A MODEL OF ACCELERATION OF ANOMALOUS COSMIC RAYS BY RECONNECTION IN THE HELIOSHEATH

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ABSTRACT

We discuss a model of cosmic ray acceleration that accounts for the observations of anomalous cosmic rays (ACRs) by Voyager 1 and 2. The model appeals to fast magnetic reconnection rather than shocks as the driver of acceleration. The ultimate source of energy is associated with magnetic field reversals that occur in the heliosheath. It is expected that the magnetic field reversals will occur throughout the heliosheath, but especially near the heliopause where the flows slow down and diverge with respect to the interstellar wind and also in the boundary sector in the heliospheric current sheet. While the first-order Fermi acceleration theory within reconnection layers is in its infancy, the predictions do not contradict the available data on ACR spectra measured by the spacecraft. We argue that the Voyager data are one of the first pieces of evidence favoring the acceleration within regions of fast magnetic reconnection, which we believe to be a widely spread astrophysical process.

Key words: interplanetary medium – MHD – solar neighborhood – solar wind – turbulence

1. INTRODUCTION

Since the crossing of the termination shock (TS) by Voyager 1 (V1) in late 2004 and by Voyager 2 (V2) in mid-2007, it became clear that several paradigms needed to be revised. Among them was the acceleration of particles. Prior to the encounter of the TS by V1, the prevailing view was that anomalous cosmic rays (ACRs) were accelerated at the TS by diffusive shock acceleration (DSA) to energies 1–300 MeV nuc−1 (e.g., Jokipii & Giacalone 1998; Cummings & Stone 1998). However, with the crossing of the TS by V1, the energy spectrum of the ACR did not unroll to the expected source shape: a power law at lower energies with a roll off at higher energies. After 2004, both the V1 spectrum in the heliosheath and the V2 spectrum upstream of the TS continued to evolve toward the expected source shape.

To explain this paradox, several models were proposed. Among them, McComas & Schwadron (2006) suggested that at a blunt shock the acceleration site for higher energy ACRs would be at the flanks of the TS, where the injection efficiency would be higher for DSA and connection times of the magnetic field lines to the shock would be longer, allowing acceleration to higher energies. Fisk et al. (2006), on the other hand, suggested that stochastic acceleration in the turbulent heliosheath would continue to accelerate ACRs and that the high-energy source region would thus be beyond the TS. Other works, such as Jokipii (2006) and Florinski & Zank (2006) tried to explain the deficit of ACRs based on a dynamic TS. Jokipii (2006) pointed out that a shock in motion on timescales of the acceleration time of the ACRs, days to months, would cause the spectrum to differ from the expected DSA shape. Florinski & Zank (2006) calculated the effect of magnetic interacting regions (MIRs) with the TS on the ACR spectral shape. They showed that there is a prolonged period of depressed intensity in mid-energies from a single MIR. Other recent works have recently included stochastic acceleration, as well as other effects (Moraal et al. 2006; Zhang 2006; Langner & Potgieter 2006; Ferreira et al. 2007). It became clear after the crossing of the TS by V2 that these models would require adjustments. The observations by V2 indicate, for example, that a transient did not cause the modulation shape of the V2 spectrum at the time of its TS crossing. When both spacecraft are in the heliosheath in late 2007, the radial gradient in the 13–19 MeV nuc−1 ions does not appear to be caused by a transient. The 60–74 MeV nuc−1 ions have no gradient, so no north–south or longitudinal asymmetry is observed in the ACR intensities at the higher energies.

Here we propose an alternative model, which explains the source of ACRs as being in the heliosheath. In what follows, we appeal to magnetic reconnection as a process that can accelerate particles. We explain the origin of the magnetic field reversals that induce magnetic reconnection in the heliosheath and heliopause in Section 2.

It is known that magnetic fields are well frozen-in into space plasma, which causes magnetic fields and plasma to move together. However, as soon as two magnetic flux tubes attempt to cross each other, magnetic field lines of different directions come sufficiently close to each other so that the magnetic field rearrangement takes place. This topological rearrangement converts free energy of the magnetic field into the energy of plasma motion and plasma heating. Many phenomena, including solar flares, the Earth magnetospheric events, γ-ray bursts, are accepted to be powered by magnetic reconnection, although the detailed mechanism of the process stayed illusive for a long time (see Biskamp 2000; Priest & Forbes 2000; Bhattacharjee 2004; Zweibel & Yamada 2009). Recent progress achieved in understanding the nature of reconnection processes allows us to appeal to the process with more confidence. In Section 3 we discuss the magnetic reconnection in the heliosheath and propose a model of reconnection, which is generated by turbulence on the large scales, but exhibits properties of collisionless reconnection on microscales.

The acceleration of energetic particles by reconnection is a subject in the state of development. In Section 4, we provide an outlook on the processes of energetic particle acceleration in reconnection regions and identify the first-order Fermi acceleration arising from energetic particle bouncing between the reconnecting fluxes as the dominant acceleration process.

We accept the exploratory nature of this work and compare our proposed solution with the alternative solutions of the
ACRs problem in Section 5. There we also discuss the existing limitations of our understanding of the complex processes that we invoke in our model and how these uncertainties affect our conclusions.

2. NATURE OF MAGNETIC FIELD REVERSALS EXPECTED IN THE HELIOSHEATH

It is well known that the magnetic field in the heliosphere changes polarity and creates current sheets. For instance, as the Sun rotates the magnetic field twists into a Parker spiral (Parker 1958) with magnetic fields separated by a current sheet (see Schatten 1971). The changes of the magnetic field are also expected due to the solar cycle activity.

The question now is at what part of the heliosheath do we expect to see reversals. The structure of the magnetic field in the solar wind is complex. The solar magnetic field lines near the TS are azimuthal and form a spiral (see Figure 1 that shows a global view of the interaction of the solar wind with the interstellar wind). The spiral solar magnetic field is shown being deflected at the heliopause (shown in dark dashed lines). The heliopause itself is deflected by an interstellar magnetic field. (figure adapted from S. Suess 2006, private communication).

