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The three-fluid structure of the particle modulated solar wind termination shock

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Abstract. In this paper we have improved our earlier physical concept to theoretically describe the particle-modulated solar wind termination shock by three dynamically interacting hydrodynamic fluids, i.e. the solar wind, the pick-up ion, and the anomalous cosmic ray (ACR) plasmas. In this earlier concept we have introduced a parametrized form to describe the injection of Fermi-I energized pick-up ions into the ACR regime in the precursor and the subshock region. With this parametrization it was possible to give solutions for the upstream-to-downstream transition of all thermodynamic fluid properties on the basis of two undefined parameters, i.e. the total injection efficiency of pick-up ions and the mean energy of post-shock ACR particles. Making use of the pick-up ion transport equation we now obtain a parameter-free description of the injection process. In addition we obtain the energy-averaged spatial diffusion coefficient as function of the upstream distance from the subshock by use of an approximated ACR energy spectrum. As a consequence of these improvements we now can give solutions for the spatial structure of the three-fluid system in the precursor and the subshock region in terms of bulk velocity, fluid pressures, and compression ratio, without undefined parameters. We discuss the observability of the derived multifluid post-shock plasma conditions by energetic neutral atoms (ENA's) and emphasize the fact that ENA's may be the best tracers to the physics of the particle-modulated solar wind termination shock.

Key words: solar system: multi-fluid shocks – cosmic rays – heliospheric termination shock – pick-up ions

1. Introductory sketch of the shock scenario

Pick-up ions are produced all over the heliosphere by ionization of interstellar neutral atoms (e.g. Fahr & Rucinski 1989, Ratkiewicz et al. 1990, Fahr et al. 1992, Rucinski et al.1993, Fahr et al. 1995, Mall et al. 1996). After pick-up by solar wind magnetic fields they undergo rapid pitch angle scattering, and on more extended time scales adiabatic deceleration and energization by Fermi-II processes and by transit time damping mechanisms (Fisk et al.1974; Fisk 1976a/b; Vasyliunas & Siscoe 1976; Klecker 1977; Möbius et al. 1985, 1988; Isenberg 1987; Fahr & Ziemkiewicz 1988; Bogdan et al. 1991, Chalov et al. 1995, 1997; Fichtner et al. 1996, Schwadron et al 1996). Results obtained by Pesses et al. (1981), Potgieter & Moraal (1988), Jokipii (1990, 1992) clearly suggest that the solar wind termination shock and the region ahead of it will act as converters of such energized pick-up ions into ACR particles by effective Fermi-I and shock drift acceleration processes. ACRs being enough energized eventually can diffuse upstream from the shock and appear in the inner solar system with typical energies of about 10 MeV/nucleon. The energy density of ACR particles close to the termination shock is apparently rather high to modify the shock structure. Two-fluid models of cosmic-ray-modified shock waves were proposed by Drury & Völk (1981), Axford et al. (1982), Achterberg et al. (1984), Kang & Jones (1990), Jones & Kang (1990), Donohue & Zank (1993), Zank et al. (1993), Krülls & Achterberg (1994), Chalov & Fahr (1994, 1995a/b), Ziemkiewicz (1994).

We have recently presented a three-fluid model of a ACR-modulated one- dimensional shock structure where we treated the pick-up ion plasma as a separate fluid which reduces the effective preshock solar wind Mach number and the strength of the shock (Chalov & Fahr 1996). This extrafluid though considered to comove with the normal solar wind has a polytropic behaviour different from that of the solar wind ions meaning that it is a thermally independent plasma component. The model includes a continuous energy injection from the pick-up fluid into the ACR fluid all over the region of decelerated plasma flow. We then could give solutions for the upstream-to-downstream transition of the coupled three-fluid system keeping two undefined parameters, namely the total energy injection into the ACR fluid and the average energy of the post-shock ACR particles. Here we have now studied more carefully the physics of this injection process on the basis of an adequate pick-up ion transport equation and thereby could remove the uncertainty in its earlier parametrized representation. Furthermore we have cal-

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culated the energy-averaged spatial ACR diffusion coefficient as a function of the upstream distance from the subshock and thus were able to present solutions for the spatial structure of the coupled three-fluid system in the precursor and subshock region.

Here we only consider the most important pick-up ion component which is pick-up protons. Their production due to H-atom ionizations has essentially been completed prior to the entrance of the solar wind into the precursor region. The production of fresh pick-up ions within the precursor (i.e. at the last $10AU$ ahead of the shock) and their effect on the dynamics of the flow has been neglected here.

2. Basic theoretical approach

2.1. Governing equations

Here we briefly review the theoretical approach used in our earlier paper (Chalov & Fahr 1996). At moderate energies the transport equation for pick- up ions can be solved under the neglect of spatial diffusion compared to the convection. Thus the solar wind and the pick-up ion plasmas are treated as distinct fluids with different thermodynamical behaviours, i.e. different temperatures, pressures and number densities, but with identical bulk velocities u. Since the radial extent of the region of an ACR-modified solar wind flow upstream of the termination shock at its solar distance of about $100AU$ is estimated by $\Delta r = 10AU$(Chalov and Fahr, 1996), we can assume that an adequate separation of space variables due to the different scales of derivatives exists. Instead of the radial symmetry we thus may simply use a one-dimensional approximation since: $(1/r^2 d/dr(r^2\rho u) = d/dr(\rho u) + 2(\rho u)/r \cong d/dr(\rho u)$. In a one-dimensional approximation the equations of conservation of mass, momentum, and energy for the mixture of the solar wind plasma, pick-up ions, and ACRs then attain the form:

$$\rho u = \mu, \tag{1}$$

$$\mu u + p_g + p_i + p_c = \text{const}, \tag{2}$$

$$\mu \left\{ u^2 + \frac{\gamma_g}{\gamma_g - 1} \left(\frac{p_{rmg}}{\rho} + \frac{p_i}{\rho} \right) \right\} + F_c = \text{const}, \tag{3}$$

where $\rho = \rho_g + \rho_i$ is the density of the mixture of the solar wind plasma and pick-up ions, u and p_g are the velocity and pressure of the solar wind plasma, p_i and p_c are the pressures of the pick-up ions and ACRs. γ_g is the adiabatic index of both the solar wind plasma and the pick-up ions. The value μ denotes the constant total mass flow across the shock structure. In addition to the conservation of the mass flux of the solar wind/pick-up ion mixture given by Eq. (1), we also have: $\mu_g = \rho_g u = \rho_{g1} u_1$, $\mu_i = \rho_i u = \rho_{i1} u_1$. The quantity F_c is the cosmic ray energy flux given by:

$$F_c = \frac{\gamma_c}{\gamma_c - 1} u p_c - \frac{\bar{\kappa}}{\gamma_c - 1} \frac{dp_c}{dx}, \tag{4}$$

where γ_c is the adiabatic index of the anomalous plasma, $\bar{\kappa}$ is the energy- averaged, scalar diffusion coefficient, and x is the space coordinate. We take the x-axis as coinciding with the solar wind bulk flow direction. The flux F_c is connected with the ACR pressure p_c by the diffusion equation:

$$\frac{dF_c}{dx} = u \frac{dp_c}{dx} + Q(x). \tag{5}$$

Here Q is the energy injection rate describing energy gains of the anomalous particle regime from pick-up ions by the spatially extended action of Fermi-I acceleration processes. We shall again assume that $Q(x)$ is given by our earlier expression (Chalov & Fahr 1996):

$$Q(x) = -\alpha p_i \frac{du}{dx}, \tag{6}$$

where α is introduced as a factor. The exact value of α is calculated in Appendix A from basic phase space transport theory.

