parameters* $D_{\mu\mu}, D_{\mu p}, D_{pp}$ for the Fokker-Planck equation and κ, V, A for the diffusion-convection equation *have to be calculated* self-consistently from the shock properties including in particular the interaction of plasma waves (transmission and reflection) with the shock. As a consequence, the transport parameters are significantly different downstream from upstream! This

finding should have consequences also for the assumptions made in relativistic shock acceleration theories.
3) The dynamics of plasma waves near shocks is crucial for particle acceleration. Especially the interaction of plasma waves with the shock wave change the cosmic transport parameters significantly from upstream to downstream. As an alternative it might be useful to study simplified models of plasma wave dynamics at the interface of outflows and interstellar medium, ignoring in the beginning the complications arising from the possible existence of a gaseous shock. One such model is the relativistic pick-up model of Pohl and Schlickeiser (2000).

All three finding should have corresponding consequences also for studying particle acceleration in relativistic shocks and may guide further analytical developments of relativistic shock acceleration theories.

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