As the HCS crosses the TS, the spiral gets tighter and tighter due to the slow-down of the solar wind flow as it approaches the heliopause (see Figure 1). The thickness of the outflow regions in the reconnection region depends on the level of turbulence. The length of the outflow regions L depends on the mean geometry of the magnetic field and turbulence.

Near the heliopause, the magnetic sectors get tighter and tighter (see Figure 3). We expect that reconnection might play a major role especially in this region, as the magnetic fields of one sector in the heliosheath are pressed against the interstellar magnetic field draped around the heliopause. For example, considering an upstream solar velocity of 450 km s\(^{-1}\) the separation between the boundary sectors upstream of the TS, \(\lambda\), will be \(\sim 7\) AU, while immediately downstream the velocity drops by approximately a factor of 3–4 to 100 km s\(^{-1}\) so \(\lambda\) will drop to 2.0 AU. As the velocity approach the heliopause, the velocity drops to zero.

Taking a velocity of 10–20 km s\(^{-1}\), the wavelength is reduced to 0.4–0.2 AU close to the heliopause (see Figure 1).

Qualitatively speaking, after the solar magnetic field lines cross the TS, the spiral gets tighter and tighter due to the slow-down of the solar wind flow as it approaches the heliopause (see Figure 1). The thickness of the outflow regions in the reconnection region depends on the level of turbulence. The length of the outflow regions L depends on the mean geometry of the magnetic field.

In situ measurements in the heliosheath (see Burlaga et al. 2009) show the presence of magnetic sectors beyond the TS. We expect as well the solar cycle to affect the magnetic sectors. In analytic studies, Nerney et al. (1995) predicted a complex region in the heliosheath due to the solar cycle effects. Note that during each solar cycle the solar global polarity reverses as well. This will create opposite polarity sectors due to solar cycle separating strongly mixed polarities (due to the tilt of the HCS discussed above).

Figure 3 shows how these two effects will affect the overall structure of the heliosheath (adapted from Nerney et al. 1995). The dark and faint green represents the opposite solar polarity sectors while the white represents the regions with strongly mixed polarities due to the varying tilt of HCS in each solar cycle. The heliopause is represented by a blue line and the TS by a red line. The heliosheath upstream of the TS is the far right of the figure. Their studies, however, were done in the kinematic approximation where the magnetic field reaction on the flow was neglected. The flow pattern in the heliosheath is expected to be very complex and can affect the overall picture above.
Another effect is related to the fact that as the solar wind smashes against the interstellar wind, the solar magnetic field increases in intensity near the solar equatorial plane creating magnetic ridges where the interstellar pressure is stronger (see Opher et al. 2003, 2004). Thus we expect the solar magnetic field to have a very complex structure exhibiting magnetic field reversals beyond the TS, in the heliosheath (see Nerney et al. 1995), as shown in Figure 3.

3. MODEL OF RECONNECTION IN THE HELIOSHEATH

Currents in resistive plasma drain their energy from the magnetic field. However, we shall discuss further that the default rates of magnetic energy conversion that can be obtained from a naive treatment of the problem are too small to be important for the acceleration of energetic particles.

For instance, the famous Sweet–Parker model of reconnection (Sweet 1958; Parker 1958; see Figure 4, upper panel) produces reconnection rates that are smaller than the Alfvén velocity by a square root of the Lundquist number, that is, by $S^{-1/2} = (LV_A/\eta)^{1/2}$, where $L$ is the length of a current sheet, $V_A$ is the Alfvén velocity, $\eta$ is the Ohmic diffusivity. The current sheet length is determined by transversal extend of the magnetic flux tubes that get into contact and for the heliopause it can be larger than 100 AU. The corresponding reconnection speed for the Sweet–Parker reconnection in this case is negligible, it being $\sim 3 \times 10^{-7}V_A$. This result can be understood intuitively, as in the Sweet–Parker reconnection plasma collected over the size $L$ should be ejected with the speed $\sim V_A$ from the thin slot $\delta_{SP} = LS^{-1/2}$, where $S$ for the heliosheath is about $10^{13}$ (see below). The disparity of sizes $L$ and $\delta_{SP}$ makes the reconnection slow and thus one can disregard the effect of the Sweet–Parker mechanism for heliospheric reconnection altogether.

The first model of fast reconnection by Petschek (1964) assumed that magnetic fluxes get into contact not along the astrophysically large scales of $L$, but instead over a scale comparable to the resistive thickness $\delta$, forming a distinct X-point, where magnetic field lines of the interacting fluxes converge at a sharp point to the reconnection spot. The stability of such a reconnection geometry in astrophysical situations is an open issue. At least for uniform resistivities, this configuration was proven to be unstable and to revert to the Sweet–Parker configuration (Biskamp 1986; Uzdensky & Kulsrud 2000).

To understand the effect of plasma instabilities (including Hall-MHD effects), we need to know the connection length $L_{SP}$ for which Hall-MHD effects should be important for the reconnection. The criterion at which the Hall-MHD term gets important for the reconnection is that the ion skin depth $\delta_{ion}$ is comparable to the Alfvén gyroradius of an ion moving at the Alfvén speed, that is, $\delta_{ion} = V_A/\omega_{ci}$, where $\omega_{ci}$ is the cyclotron frequency of an ion.

For the parameters of the heliosheath given in Table 1, we find that for a proton it is $\delta_{ion} \sim 10^3$ km. Thus, one can get the constraint on the scale $L$ for which Hall-MHD effects should dominate the reconnection:

$$\frac{\delta_{SP}}{\delta_{ion}} \approx 0.2 \left( \frac{L}{\lambda_{mfp}} \right)^{1/2} \rho_{pl}^{1/4} \ll 1,$$

where $\rho_{pl}$ is the electron plasma density, and $\lambda_{mfp}$ is the mean free path of a proton. In general, the reconnection is termed fast when the reconnection velocity does not depend on the Lundquist number $S$ or if it depends on $\ln(S)$. In all other cases, the large values of $S$ make reconnection too slow for most of astrophysical applications.