2.2. The precursor flow

We assume that the local energy input to ACRs as given by Eq. (6) is extracted from pick-up ion energies, while the flow of the solar wind fluid behaves adiabatic:

$$p_g = p_{g1} / f^{\gamma_g}, \tag{7}$$

where $f = u/u_1$, and u_1 and p_{g1} are the velocity and the pressure at an arbitrary point of the flow. Then from Eqs. (2), (3), (5), and (7) one obtains an equation for the pick-up ion pressure:

$$u \frac{dp_i}{dx} + \gamma_g p_i \frac{du}{dx} = -(\gamma_g - 1) Q(x). \tag{8}$$

From this equation we obviously can derive a polytropic relation for pickup ions in the form (see Chalov & Fahr, 1995a/b, 1996)

$$p_i = p_{i1} / f^{\sigma}, \tag{9}$$

with the index σ given by

$$\sigma = \gamma_g - \alpha(\gamma_g - 1). \tag{10}$$

As evident σ is directly connected with the efficiency factor α used in Eq. (6) to describe a spatially extended energy injection into the ACR regime.

Furthermore we suppose that the most energetic ACRs (ACR's above a certain energy threshold) leak out from the shock structure in upstream direction assuming that these particles lose dynamical contact to the shock-induced turbulences (i.e. $\kappa(E) \to \infty$) and thus to the thermal plasma. To account for this process we require that within the concept of a finite extent of the precursor at its entrance $x = x_1$ the following relation holds:

$$\begin{aligned} F_{c1} &= \lim_{x \to x_1} \left[\frac{\gamma_c}{\gamma_c - 1} u p_c - \frac{\bar{\kappa}}{\gamma_c - 1} \frac{dp_c}{dx} \right] \\ &\cong -\frac{\kappa_{\max}}{\gamma_c - 1} | \frac{dp_c}{dx} |_{x_1}. \end{aligned} \tag{11}$$

Consequently at the point x_1 there must exist a diffusive ACR energy flux F_{c1} directed upstream and representing an energy loss from the shock system.

Locating the subshock at $x = 0$, then one can estimate: $|x_1| \cong \kappa(E_{max})/u_1$, where $\kappa(E_{max})$ is the kinetic, i.e energy-dependent spatial diffusion coefficient and E_{max} is the lower energy threshold of particles which escape from the shock structure. In planar approximation the functions ρ, u, p_g, and p_i can be taken to be constants at $x \leq x_1$, adopting values ρ_1, u_1, p_{g1}, and p_{i1} there. With Eqs. (2) and (11) we obtain:

$$p_c = \mu u_1(1-f) + p_{g1}(1-f^{-\gamma_g}) + p_{i1}(1-f^{-\sigma}). \qquad (12)$$

and with Eqs. (2), (3), (7), (9), and (12) we are led to the following differential equation for the function f:

$$\frac{df}{dx} = -\frac{u_1}{\bar{\kappa}(x)} \frac{G(f)}{\left[1 - 1/(M_{g1}^{*2} f^{1+\gamma_g}) - \sigma/(\gamma_g M_{i1}^{*2} f^{1+\sigma})\right]}, \qquad (13)$$

where

$$\begin{aligned} G(f) = & \frac{\gamma_c(1-f)}{\gamma_g M_{g1}^{*2}} \left\{ \gamma_g M_{g1}^{*2} \frac{\gamma_c+1}{2\gamma_c}(f-\beta_c) \right. \\ & + \frac{\gamma_g-\gamma_c}{\gamma_c(\gamma_g-1)} \frac{1-f^{1-\gamma_g}}{1-f} - 1 \\ & \left. + \frac{M_{g1}^{*2}}{M_{i1}^{*2}} \left[\frac{\gamma_g-\gamma_c}{\gamma_c(\gamma_g-1)} \frac{1-f^{1-\sigma}}{1-f} - 1 \right] \right\} \\ & - \frac{(\gamma_c-1)}{2} \frac{F_{c1}}{\mu u_1^2/2}. \end{aligned}$$

In Eq. (13) $\beta_c = (\gamma_c - 1)/(\gamma_c + 1)$, M_{g1}^* and M_{i1}^* are the solar wind and pick-up ion effective Mach numbers given by the formulae:

$$M_{g1}^{*2} = u_1\mu/\gamma_g p_{g1}, \, M_{i1}^{*2} = u_1\mu/\gamma_g p_{i1}.$$

The differential Eq. (13) has to be solved with the boundary condition $f(x_1) = 1$. Eq. (13) now contains an unknown energy loss term

$$\epsilon = -\frac{F_{c1}}{\mu u_1^2/2}, \qquad (14)$$

which has to be determined within the solution of a closed system of governing equations.

The kinetic diffusion coefficient $\kappa = \kappa(E)$ of the ACRs is an increasing function of the kinetic energy E. This results in a spatial dependence of the ACR energy spectrum in the precursor (e.g. Eichler 1984, Lee 1982, le Roux & Fichtner 1997). As a consequence the energy-averaged diffusion coefficient $\bar{\kappa}$ and the adiabatic index γ_c are functions of the upstream subshock distance x. As already stated in Chalov & Fahr (1996) this spatial dependence of $\bar{\kappa}$ and γ_c, can, however, not consistently be determined within the framework of a purely hydrodynamic description. In the present paper we nevertheless obtain an approximate solution for $\bar{\kappa} = \bar{\kappa}(x)$ when using some available theoretical knowledge on the expected ACR velocity distribution function in the precursor region (see Appendix B). In the case of a completely smooth shock transition, the downstream velocity far from the transition must become constant, i.e. $|df/dx|_\infty = 0$, and thus is found by:

$$G(f_2) = 0. \qquad (15)$$

Our earlier calculations did show that the solar wind termination flow contains a subshock structure under all reasonable sets of parameters. Hence we shall consider only transitions including a subshock in this paper here.

2.3. Subshock transition relations

If the denominator at the rhs of Eq. (13) can vanish within the interval $f_2 \leq f \leq 1$, where the value f_2 is a solution of Eq. (15), then no smooth transition exists. The flow thus must undergo a dissipative subshock. If we integrate Eqs. (1) - (3), and (5) in an infinitesimal vicinity of the subshock and take into account that the ACR pressure should behave continuously, then we obtain the following jump relations at the subshock:

$$[\mu] = 0, \qquad (16)$$

$$[\mu u + p_g + p_i] = 0, \qquad (17)$$

$$\mu\left[u^2/2 + \frac{\gamma_g}{\gamma_g-1}\frac{p_g+p_i}{\rho}\right] + [F_c] = 0, \qquad (18)$$

$$p_c = 0, \qquad (19)$$

$$F_c = 0, \qquad (20)$$

where

$$q = \lim_{\in\to 0}\int_{-\in}^{+\in} Q(x)\mathrm{d}x = (\alpha/2)(u_0 - u_2)(p_{i0} + p_{i2}). \qquad (21)$$

Here the subshock location is defined as $x = 0; u_0(u_2)$ and $p_{i0}(p_{i2})$ are the plasma velocity and pick-up ion pressure in front of (behind) the subshock.