### Table 1

Parameters of the Heliosheath (from Burlaga et al. 2009)

| Symbol | Name                  | Magnetic Field | Temperature | Density | Velocity |
|--------|-----------------------|----------------|-------------|---------|----------|
| $B$    | $B$                   | 0.1 nT         | $1.4 \times 10^5$ K | $0.2 \times 10^{-3}$ cm$^{-3}$ | $150$ km s$^{-1}$ |

**Note.** The velocity in the table corresponds to the “bulk velocity” of the flow measured by V2 near the downstream of the TS.
where $\lambda_{\text{mfp}}$ is the electron mean free path and $\beta_0$ is the ratio of thermal pressure to magnetic pressure (see more discussion in Yamada et al. 2006).

It is clear from the discussion above that the scales of magnetic reconnection dominated by the Hall-MHD term correspond to less than AU (see the appropriate parameters in Table 1). We note that the reconnection of with $\delta_{\text{inj}} < 1$ is called “collisionless.” This may be misleading, as, for instance, a usual definition for being collisionless for magnetized plasma in the interstellar medium (ISM) is to have many Larmor rotations per time between collisions. In this sense the aforementioned requirement for the Hall-MHD reconnection is a requirement to be “super-collisionless,” as the constraint on the number of collisions is a factor $L/d_{\text{inj}} \gg 1$ more stringent than the usually adopted one. Thus, magnetic reconnection in most phases of the ISM (see Table 1 in Draine & Lazarian 1998) is “collisional.”

While the actual scales of relevant magnetic reconnection in the heliosheath and near the heliopause are expected to be large (e.g., $L \sim 100$ AU), it will be clear from the further discussion that this does not actually matter. The Hall-MHD reconnection was demonstrated to provide stable Petschek-type4 configuration characterized by X-points. However, in the presence of external forcing in the heliosheath, for example, the forcing due to variations of solar wind pressure and turbulence, one would expect to observe a collapse of X-points and formation of extended thick outflow regions instead. Incidentally, such regions have been identified by Ciavarella & Raymond (2008) via multi-frequency observations of solar flares. The reconnection within such outflow regions is expected to happen according to a different scheme.

Most astrophysical fluids are turbulent and the heliosheath is not an exception5 with the temporal flows as shown by Richardson et al. (2008). The flows will not only deflect near the heliopause but are expected to be affected near the magnetic ridges that form close to the heliopause (see Opher et al. 2004). In addition, various instabilities are likely to inject turbulent energy. For example, Opher et al. (2003, 2004) suggested the possibility of instabilities near the HCS, with a narrow jet of high-speed flow, strong wrapping of the HCS, and movement away from the ecliptic. The instability has a characteristic wavelength of tens of AU. Thus, the region near the current sheet in the HCS might be unstable and dynamic on the scale of the HCS. Such instabilities could produce high levels of turbulence, back flows and gradients of density and pressure (see also Section 5.2).

What are the characteristics of magnetic reconnection that we can expect in the heliosheath? We argue in Section 5.2 that the characteristics of turbulence may change, especially as a result of flow slowing down and more active reconnection taking place. The characteristic Alfvén velocity $V_A$ in the flow downstream of the TS is $\sim 50$ km s$^{-1}$ if we use the data in Table 1. Using the John Richardson data from the V2 flow instrument, of the flows in the heliosheath (which is presented in Opher et al. 2009), we can say that $v_1 / V_A \sim 1$. The injection scale of the turbulence $\ell$ is uncertain. If the flow gets globally turbulent, the characteristic size of the flow can provide an estimate of $\ell$.

We explain below that for our discussion the most important point for our discussion are the diffusivities and resistivities perpendicular to the magnetic field. Thus we adopt the resistivity coefficient $\eta_L \approx 1.3 \times 10^{13} \ln \Lambda / T^{3/2} \text{cm}^2 \text{s}^{-1}$, where $\ln \Lambda \approx 30$ for the parameters in Table 1, and the viscosity coefficient $\nu_L \approx 1.7 \times 10^{-2} n \ln \Lambda / (T^{1/2} B^2)$ (see Spitzer 1962). The rough estimates of the Lundquist and Reynolds numbers (Re $\equiv v_L / \nu$) are $S \sim 10^{13}$ and $Re \sim 10^{14}$. The exact values of $S$ and $Re$ are irrelevant and the most important message from the above exercise is that the flows are expected to be turbulent (as Re $\gg 1$, see also Section 5.2) and the Sweet–Parker reconnection is expected to be negligible (see above).

Turbulence in the heliosheath is expected to make magnetic fields at least weakly stochastic. A model of fast three-dimensional reconnection that generalizes the Sweet–Parker scheme for the case of a weakly stochastic magnetic field was proposed by LV99. While the notion of turbulence affecting the reconnection rate was not unprecedented (see Speiser 1970; Strauss 1988; Matthaeus & Lamkin 1985, 1986), the LV99 model was the first that predicted the change of reconnection rates as the function of turbulence intensity and the turbulence injection scale. It also proved that in the presence of turbulence three-dimensional magnetic reconnection is fast, that is, independent of $S$, even if only Ohmic resistivities are considered (see also Kowal et al. 2009).

The middle and lower panels of Figure 4 illustrate the key components of the LV99 model.6 The reconnection events happen on small scales $\lambda_1$ where magnetic field lines get into contact. As the number of independent reconnection events taking place simultaneously is $L/\lambda_1 \gg 1$, the resulting reconnection speed is not limited by the speed of individual events on the scale $\lambda_1$. Instead, the constraint on the reconnection speed comes from the thickness of the outflow reconnection region $\Delta$, which is determined by the magnetic field wandering in a turbulent fluid. The model is intrinsically three dimensional7 since both field wandering and simultaneous entry of many independent field patches, as shown in Figure 4, are three-dimensional effects. The magnetic reconnection speed gets comparable with $V_A$ when the scale of magnetic field wandering8 $\Delta$ gets comparable with $L$.

For a quantitative description of the reconnection, one should adopt a model of MHD turbulence (see Iroshnikov 1963; Kraichnan 1965; Dobrowolny et al. 1980; Shebalin et al. 1983; Montgomery & Turner 1981; Higdon 1984). The most important for magnetic field wandering is the Alfvénic component (see a discussion of the decomposition in Section 5.2). Adopting the Golde & Sridhar (1995, henceforth GS95) scaling of the Alfvénic component of MHD turbulence extended to include the case of the weak turbulence (see the Appendix), LV99 predicted that the reconnection speed in a weakly turbulent magnetic field

\footnote{Whether this is really the Petschek reconnection is still a subject of debates. For instance, slow shocks, which are an essential feature of the Petschek model, are not observed in the simulations (see Zweibel & Yamada 2009).}

\footnote{If turbulence is very weak, the injection of energy due to reconnection can speed up the reconnection, resulting in flares (Lazarian & Vishniac 1999, henceforth LV99; Lazarian et al. 2009b; T. Hoang et al. 2009, in preparation). Thus it is plausible that collisionless Hall-MHD effects can enhance the level of turbulence for the situation when magnetic fields are almost laminar initially. This is a rather unlikely scenario for the heliosheath, however (see more discussion in Section 5.2).}

\footnote{Our two-dimensional numerical simulations of turbulent reconnection show that the reconnection is not fast in this case.}

\footnote{Another process that is determined by magnetic field wandering in the diffusion of energetic particles parallel to the mean magnetic field. Indeed, the diffusion coefficients of perpendicular diffusion in the Milky Way is just a factor or order unity less than the coefficient of the diffusion parallel to magnetic field (see Jokipii 1999 and references therein).}
which the Alfvénic turbulence transfers from the weak to strong regimes. Thus Equation (3) can also be rewritten as

\[ V_R = V_A(l/L)^{1/2}(v_l/V_A)^2, \]  

(3)

where the level of turbulence is parameterized by the injection velocity \( v_l < V_A \) and the turbulence injection scale \( l \).

The scaling predictions given by Equation (3) have been tested successfully by three-dimensional MHD numerical simulations in Kowal et al. (2009). This stimulates us to adopt the LV99 model as a starting point for our discussion of magnetic reconnection in the heliosphere.

How can \( \lambda_{||} \) be determined? In the LV99 model as many as \( L^2/\lambda_{\perp}\lambda_{||} \) localized reconnection events take place each of which reconnects the flux at the rate \( V_{\text{rec,local}}/\lambda_{\perp} \), where the velocity of local reconnection events at the scale \( \lambda_{||} \) is \( V_{\text{rec,local}} \). The individual reconnection events contribute to the global reconnection rate, which in three dimensions gets a factor of \( L/\lambda_{||} \) larger, that is,

\[ V_{\text{rec,global}} \approx L/\lambda_{||} V_{\text{rec,local}}. \]  

(4)

The local reconnection speed, conservatively assuming that the local events happen at the Sweet–Parker rate, can be easily obtained by identifying the local resistive region \( \delta_{\text{SP}} \) with \( \lambda_{\perp} \) associated with \( \lambda_{||} \) and using the relations between \( \lambda_{||} \) and \( \lambda_{\perp} \) that follow from the MHD turbulence model (see the Appendix, Equation (A2)). The corresponding calculations in LV99 provided the local reconnection rate \( v_lS^{-1/4} \). Substituting this local reconnection rate in Equation (4), one gets the estimate of the global reconnection speed, if \textit{this speed were limited by Ohmic resistivity}, which is larger than \( V_A \) by a large factor \( S^{1/4} \).

As a result, one has to conclude that the reconnection does not depend on resistivity.

However, it is possible to invert the above arguments and search for the largest scale \( \lambda_{||} \) at which the Sweet–Parker reconnection can reconnect the magnetic field bundles provided the reconnection happens with the speed given by Equation (3). Evidently, the reconnection speed at this scale is \( V_{\text{rec,int}} \approx V_A\delta_{\text{SP,modified}}/\lambda_{||} \), where \( \delta_{\text{SP,modified}} \) is a new Sweet–Parker scale related to the reconnection events at the scale \( \lambda_{||} \). Equation (4) is valid for all the scales and not only for the smallest one (Lazarian et al. 2004). Thus substituting \( V_{\text{rec,int}} = V_AS^{-1/2}(L/\lambda_{||})^{1/2} \) there and combining the result with Equation (3), one gets the fact that the largest scale Sweet–Parker reconnection events take place at the scale

\[ \lambda_{||} = L S^{-1/3} \left( \frac{L}{T} \right)^{1/3} \left( \frac{V_A}{V_l} \right)^{4/3}, \]  

(5)

and the corresponding thickness of the Sweet–Parker current sheet is

\[ \delta_{\text{SP,modified}} = L S^{-2/3} \left( \frac{L}{T} \right)^{1/6} \left( \frac{V_A}{V_l} \right)^{2/3}. \]  

(6)

The scale \( \lambda_{||} \) given by Equation (5) is much smaller than \( L \) and it is evident that for this scale the constraint given by Equation (2) is well satisfied for the heliosphere. Therefore, for all the scales involved in the heliospheric reconnection the \textit{local} reconnection takes place in the regime when plasma can be considered collisionless.

What changes if plasma gets collisionless on scales \( \lambda_{||} \ll L \)? As the local reconnection speed does not limit the speed of reconnection in the LV99 model, one does not expect the change of the rate given by Equation (3).