The above system of jump relations (16) - (20) is not yet closed. For the sake of its closure one needs an additional relation derivable with the following physical argument (see also Chalov & Fahr 1996): In view of the fact that the resulting subshock, due to the small pick-up ion Mach number of $M_i = (\mu_i u/\gamma_g p_i)^{1/2} \cong 1$ (see Sect. 3), for pick-up ions represents a very weak shock (or not shock at all) it seems reasonable to assume that the pick-up ion flow passes isentropic over the subshock when no injection into the ACR regime takes place at the subshock, i.e. when "$q = 0$" or $\alpha = 0$ is valid. This assumption seems to be confirmed by hybrid simulations of the interaction of pick-up ions with the solar wind termination shock as run e.g. by Liewer et al. (1993), or Kucharek & Scholer (1995). As could be proven by the results of these simulations which due to limitations in available computer times are not run up to the occurence of ACR-typical particle energies (10 MeV/nucleon) the downstream pick-up ion temperatures are even less increased as should be expected from an adiabatic reaction on the basis of the increased downstream densities.

In case,however, of an injection of the more energetic pick-up ions into the ACR regime due to consecutive Fermi-1 acceleration processes, i.e. in our treatment for the case "$q > 0$", then the passage of pick-up ions over the subshock consequently has to result even in a decrease of the pick-up ion entropy $s_{\rm i}$ at the subshock. This conclusion also is, if not supported, at least is not invalidated by test particle calculations of shockreflected particles by Zank et al. (1996), Lee et al.(1996), or Gedalin (1996). Concerning the latter, his post-shock pick-up ion component by its velocity distribution (see his Fig. 3) does not seem to show anything else but an adiabatic reaction of the downstream pick-up ions connected with their density increase by a factor of $n_{\rm i2}/n_{\rm i1} = B_2/B_1 = 3$.

For the entropy change per mass one thus can write:

$$ds_{\rm i} = dq^*/T_{\rm i}. \tag{22}$$

In Eq. (22) $dq^* = -q/\mu_{\rm i}$ is the heat energy exchanged between pick-up ions and ACRs due to injection at the subshock, and the entropy $s_{\rm i}$ per mass is given by:

$$s_{\rm i} = c_{\rm v} \log(p_{\rm i}/\rho_{\rm i}^{\gamma_{\rm g}}) + \text{const.}$$

Integrating Eq. (22) over the subshock thus gives

$$\lim_{\in \to 0} \int_{-\in}^{+\in} T_{\rm i}(x) \Delta s_{\rm i}(x) \delta(x) dx = -q/\mu_{\rm i}, \tag{23}$$

where $\Delta s_{\rm i}$ is the jump in the pick-up ion entropy through the subshock and $\delta(x)$ is the Dirac delta function. Now, if we assume that $p_{\rm i} = \rho_{\rm i} R T_{\rm i}$, where $R = c_{\rm p} - c_{\rm v}$, we obtain from Eq. (23):

$$\log \left[\frac{p_{\rm i2}}{p_{\rm i0}} \left(\frac{u_2}{u_0} \right)^{\gamma_{\rm g}} \right] \left(\frac{p_{\rm i0}}{\rho_{\rm i0}} + \frac{p_{\rm i2}}{\rho_{\rm i2}} \right) + \frac{2(\gamma_{\rm g} - 1)q}{\mu_{\rm i}} = 0. \tag{24}$$

The flow of the solar wind, pick-up ions, and ACRs in the precursor is described by Eqs. (7), (9), (11), (12), and (13). From Eqs. (1), (12), (16), and (19) we obtain:

$$\rho_2 = \rho_1/z f_0, \tag{25}$$

$$p_{\rm c2} = \mu u_1 (1 - f_0) + p_{\rm g1}(1 - f_0^{-\gamma_{\rm g}}) + p_{\rm i1}(1 - f_0^{-\sigma}), \tag{26}$$

where $z = u_2/u_0$ and $f_0 = u_0/u_1$.

If we introduce a new variable λ by

$$p_{\rm i2} = \lambda p_{\rm i0}, \tag{27}$$

then the post-shock pressure of the solar wind plasma with the use of Eqs. (7), (9), and (17) is given by

$$p_{\rm g2} = \mu u_1 (1 - z) f_0 + \frac{p_{\rm g1}}{f^{\gamma_{\rm g}}} + (1 - \lambda) \frac{p_{\rm i1}}{f^{\sigma}}. \tag{28}$$

Inserting $p_{\rm i2}$ from Eq. (27) into Eq. (24) we can obtain the following equation for λ:

$$\alpha(\gamma_{\rm g} - 1)(1 - z)(1 + \lambda) + (1 + \lambda z) \log(\lambda z^{\gamma_{\rm g}}) = 0. \tag{29}$$

From Eq. (18) we obtain the equation for the velocity jump z across the precursor:

$$z = \frac{\gamma_{\rm g}-1}{\gamma_{\rm g}+1} \left\{ 1 + \left[\frac{2\gamma_{\rm g}}{\gamma_{\rm g} - 1} - \alpha(1 + \lambda) \right] \frac{1}{\gamma_{\rm g} M_{\rm i1}^{*2} f_0^{\sigma+1}} + \frac{2}{\gamma_{\rm g} - 1} \frac{1}{M_{\rm g1}^{*2} f_0^{\gamma_{\rm g}+1}} \right\}. \tag{30}$$

From Eq. (20) we have:

$$(1 - f_0) \left[1 + f_0 - \frac{2\gamma_{\rm c2}}{\gamma_{\rm c2} - 1} z \left(f_0 - \frac{1}{\gamma_{\rm g} M_1^{*2}} \right) \right] + 2 \left(\frac{\gamma_{\rm g}}{\gamma_{\rm g} - 1} - \frac{\gamma_{\rm c2}}{\gamma_{\rm c2} - 1} z \right) \left[\frac{1 - f_0^{1-\gamma_{\rm g}}}{\gamma_{\rm g} M_{\rm g1}^{*2}} + \frac{1 - f_0^{1-\sigma}}{\gamma_{\rm g} M_{\rm i1}^{*2}} \right] + \frac{\alpha(1 - z)(1 + \lambda)}{\gamma_{\rm g} M_{\rm i1}^{*2} f_0^{\sigma-1}} - \epsilon = 0, \tag{31}$$

where

$$\frac{1}{M_1^{*2}} = \frac{1}{M_{\rm g1}^{*2}} + \frac{1}{M_{\rm i1}^{*2}}.$$

For the derivation of Eq. (31) we assumed that the flow behind the subshock is uniform, i.e has constant velocity, density, and pressures. On the other hand $dp_{\rm c0}/dx$ appearing on the lhs of Eq. (20) was found from Eqs. (12) and (13). If we now insert z given by Eq. (30) into Eqs. (29), (31), we obtain two equations for the three unknown quantities, namely f_0 (the compression across the precursor), ϵ (the energy loss), and λ. These equations also contain free parameters: $M_{\rm g1}^*$, $M_{\rm i1}^*$, $\gamma_{\rm c2}$, and α. In Appendix B we obtain a third additional relation connecting f_0, ϵ, and λ, allowing a closure of the system of Eqs. (29), (30), (31). In addition in Appendix A we estimate the value of the injection efficiency α by use of results concerning stochastic acceleration of pick-up ions in the solar wind obtained by Chalov et al. (1997).