The LV99 model is applicable to the situations when plasma effects are included, for example, Hall-MHD effect, which increases effective resistivities for local reconnection events. While the latter point is difficult to test directly with the existing plasma codes, for example, with PIC codes, due to the necessity of simulating both plasma microphysics effects as well as macrophysical effects of magnetic turbulence, Kowal et al. (2009) simulated the action of plasma effects by parameterizing them via anomalous resistivities. The values of such resistivities are a steep function of the separation between the oppositely directed magnetic field lines, which also determine the current separating magnetic fluxes. With anomalous resistivities the structure of the fractal current sheet of the turbulent reconnection changed substantially, but no significant changes of the reconnection rate were reported, which agrees well with the theoretical expectations of the LV99 model. Within the model the explanation of this stems from the fact that the reconnection is already fast (i.e., independent of resistivity) even when small-scale reconnection events are mediated by the Ohmic resistivity, while the bottleneck for the reconnection process is provided by magnetic field wandering. Thus, the increase in the local reconnection rate does not increase the global reconnection speed.

As a result, in what follows to describe magnetic reconnection in the heliosphere we adopt the LV99 model but with the reconnection events on the scale \( \lambda_{||} \) happening in a collisionless fashion. The latter may have important consequences for the acceleration of the electrons that we discuss below, but is unlikely to affect heavier species.

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9 As we discuss in the Appendix, the combination \( V_A(v_l/V_A)^2 \) is the velocity of largest strong turbulence eddies \( V_{\text{strong}} \), that is, the velocity at the scale at which the Alfvénic turbulence transfers from the weak to strong regimes. Thus Equation (3) can also be rewritten as \( V_R = V_{\text{strong}}(l/L)^{1/2} \).
The electric field associated with reconnection events can accelerate energetic particles. For a particle of charge \( q \) the typical energy gained in such a process is of the order of
\[
q(V_R/c)B\lambda_q,
\]
where \( \lambda_q \) is the coherence length of a particle within the reconnection layer. To accelerate this way, one requires to have both \( V_R \) and \( \lambda_q \) to be large. Reconnection as a process of accelerating energetic particles via the electric field within the Sweet–Parker reconnection model (see Haswell et al. 1992; Litvinenko 1996) is inefficient due to slow reconnection rates. Moreover, due to the tiny speed of reconnection only a very small fraction of the magnetic energy can be potentially used for driving the acceleration.

Electric fields are much stronger in the Petschek (1964) reconnection. However, in the Petschek model \( \lambda_q \) gets small, which does not allow efficient acceleration either. We may state that in general in any fast reconnection scheme whether it is the Petschek, collisionless, LV99, or any other, the fraction of the volume that is being subject to resistive effects and reveals strong electric fields is small and most of the magnetic energy is converted into kinetic energy. Thus, only a small fraction of energy can be transferred through any fast reconnection process to energetic particles if direct electric field acceleration is involved. Therefore, we shall ignore this process dealing with ACRs in the heliosphere.

In what follows, for the sake of simplicity, we consider the acceleration of ACRs in an individual reconnection layer. The situations of multiple reconnection layers may present additional effects arising from particles interacting with several layers. However, these situations are not considered in this paper.

In the LV99 scheme, the reconnection velocities can approach \( V_A \) and therefore be appreciable. Thus one visualizes a different scheme of reconnection, which is similar to the one involving shocks. Indeed we showed that, due to high speeds of stochastic reconnection, particles entrained on reconnecting field lines bounce back and forth between magnetic walls while staying on the field lines that contract. This results in a systematic increase of the velocity with every bouncing of energetic particles. Such a model was discussed in de Gouveia Dal Pino & Lazarian (2003, henceforth GL03, 2005; see also Lazarian 2005), where it was showed that the reconnection induces first-order Fermi acceleration of the particles entrained. Figure 5 shows the cross sections of the numerical \( 256^2 \times 512 \) box with the magnetic field perturbed sub-Alfvénically, that is, with the turbulent injection velocity \( v_I < V_A \). The field loops of the reconnected magnetic flux are clearly visible. Energetic particles are being accelerated as the magnetic field lines of the three-dimensional loops shrink and the particles bounce back and forth between converging magnetic fluxes.

Figure 6 exemplifies the simplest realization of the acceleration within the reconnection region expected within the LV99 model. As a particle bounces back and forth between converging magnetic fluxes, it gains energy.

The process of acceleration depicted in Figure 6 is easy to quantify. If an energetic particle bounces back and forth between the magnetic mirrors created by magnetic flux, such a particle having energy \( E \) will in every collision gain energy
\[
\sim V_R/cE.
\]

The process would continue till the particle either diffuses perpendicular to the reconnection flux or gets ejected by the outflow together with the plasma and reconnected magnetic flux. The latter possibility was considered in GL03 and de Gouveia Dal Pino & Lazarian (2005). Formally, the physical setup considered there corresponds, first of all, to the assumption of the \( y \)-component of the diffusion velocity of cosmic rays:

\[
V_{y,\text{diff}} = \kappa_{yy}/\Delta,
\]

where \( \Delta \) is the thickness of the outflow reconnection region.
and \( \kappa_{yy} \) is the diffusion coefficient perpendicular to the average magnetic field, it being slow compared to \( V_R \). This assumption implies that the scattering of energetic particles is efficient and the diffusion coefficient \( \kappa_{yy} \) is small. In addition, it assumes that the diffusion of energetic particles parallel to the magnetic field

\[
V_{x,\text{diff}} = \kappa_{xx}/L, \tag{8}
\]

where \( \kappa_{xx} \) is a coefficient of parallel diffusion, does not exceed \( V_A \).

The diffusion coefficient \( \kappa_{yy} \) arises from magnetic field wandering as particles diffuse marginally perpendicular to the local magnetic field. The theory-motivated rates of magnetic field wandering were calculated in LV99 and later used in Lazarian (2006) to understand heat diffusion and in Yan & Lazarian (2008) to describe cosmic ray propagation. One may argue that the measurements of diffusion parallel and perpendicular to the mean magnetic field are ill-motivated, as the energetic particles do not feel any other field apart from the local field they interact with.

We believe that all the quantities of the future scattering and acceleration theories should be formulated in terms of local fields, the same way as the turbulent theories are formulated (see the Appendix). If, nevertheless, one adopts a conventional approach, for the developed strong turbulence at scales \( \Delta \) less than the turbulence injection scale \( l \) the perpendicular diffusion coefficient for energetic particles is \( \kappa_{yy} \approx (\Delta/L)(v_t/V_A)\kappa_{xx} \) (Yan & Lazarian 2008). Substituting the latter expression in Equation (7) and comparing \( V_{y,\text{diff}} \) with \( V_R \) given by Equation (3), one may note that for sub-Alfvénic turbulence with \( l \sim L \) this constraint is less restrictive than the constraint arising from \( V_{x,\text{diff}} < V_A \). For the sake of simplicity, in what follows we shall refer to \( V_{\text{diff}} \) as the most restrictive of the two processes.

If one uses the experimentally measured parallel diffusion coefficient suggested by the Palmer consensus (Palmer 1982), namely the diffusion coefficient corresponding to the parallel mean free path for particles between 0.5 MeV and 5 GeV being from 0.08 AU to 0.3 AU, \( V_{\text{diff}} \) given gets larger than the corresponding velocity of advection of cosmic rays with the reconnected magnetic field. However, for the diffusion along magnetic field lines undergoing reconnection the aforementioned estimate for the diffusion speed is definitely an overestimate, as both magnetic reconnection and streaming particles should create magnetic perturbations that efficiently scatter particles back (see Cesarsky 1980). The former effect of modification of turbulence in turbulent reconnection was predicted in LV99 and observed in simulations by Kowal et al. (2009).10 Thus, we expect the particles to be both reflected back by magnetic bottles and efficiently scattered by the modified turbulence.

To derive the energy spectrum of particles, one can use the routine way of dealing with first-order Fermi acceleration (see Longair 1992). Consider the process of acceleration of \( M_0 \) particles with the initial energy \( E_0 \). If a particle gets energy \( \beta E_0 \) after a collision, its energy after \( m \) collisions is \( \beta^m E_0 \). At the same time if the probability of a particle to remain within the accelerating region is \( P \), after \( m \) collisions the number of particles gets \( P^m M_0 \). Thus, \( \ln(M/M_0)/\ln(E/E_0) = \ln P/\ln \beta \) and

\[
M/M_0 = \left( \frac{E}{E_0} \right)^{\ln P/\ln \beta}. \tag{9}
\]

For the stationary state of accelerated particles the number \( M \) is the number of particles having energy equal to or larger than \( E \), as some of these particles are not lost and are accelerated further. Therefore,

\[
N(E)dE = \text{const} \times E^{-1+4(\ln P/\ln \beta)}dE. \tag{10}
\]

To determine \( P \) and \( \beta \), consider the following process. The particles from the upper reconnection region see the lower reconnection region moving toward them with the velocity \( 2V_R \) (see Figure 6). If a particle from the upper region enters at an angle \( \theta \) into the lower region, the expected energy gain of the particle is \( \delta E/E = 2V_R \cos \theta/c \). For isotropic distribution of particles their probability function is \( p(\theta) = 2\sin \theta \cos \theta d\theta \) and therefore the average energy gain per crossing of the reconnection region is

\[
(\delta E/E) = \frac{V_R}{c} \int_0^{\pi/2} 2\cos^2 \theta \sin \theta d\theta = 4/3 \frac{V_R}{c}. \tag{11}
\]

An acceleration cycle is when the particles return back to the upper reconnection region. Being in the lower reconnection region, the particles see the upper reconnection region moving the speed \( V_R \). As a result, the reconnection cycle provides the energy increase \( (\delta E/E)_{\text{cycle}} = 8/3(V_R/c) \) and

\[
\beta = E/E_0 = 1 + 8/3(V_R/c). \tag{12}
\]

Consider the case of \( V_{\text{diff}} \ll V_R \). The total number of particles crossing the boundaries of the upper and lower fluxes is \( 2 \times 1/(4n) \), where \( n \) is the number density of particles. With our assumption that the particles are advected out of the reconnection region with the magnetized plasma outflow, the loss of the energetic particles is \( 2 \times V_R n \). Therefore, the fraction of energetic particles lost in a cycle is \( V_R n/[1/4(n)c] = 4V_R/c \) and

\[
P = 1 - 4V_R/c. \tag{13}
\]
Combining Equations (10), (12) and (13), one gets

\[ N(E)dE = \text{const}_1 E^{-5/2}dE, \]  

(14)

which is the spectrum of accelerated energetic particles for the case when the backreaction is negligible (see also GL03).\(^{11}\)

We note that Equation (14) provides the estimate of the expected spectrum for a rather idealized situation. It is important to understand what modifications of the spectrum we expect to observe under more realistic circumstances. First of all, the above derivation considers only particles bouncing between lower and upper reconnecting fluxes. The actual picture of the stochastic reconnection in LV99 includes many reconnection events happening at different scales (see Figure 5). However, each of these events can be viewed as a repetition of the large-scale reconnection event and should provide the same type of power spectrum. In fact, the efficiency of acceleration increases with the decrease of the scale of reconnection (see Section 4.3).