An interesting quantity in question is η describing the effective number of pick-up ions injected into the ACR regime per unit of time. For the sake of a determination of this quantity we use the fact that the integral $\int_{x_1}^{\infty} Q dx$ gives the total energy input to ACRs per unit of time. Thus the total flux of injected particles is given by

$$F_{\rm inj} = \int_{x_1}^{\infty} Q dx / E_{\rm inj}, \tag{32}$$

where $E_{\rm inj}$ is the injection energy. From Eq. (32) one can find the fraction of pick-up ions converted to ACR:

$$\eta = m_{\rm p} \int_{x_1}^{\infty} Q dx / \mu_{\rm i} E_{\rm inj} = m_{\rm p} \left[\lim_{\in \to 0} \int_{x_1}^{-\in} Q dx + q \right] / \mu_{\rm i} E_{\rm inj}. \tag{33}$$

In Eq. (33) $m_{\rm p}$ is the proton mass and $\mu_{\rm i} = \rho_{\rm i} u = \rho_{\rm i1} u_1$.

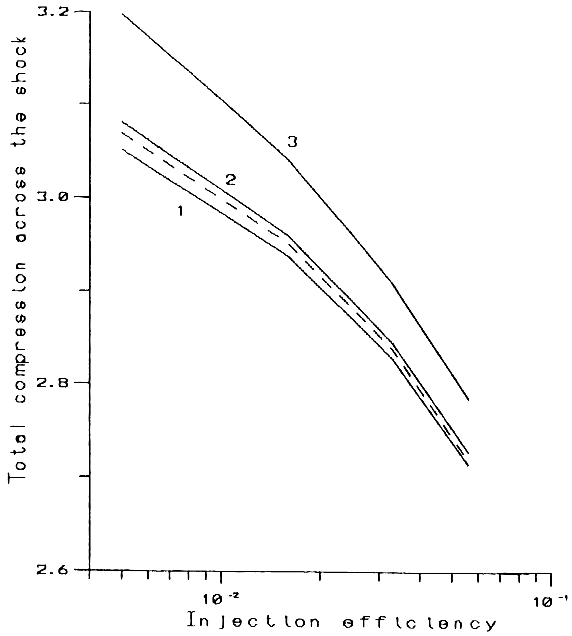

Fig. 1. The total compression ratio ρ_2/ρ_1 across the shock as a function of the injection efficiency α for different forms of the momentum dependence of the kinetic diffusion coefficient κ : **1** : $-\delta$ = 2.0, **2** : $-\delta$ = 1.0, **3** : $-\delta$ = 0.5. The dashed curve corresponds to κ given by Eq. (37).

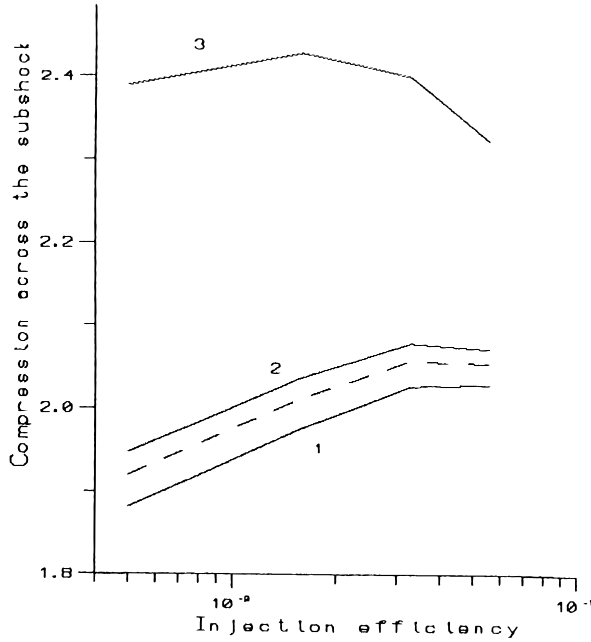

Fig. 2. The compression across the subshock as a function of the injection efficiency α for different forms of the momentum dependence of κ(see the caption in Fig. 1).

3. Numerical results and conclusions

In this paper we assume that $E_{\rm max}$ = 300 MeV and $M^*_{\rm g1}$ = 100. It was recently shown by Chalov et al. (1997) that stochastic acceleration of pick-up ions through the heliosphere by Alfvenic turbulence and large-scale fluctuations of the solar wind including travelling interplanetary shock waves results in formation of a high energy tail in the energy distribution of the ions at their travel to the outer parts of the heliosphere. In front of the termination shock the energy of the accelerated pick-up ions can reach several hundred KeV. At these energies the spatial diffusion becomes important for the pick-up ion transport in the precursor region and we shall consider particles from this high energy tail as a seed population for ACRs. Thus it seems to be reasonable to admit here that $E_{\rm inj}$ = 100 keV. What concerns the solar wind velocity upstream of the precursor of the termination shock we take a value of $u_1 = 387 km/s$ resulting from calculations of Baranov & Malama (1993, 1995) for the case that both the electron and neutral hydrogen number densities in the LISM are adopted with $0.14 cm^{-3}$, and that the electron number density and the solar wind velocity at the Earth's orbit are $7 cm^{-3}$ and $450 km/s$, respectively. Such deceleration of the supersonic solar wind is connected with the influence of interstellar neutrals on the plasma flow (see also Isenberg 1986). The ratio of the number densities of pick-up ions and solar protons is then consistently obtained with: $\chi = n_{\rm i1}/n_{\rm g1} = 0.17$.

Taking into account that $p_i = \rho_i < v_{\rm i}^2 > /3$, where $< v_{\rm i}^2 >$ is the mean square of the thermal velocity of the pick-up ions, we then can write:

$$M_{\rm i1} = \left[\frac{3}{\gamma_{\rm g} \frac{u_1^2}{<v_{i1}^2>}} \right]^{1/2} \tag{34}$$

The pick-up Mach number $M_{\rm i1}$ is connected with the effective Mach number $M^*_{\rm i1}$ by the relation:

$$M^*_1 = M_{\rm i1} \sqrt{1 + \frac{1}{\chi}} \tag{35}$$

We have mentioned above that in the paper by Chalov et al.(1997) the stochastic acceleration of pick-up ions by solar wind turbulences through the heliosphere has been investigated and energy spectra of these ions in front of the termination shock have been calculated.From those calculations we can obtain the mean energy of pick-up ions $< E_{\rm i1} >$ in front of the termination shock as variable with the intensity of the mean Alfvenic turbulence at the Earth's orbit. Then we can calculate the Mach number $M_{\rm i1}$ by use of Eq. (34) ($< E_{\rm i1} >= m_{\rm p} < v_{\rm i1}^2 >$) and injection efficiency α by use of Eq. (A9). The results of these calculations are presented in Table 1. One can see there that in case of strong stochastic acceleration of pick-up ions in the solar wind the flow of pick-up ions close to the termination shock is classified as subsonic.