Dealing with the reconnection, we have another limiting case which can be easily treated. In the case when \( V_{\text{y, diff}} \ll V_{\text{x, diff}} \) and \( V_{\gamma, \text{diff}} \gg V_{\phi} \) one can formally take \( P \approx 1 \), as the ejection of the particles is negligible. This situation corresponds to an infinite reconnection region where particles may diffuse into the magnetic flux, but do not escape. In this situation, Equation (10) provides

\[ N(E)dE = \text{const}_2 E^{-1}dE, \]  

(15)

which coincides with the spectrum obtained by J. R. Jokipii (2009, private communication) by solving the one-dimensional Parker equation for the Sweet–Parker reconnection process. Naturally, the model of particle acceleration which does not allow for particle escape is rather artificial. However, our discussion of this situation exemplifies the fact that the decrease of the particle escape makes the spectrum of accelerated particles more shallow. Similarly, one can show that the diffusion along magnetic field lines in the direction, which enhances particle escape, should make the spectrum of the energetic particles steeper.

4.2. Particle Backreaction

The backreaction of accelerated particles is known to be important for the processes of shock acceleration, although the consensus on the quantitative description of the process is still missing (see Malkov & Diamond 2009). Thus, it is not surprising that the situation with the backreaction of particles in the reconnection sites is rather unclear. The only work attempting to address this problem that we are aware of is that by Drake et al. (2006, henceforth DX06). DX06 repeated the claim in GL03 that the first-order Fermi acceleration should happen as a result of interaction of energetic particles with reconnecting magnetic fields. However, the study is intended for the acceleration of electrons and the model is of two-dimensional collisionless reconnection. In this scheme of reconnection, the production of contracting magnetic loops is expected. Within these loops, the energetic particles are expected to undergo acceleration.

The acceleration of particles other than electrons was not discussed by DX06 within their model. For the small-scale loops considered by DX06, the Hall term dynamics beyond the ion skin depth is important. In both cases, the contracting magnetic loops increase the energy of entrained particles. In fact, the evidence for such a process can be seen in the simulations of test particles in the magnetotail (Birn et al. 2004) and traced back even further to the test particles studies in MHD models with magnetic islands (Matthaeus et al. 1984; Kliem 1994).

DX06 appealed to physical arguments rather than to direct numerical calculations to justify the advocated picture of backreaction. In particular, the backreaction is introduced by the term \( \left(1–8\pi\bar{\epsilon}_1/B^2\right) \), where \( \bar{\epsilon}_1 \) is the parallel energy of energetic particles averaged over the distribution of particle velocities. This term gets negative to represent the halt in the contraction of the magnetic loops. The generalization of this picture for open loops expected for a generic three-dimensional reconnection is feasible, but was not performed by DX06.

We feel that the DX06 model, although formulated in terms of contracting two-dimensional loops, can potentially be generalized for contracting three-dimensional spirals, in which the acceleration processes in GL03 and DX06 are similar.\(^{12}\) The notable difference is that with the Hall-MHD effect included the loops can preferentially accelerate electrons on the scales less than the ion skin depth \( d_{\text{ion}} \).

In terms of the nature of the particle backreaction, electrons and protons should act similarly as soon as their energy gets comparable with the energy of the magnetic field that drives them. This, according to DX06, may result in a more shallow spectrum, with the index \(-3/2\) rather than \(-5/2\) predicted by Equation (14).

It is not clear to what extent the backreaction is important for the energetic particles accelerated from low energy (around keVs) into the MeVs ACRs (from 0.04 to 4 MeV) in the zones of reconnection in the heliosheath. Thus, depending on the importance of the yet unclear backreaction of the energetic particles, the actual spectral index may vary from \(-5/2\) to \(-3/2\), which encompasses the value of \(-5/3\) that was observed by the Voyagers. The fact that the observed value is close to the low boundary of the expected indices means that the backreaction of ACRs to magnetic loops contracting in the process of reconnection is probably important.

Contracting loops are not a part of the picture of reconnection in the limits \( V_{\gamma, \text{diff}} > V_{\phi} \) or and \( V_{\gamma, \text{diff}} > V_{\phi} \). For these cases, the backreaction of the particles may be very different. We expect both parallel diffusion to increase and the perpendicular diffusion to decrease with the increase of the strength of the shared component of the magnetic field. We also expect that for higher energy particles the confinement within the reconnection region gets problematic, which limits the energies to which the energetic particles can be accelerated. For instance, the upper boundary on the energies of the accelerated energetic particles comes from the requirement that the particle Larmor radius \( R_l \) is less than the transverse dimension of the magnetic flux undergoing reconnection. Assuming that the latter is about 1 AU.

\(^{11}\) The obtained spectral index is similar to that of Galactic cosmic rays.

\(^{12}\) We feel that the potential deficiency of considering closed loops compared to the open loops in the acceleration picture advocated in this paper is that in the incompressible limit and in the presence of particle scattering we expect to see the cancellation of the increase of the total particle momentum, the increase of the particle momentum \( p_t \) through magnetic loop contraction considered by DX06, and the decrease of the momentum perpendicular to the magnetic field \( p_t \) through the betatron effect, as was shown for incompressible MHD turbulence in Cho & Lazarian (2006). In the absence of the collisions, the increase of the parallel momentum for a contracting loop is limited by the relation on the particle energy \( E \) and the loop length \( L_{\text{loop}} \) of the form \( E f_{\text{loop}} = \text{const} \). It is easy to see that the escape of particles from one loop and capturing into another should result in random changes in the projections of particle momentum and the effect of this loop changing is similar to collisions, which makes the limitations on acceleration derived in Cho & Lazarian (2006) again applicable. These limitations are not applicable to the case of the open loops, whose ends are connected with converging fluxes that we discussed above (see also Somov & Kosugi 1997; Giuliani et al. 2005).
we get the maximal energy of particles the magnetic field of 0.1 nT (see Table 1) of the order of 10^3 MeV, which is higher than the energies of the ACRs. However, as the particles approach this energy their diffusion and escape in the x-direction increases, changing the nature of acceleration and decreasing its efficiency. Needless to say that these and similar issues require a further quantitative investigation.