As explained in Appendix B we need to specify the energy-dependence of the kinetic diffusion coefficient of ACRs to solve Eq. ($B10$). We consider the two following cases:

$$a) \ \kappa \propto p^{\delta} \tag{36}$$

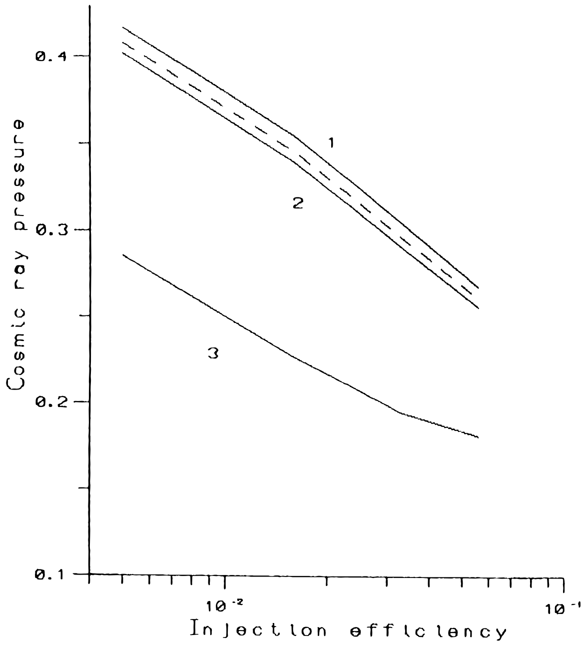

Fig. 3. The normalized cosmic ray pressure behind the shock $p_{cN} = p_{c2}/(p_{g2} + p_{i2} + p_{c2})$ as a function of the injection efficiency α for different forms of the momentum dependence of κ (see the caption in Fig. 1).

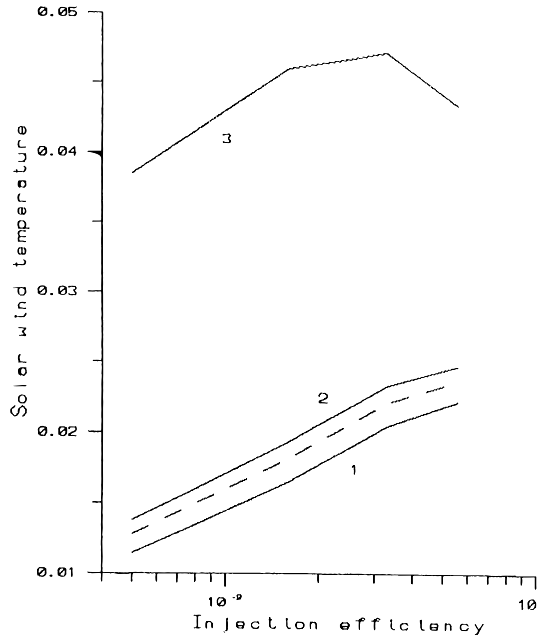

Fig. 4. The normalized post-shock temperature of the solar wind $2kT_{g2}/m_p u_1^2$ as a function of the injection efficiency α for different forms of the momentum dependence of κ (see the caption in Fig. 1).

Table 1. The mean energy, Mach number, and injection efficiency of pick-up ions as functions of the intensity of Alfvenic turbulence

$\langle \tilde{B}_E^2/B_E^2 \rangle$	$< E_{i1} >$ (keV)	M_{i1}	α
0.1	1.00	1.18	0.005
0.2	1.16	1.10	0.016
0.3	1.33	1.03	0.033
0.4	1.52	0.96	0.056

and

$$b)\ \kappa \propto \beta \begin{cases} P/1\text{GV} & \text{at } P \geq 0.4\text{GV} \\ 0.4 & \text{at } P < 0.4\text{GV} \end{cases} \quad (37)$$

where $\beta = v/c$, P is the rigidity of particles. The diffusion coefficient in the form given by Eq. (37) has been used by Potgieter & Moraal (1988) to study acceleration and modulation of ACRs in the heliosphere.

Figs. 1 through 4 show - the total compression across the shock, $zf_0 = \rho_2/\rho_1$, the partial compression z across the subshock, the normalized cosmic ray pressure behind the shock, $p_{cN} = p_{c2}/(p_{g2} + p_{i2} + p_{c2})$, and the normalized post-shock temperature of the solar wind $2kT_{g2}/m_p u_1^2$, all as functions of the injection efficiency α for different forms of a momentum-dependence of the kinetic diffusion coefficient $\kappa(p)$. *The curves 1, 2, and 3 correspond to* $\delta = 2.0, 1.0,$ *and 0.5, respectively (see Eq. (36)).* The dashed curves correspond to the diffusion coefficient given by Eq. (37). As we have mentioned above, the ratio of the pick-up ion to solar wind number densities is $\chi = 0.17$.

One can see from Fig. 1 that the total compression across the termination shock ranges from 2.7 to 3.2 in dependence on α (or on the intensity of solar wind turbulence, see Table 1). This relatively low value for the compression is directly connected with the influence of pick-up ions considerably decreasing the effective Mach number of the flow, i.e. weakening the shock (Baranov & Malama 1993, 1995). In this respect it is interesting to notice that in a recent paper Stone et al.(1996), by examining the ACR energy spectra measured at different distances from the shock by the Voyager and Pioneer spacecraft in the period from 1993 and 1994, have concluded that to explain these spectra with modulation theory the strength of the termination shock had to be 2.63 ± 0.14 over this period. This estimation is surprisingly close to our present theoretical values.

The role of ACRs consists in the formation of a smooth precursor in front of the dissipative subshock. When the kinetic diffusion coefficient is given by Eq. (37) which seems to be more adequate to describe the ACR transport the compression across the subshock is close to 2 (dashed curve in Fig. 2). It may appear strange at first glance that the cosmic ray pressure decreases when the injection efficiency increases (see Fig. 3). However, the explanation for this outcome is fairly simple: The mean energy and size of the high energy tail of pick-up ions depend on the level of solar wind turbulence which can effectively accelerate particles (Chalov et al. 1997). The increase of the magnitude of solar wind turbulences, on one hand, results in the increase of the injection efficiency of the pick-up ions into the ACR population. On the other hand, it also reduces the effective Mach number of the flow and hence the total compression at the shock (see Fig. 1). Decrease of the shock strength, in turn, reduces the efficiency of the acceleration of ACRs. It follows from our calculations that the pressure of ACRs represents between 25% to 40% of the total pressure of the three-fluid mixture.

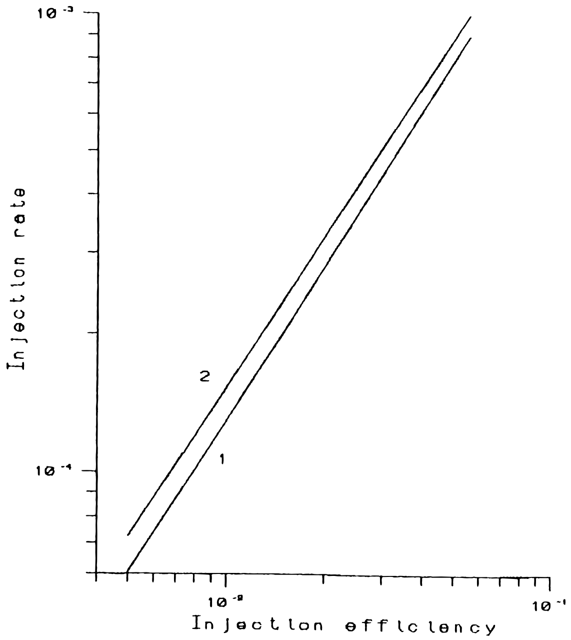

Fig. 5. The fraction of pick-up ions converted to ACR η as a function of the injection efficiency α for different forms of the momentum dependence of the kinetic diffusion coefficient κ : **1** : $-\delta$ = 2.0, **2** : $-\delta$ = 0.5.

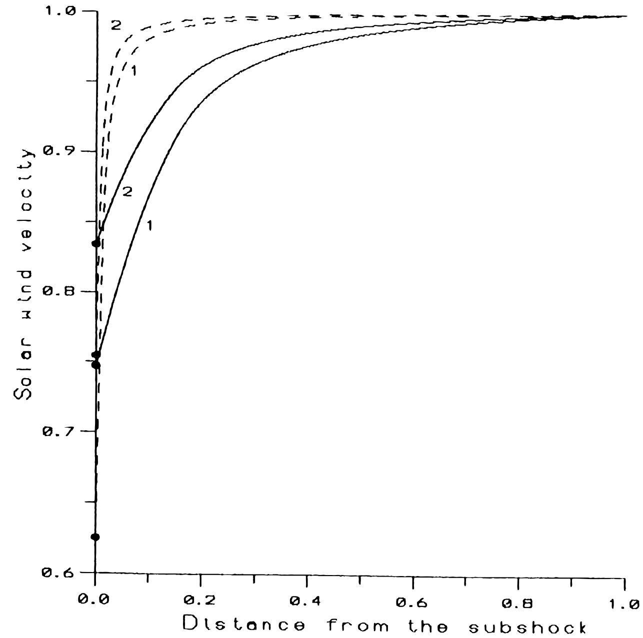

Fig. 6. The solar wind velocity in the precursor $f = u/u_1$ as a function of the distance from the subshock x/x_1. Solid curves correspond to κ given by Eq. (36) with δ = 0.5; dashed ones correspond to κ given by Eq. (37). **1** : $-\alpha$ = 0.005, **2** : $-\alpha$ = 0.056.

Table 2. Downstream parameters of the termination shock with vanishing cosmic ray pressure

M_{i1}	ρ_2/ρ_1	$2kT_{g2}/m_p u_1^2$
1.18	3.05	0.129
1.10	2.94	0.118
1.03	2.83	0.108
0.96	2.71	0.097

The deceleration of the plasma flow in the precursor by the ACR pressure gradient decreases the velocity jump in the dissipative subshock. This results in the strong decrease of the post-shock temperature of the solar wind plasma. In Table 2 we present downstream parameters of the termination shock for the case of vanishing ACR pressure (i.e.there is no injection of pick-up ions into the ACR regime) to compare ACR-modulated shocks with classical, purely hydrodynamical ones. If we compare the results from Table 2 with those shown in Fig. 4 we can conclude that the post-shock temperature of the solar plasma in the case of an ACR-modulated shock decreases by a factor of 5 to 10 compared to the value of the purely hydrodynamical shock wave.

Fig. 5 shows the fraction of pick-up ions converted to ACRs (injection rate) η as a function of the injection efficiency α for δ = 2.0 (curve 1) and δ = 0.5 (curve 2). From this Figure it follows that the relative number of injected particles varies from $6 \cdot 10^{-5}$ to 10^{-3} dependent on the level of Alfvenic turbulence. Stone et al. (1996) propose $(2.4 \pm 1.1) \cdot 10^{-4}$ for the hydrogen injection rate in 1993 and 1994.

Fig. 6 shows the solar wind velocity in the precursor, by $f = u/u_1$, as a function of the distance from the subshock x/x_1. Solid curves correspond to the kinetic diffusion coefficient given by Eq. (36) with δ = 0.5 and dashed ones correspond to κ given by Eq. (37). Curves labelled by **1** are plotted for α = 0.005, and by **2** for α = 0.056. We can see from these velocity profiles that in that case, when the kinetic diffusion coefficient is rapidly increasing with energy, the deceleration of the plasma flow in the precursor mainly takes place in a narrow region close to the subshock (dashed curves).

It is now interesting to compare in a little more detail differences between the classical and "non-classical" post-shock plasma conditions connected with the particle-modulated three-fluid shock. The best possibility to open up an observational access to these distant post-shock plasma conditions, before the deep space probes may enter into this region at perhaps some time in the near future, is to look for ENA (energetic neutral atom) particles produced by charge exchange reactions in this post-shock region. Here mainly post-shock solar wind ions and pick-up ions are decharged by H-atoms by means of charge exchange processes. Gruntman (1992) has calculated such ENA fluxes (called HELENA, s) arriving at the Earth's orbit. His calculations were made on the basis of the post-shock plasma conditions derived within the frame of the hydrodynamical twin-shock model by Baranov (1990). This latter model treats the solar wind plasma as a hydrodynamical monofluid entering the shock with a high pre-shock Mach number. In this model, as a consequence of the Rankine-Hugoniot relations applied, the post-shock plasma appears as a very hot plasma com-

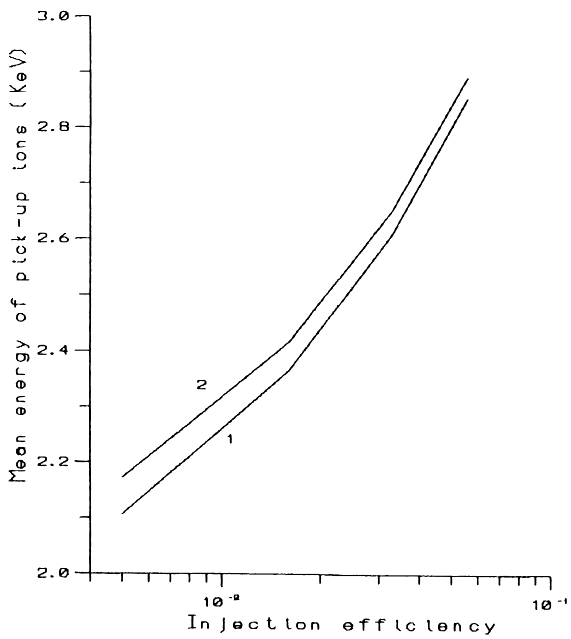

Fig. 7. The mean energy of pick-up ions (keV) behind the shock as a function of the injection efficiency α for different forms of the momentum dependence of the kinetic diffusion coefficient κ : **1** : $-\delta = 2.0$, **2** : $-\delta = 0.5$.

ponent. From this plasma Gruntman (1992) calculates ENA-spectra which peak at energies of about 200 eV.

Now, in view of the above mentioned results for the particle-modulated three-fluid shock the properties of the post-shock plasma appear substantially changed as compared to the monofluid post-shock conditions. It may hence be interesting to briefly study in qualitative terms the expected influences on the ENA-spectra resulting from these new, "non-classical" conditions. As we have shown in Fig. 4, the consistent inclusion of the dynamical action of pick-up ions and ACR's substantially modifies the shock transition. Due to both the decrease of the effective solar wind Mach number by the pick-up ion presence and to the smooth deceleration of the solar wind by the ACR-pressure gradient in the precursor, the subshock now becomes much weaker as in the classical case and, as a consequence, there is less dissipative heating and hence the resulting post-shock solar wind proton temperatures are much lower, amounting only to 5 to 10 percent of the classical values. The ENA-particles originating from the decharging of this plasma component thus will show energy spectra peaking at much lower energies as compared to Gruntman's calculations (i.e. less than 20eV). On the other hand, amongst the multifluid post-shock plasma constituents there are also post-shock pick-up ions which have non-negligible relative abundances (about 20%). As we have shown in Fig. 7 these pick-up ions have mean energies of the order of a few KeV which means that their decharging by charge exchange leads to an additional form of ENA's at this energy level, i.e. leading to ENA spectra which peak at the KeV-level. This in fact may very much advice to look for such ENA's since they could be the only clear tracer at the moment for the multi-fluid shock transition phenomena which were described above. In addition the postshock ACR's will also be reflected by ENA particles which was already discussed earlier in the literature (see Grzedzielski et al., 1993).

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Appendix A: evaluation of the injection efficiency

Stochastic acceleration of pick-up ions in the heliosphere is a very important process (see Isenberg 1987, Chalov et al. 1995, 1997, Fichtner et al. 1996) by which pick-up ions with energies of even up to some 100 KeV can be produced (Chalov et al. 1997) and we shall consider these energetic pick-up ions as a seed population for ACRs. However, the process of stochastic acceleration operates on longer time scales and thus effectively only acts over large heliospheric scales. In the present approach we assume that stochastic preacceleration of pick-up ions has been completed upstream of the precursor. Pick-up ions entering the precursor are here only subject to Fermi-I acceleration processes and thereby produce ACRs. In the precursor region described in the one-dimensional approach the angle-integrated differential number density $f_i(t, x, v)$ of pitch-angle isotropized pick-up ions in phase space, at the absence of energy and spatial diffusion, is described by the following kinetic equation:

$$\frac{\partial f_i}{\partial t} + u\frac{\partial f_i}{\partial x} = \frac{1}{3}\frac{\partial u}{\partial x}v\frac{\partial f_i}{\partial v}, \tag{A1}$$

where v is the speed of isotropized pick-up ions in the solar wind rest frame. The pressure of pick-up ions is then given by (Drury & Völk 1981):

$$p_{\mathrm{i}} = \frac{1}{3}\int_0^{p_{\mathrm{inj}}} vp^3 f_{\mathrm{i}} dp, \tag{A2}$$

where p is the pick-up ion momentum, and p_{inj} is the lower threshold for the momentum needed for injection into the ACR regime. Studying the work of Lee (1982) and Potgieter and Moraal (1988) one may conclude that this threshold momentum is connected with the subshock compression ratio z, but in any case can be considered as a quantity which is constant over the precursor.

Therefore by partial integration in Eq. (A1) one can derive in the case of a steady-state flow the following differential equation for p_i as velocity moment of f_{i}:

$$u\frac{dp_{\mathrm{i}}}{dx} + \frac{5}{3}p_{\mathrm{i}}\frac{du}{dx} = \frac{du}{dx}\frac{f_{\mathrm{i}}(x, p_{\mathrm{inj}})}{9}p_{\mathrm{inj}}^4 v_{\mathrm{inj}}. \tag{A3}$$

If we compare now Eq. (A3) with Eq. (8) we can write:

$$Q(x) = -\frac{f_{\rm i}(x, p_{\rm inj})}{9(\gamma_{\rm g} - 1)} p_{\rm inj}^4 v_{\rm inj} \frac{du}{dx}. \quad (A4)$$

On the other hand we have given Q by Eq. (6) introducing there the undefined factor α. Thus from Eqs. (6) and (A4) one now can obtain the following relation for this factor:

$$\alpha = \frac{f_{\rm i}(x, p_{\rm inj})}{9(\gamma_{\rm g} - 1) p_{\rm i}(x)} p_{\rm inj}^4 v_{\rm inj}. \quad (A5)$$

Expressing the integral in Eq. (A2) by the help of the mean-value theorem we obtain:

$$p_i(x) \simeq \frac{1}{3}(1/m_{\rm i}) < p_i(x) >^5 f_{\rm i}(x, < p_{\rm i} >), \quad (A6)$$

where the expressions in brackets, i.e.$< .. >$, mean appropriate average values. Inserting $p_{\rm i}$ from Eq. (A6) into Eq. (A5) then leads to:

$$\alpha = \frac{1}{3(\gamma_{\rm g} - 1)} \left(\frac{E_{\rm inj}}{< E_{\rm i} >} \right)^{5/2} \frac{f_{\rm i}(x, p_{\rm inj})}{f_{\rm i}(x, < p_{\rm i} >)}. \quad (A7)$$

One can see from Eq. (A7) that rigorously taken α appears as a function of the variable x. In Sect. 2 we did, however, derive our system of equations adopting a constant or weakly variable value for α which is justified if:

$$\frac{d}{dx} \ln n_i \gg \frac{\gamma_{\rm g} - 1}{\gamma_{\rm g} - \alpha(\gamma_{\rm g} - 1)} \frac{d\alpha}{dx}$$

is valid. To obtain an appropriate value for α from Eq. (A7) we shall now evaluate the right hand side of this equation at the point x_1, i.e. at the entrance to the precursor. Instead of $f_{\rm i}$ which is not available we prefer to use here differential number densities of pick-up ions as they are directly measured by plasma analyzers on board of space probes. In the paper by Chalov et al. (1997) differential number densities $N_{\rm i}(E)$ of accelerated pick-up ions in the region upstream of the termination shock have been calculated (e.g.see their Fig. 6). These number densities $N_{\rm i}(E) = p^2 f_{\rm i}(dp/dE)$ are related to $f_{\rm i}(p)$ by:

$$\frac{f_{\rm i}(p_{\rm inj})}{f_{\rm i}(< p_{\rm i} >)} = \left(\frac{E_{\rm inj}}{< E_{\rm i} >} \right)^{-1/2} \frac{N_{\rm i}(E_{\rm inj})}{N_{\rm i}(< E_{\rm i} >)}. \quad (A8)$$

To derive Eq. (A8) we used the nonrelativistic $p - E$ relation which is accurate for energies in question here. Thus finally we obtain α by:

$$\alpha = \frac{1}{3(\gamma_{\rm g} - 1)} \left(\frac{E_{\rm inj}}{< E_{\rm i} >} \right)^2 \frac{N_i(E_{\rm inj})}{N_{\rm i}(< E_{\rm i} >)}. \quad (A9)$$

Appendix B: determination of the local diffusion coefficient

As one notices Eqs. (29) to (31) still contain the free parameter ϵ, i.e. the upstream energy loss from the shock. The occurence of this free parameter is caused by the fact that we could not yet solve Eq. (13) since it contains the unknown functions $\bar\kappa(x)$ and $\gamma_{\rm c}(x)$. Concerning $\gamma_{\rm c}$ we can assume here that $\gamma_{\rm c} = 5/3$, since the non-relativistic limit is a fairly accurate approximation for ACR energies all over the precursor. To obtain a reasonable representation for $\bar\kappa(x)$ we shall go back to the rigorous definition of the energy-averaged diffusion term in the energy-averaged transport equation written in Eq. (4). According to this definition the diffusion coefficient $\bar\kappa(x)$ is given by:

$$\bar\kappa(x) = \frac{\int_{p_{\rm inf}}^{p_{\rm max}} \kappa(p) p^3 v \frac{\partial f_{\rm c}}{\partial x} dp}{\int_{p_{\rm inj}}^{p_{\rm max}} p^3 v \frac{\partial f_{\rm c}}{\partial x} dp}, \quad (B1)$$

where $f_{\rm c}(x, p)$ is the differential number density of cosmic rays in phase space and $\kappa(p)$ is the energy-dependent diffusion coefficient. Following Eichler (1984, 1985), Krymsky (1984), and Ellison & Eichler (1985) we assume that the phase space distribution of cosmic rays in the precursor can be written in the form:

$$f_{\rm c}(x, p) = g(p) H[x - \bar x(p)], \quad (B2)$$

where $H[..]$ is the Heaviside step function and $|\bar x(p)|$ is the mean penetration depth of particles with the momentum p diffusing from the subshock upstream into the precursor region (n.b.: $x; \bar x < 0$). The representation of $f_c(x, p)$ by Eq. (B2) is appropriate for the sake of determining the coefficient $\bar\kappa(x)$, since $\kappa(p)$ is rapidly increasing with energy. Since $\bar x = \bar x(p)$ is monotonically increasing with p, one can also obtain the inverse function: $p = p(\bar x)$ in a unique form. With this function one then obtains:

$$\int_{p_{\rm inj}}^{p_{\rm max}} \kappa(p) p^3 v \frac{\partial f_{\rm c}}{\partial x} {\rm d}p = \int_{p_{\rm inj}}^{p_{\rm max}} \kappa(p) p^3 v g(p) \delta[x - \bar x(p)] {\rm d}p =$$

$$\int_{\bar x_{\rm inj}}^{\bar x_{\rm max}} \kappa(\bar x) p^3(\bar x) v(\bar x) g(\bar x) \delta[x - \bar x(p)] \frac{\partial p}{\partial \bar x} {\rm d}\bar x =$$

$$\kappa(x) p^3(x) v(x) g(x) \frac{\partial p}{\partial \bar x}|_{\bar x = x}$$

and proceeding in analogous manner for the nominator of Eq. (B1)

$$\int_{p_{\rm inj}}^{p_{\rm max}} p^3 v \frac{\partial f_{\rm c}}{\partial x} {\rm d}p = \int_{p_{\rm inj}}^{p_{\rm max}} p^3 v g(p) \delta[x - \bar x(p)] {\rm d}p =$$

$$\int_{\bar x_{\rm inj}}^{\bar x_{\rm max}} p^3(\bar x) v(\bar x) g(\bar x) \delta[x - \bar x(p)] \frac{\partial p}{\partial \bar x} {\rm d}\bar x =$$

$$p^3(x) v(x) g(x) \frac{\partial p}{\partial \bar x}|_{\bar x = x} \quad (B3)$$

where $\bar x_{\rm max} = \bar x(p_{\rm max}), \bar x_{\rm inj} = \bar x(p_{\rm inj})$. The relations (B3) both are valid for $x < \bar x_{\rm inj}$, i.e the mean upstream penetration depth of ACRs injected at the subshock with $p_{\rm inj}$.

Thus with Eq. (B1) and the above relations (B3) we obtain for $x < \bar x_{\rm inj}$:

$$\bar\kappa(x) = \kappa[p(\bar x = x)]. \quad (B4)$$

The mean penetration depth of particles in the upstream direction can now be found as the solution of the following equation (i.e diffusion-convection limit):

$$\bar{x} = -\kappa(p)/u(\bar{x}). \tag{B5}$$

If we insert now $\kappa(p)$ from Eq. (B5) into the right hand side of Eq. (B4) changing $\bar{x}$ by x, then we finally obtain:

$$\bar{\kappa}(x) = -u_1 f(x) x \text{ for} : x < \bar{x}_{\text{inj}}. \tag{B6}$$

Since we assume here that the diffusion coefficient κ does not depend on the spatial variable x, it thus can be assumed for $x_{\text{inj}} \leq x$ that also the energy-averaged expression $\bar{\kappa}$ does not depend on x in this region. Instead one shall admit that

$$\bar{\kappa}(x) = \kappa(p_{\text{inj}}) \qquad (\text{at } x > \bar{x}_{\text{inj}}). \tag{B7}$$

The value of $\bar{x}_{\text{inj}}$ is found from the equation which follows directly from Eq. (B5):

$$\bar{x}_{\text{inj}} f(\bar{x}_{\text{inj}}) = -\kappa(p_{\text{inj}})/u_1. \tag{B8}$$

The length of the precursor can be found from the relation:

$$\bar{x}_{\text{inj}} f(\bar{x}_{\text{inj}}) = -\kappa(p_{\text{inj}})/u_1. \tag{B9}$$

Let us introduce now a dimensionless spatial variable $\xi = x/x_1 > 0$. Then Eq. (13) can be rewritten in the form:

$$\frac{\mathrm{d}f}{\mathrm{d}\xi} = \frac{G(f)}{H(\xi, f)[1 - 1/(M_{\text{g1}}^{*2} f^{1+\gamma_{\text{g}}}) - \sigma/(\gamma_{\text{g}} M_{\text{i1}}^{*2} f^{1+\sigma})]}, \tag{B10}$$

where

$$H(\xi, f) = \begin{cases} \xi f(\xi) & \text{at } \xi > \xi_{\text{inj}} \\ \kappa(p_{\text{inj}})/\kappa(p_{\max}) & \text{at } \xi < \xi_{\text{inj}} \end{cases}. \tag{B11}$$

To obtain Eq. (B11) we have used Eqs. (B6) to (B9). In Eq. (B11) ξ_{inj} is the solution of the equation:

$$\xi_{\text{inj}} f(\xi_{\text{inj}}) = \kappa(p_{\text{inj}})/\kappa(p_{\max}). \tag{B12}$$

The differential Eq. (B10) has to be solved with the boundary condition: $f(1) = 1$. If in addition we demand that $f(0) = f_0$, then by Eq. (B10) one gains an additional relation which closes the system of Eqs. (29)- (31). For this purpose, however, one has to use the knowledge on the momentum-dependence of the diffusion coefficient $\kappa(p)$ (see Kota & Jokipii 1993, le Roux & Potgieter 1993).

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