4.3. Maximal Rates of Acceleration

The acceleration efficiencies for ACRs are estimated to be high with the acceleration time being less than a year for a roughly 100 MeV ACR (Mewaldt et al. 1996). This induces constraints on the possible acceleration mechanism responsible for the origin of these particles. Appealing to an analogy of the first-order Fermi acceleration at reconnection sites and magnetized shocks, one can use the relevant results in Jokipii (1992).

Indeed, the most efficient acceleration of ACRs is expected when they experience bouncing of the converging mirrors every Larmor period. Such mirrors are expected to be created by small-scale reconnection events arising from oppositely moving flux tubes within the reconnection outflow region. Such events are predicted in LV99 and can be seen in the visualizations of the reconnection regions in the presence of turbulence (see Figure 5). The particle within these microscale reconnection regions behaves analogous to a particle accelerated in a perpendicular shock considered in Jokipii (1992). For a particle of the Larmor radius \( R_L \), one can roughly estimate the corresponding acceleration time as \( t_{accel} > 2 R_L / V_A \). The latter estimate is a factor of unity different from the estimate obtained by Jokipii (1992) for a perpendicular shock, where \( 8 R_L / V_{shock} \) was obtained. Thus the estimates of a minimal acceleration time, which is less than a month in Jokipii (1992), are applicable to the process we consider in the paper.

5. DISCUSSION AND SUMMARY

This paper is of exploratory nature. We propose an alternative mechanism for explaining the origin of ACRs appealing to fast magnetic reconnection. Some time ago, such an appeal would sound completely speculative due to the notoriously enigmatic nature of magnetic reconnection. However, as a result of recent progress in understanding of magnetic reconnection, we believe that the time is right to start discussing the consequences of the process.

5.1. Expected Sites of Acceleration

Current sheets are common in the heliosphere. However, not all current sheets are expected to be associated with particle acceleration. For tangible particle acceleration, the reconnection rates should be high, that is, comparable with \( V_A \). In this respect, LV99 model predicts that the reconnection rate may be low or high depending on the level of turbulence (see Equation (3)).

When the reconnection velocity is a small fraction of the solar wind speed (e.g., as this is the case prior to the TS shown in Figure 2), then the reconnection is expected to play a marginal role in the particle acceleration. However, as the flow slows down after the TS, magnetic field reversals come closer and get crowded, which is illustrated in Figure 2 by current sheets crowding after the TS. As this happens, we expect a larger portion of the magnetic flux to be consumed by reconnection per unit time and a larger density of the accelerated energetic particles.

In addition, the large-scale variations of the magnetic field associated with solar cycles are expected to contribute to the acceleration of energetic particles. It is clear from Figure 3 that the corresponding effect is negligible within the TS. In fact, “touching” X-point Petschek-type events are expected before the TS (see Gosling et al. 2005a). However, it is also clear that magnetic fields associated with different cycles are pressed together as we approach the heliopause. This should shift the sites of the energetic particle acceleration toward the heliopause.

We should mention that the direct in situ studies of reconnection (see Gosling et al. 2005a, 2007; Phan et al. 2009) have not revealed the excess of the energetic particles associated with magnetic reconnection exhausts in the solar wind (Gosling et al. 2005b). We believe that this can be due to the fact that the X-point Petschek-type reconnection, signatures of which were searched for, is inefficient in accelerating energetic particles as we discussed in Section 4. In the paper, we appeal to thick turbulent reconnection regions (see Figure 4) as the sites of acceleration.14 We expect such regions to emerge as magnetic field reversals get crowded and magnetic field lines of different direction press against each other making the X-type opening of reconnection regions problematic (see the discussion in Section 2).

Finally, let us mention the current sheets associated with magnetic turbulence in the absence of large-scale magnetic field reversals. The variations of magnetic field directions induced by turbulent motions induce their own small-scale current sheets. However, for sub-Alfvénic driving the variations of the magnetic field induced by turbulence are not large. Only the oppositely directed components of the magnetic field are involved in the reconnection process and that it is they that determine the Alfvén speed that enters Equation (3). Thus, the large-scale reversals are much more preferable sites for energetic particle acceleration. However, the first-order Fermi acceleration within reconnection layers created by MHD turbulence may not always be neglected, especially when compared to the second-order Fermi acceleration by turbulence (see Section 5.5).

5.2. Parameters of Turbulence

According to Equation (3) the rate of reconnection depends on the level of turbulence, that is, on \( \nu_i \), and the injection scale of the turbulence \( l \). The two parameters are still uncertain in depth of the heliosheath toward the heliopause, where we expect most of the magnetic reconnection events to take place. However, even very weak turbulent driving corresponding to a small fraction of \( V_A \) provides substantial reconnection rates incomparably larger than the Sweet–Parker prediction (see Equation (3)). The dependence on \( l \) is rather weak in Equation (3) and therefore an estimate \( l \sim L_s \) should be acceptable.

Why are we sure of the presence of turbulence at the distances which have not been probed by Voyagers so far? One may claim on very general grounds that astrophysical turbulence is the consequence of high Reynolds numbers of astrophysical flows (see a discussion in Lazarian et al. 2009a). Generically, hydrodynamic flows get turbulent for \( \text{Re} \sim 10 \) or 100.

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13 An additional dependence comes from the boundary conditions. If plasma and shared magnetic flux, that is, the magnetic flux associated with the guide field, are constrained from freely leaving the reconnection zone, the reconnection rates decrease compared to those given by Equation (3).

14 Alternatively, in the two-dimensional picture employed in DX06, the loops within the thick reconnection regions rather than X-points are the engines of the accelerated particles.
A notable exception of this rule are Keplerian flows in accretion disks, which, however, also get turbulent for large Re in the presence of weak magnetic fields (Velikhov 1959; Chandrasekhar 1960; Balbus & Hawley 1991). For the magnetized flows in the heliosheath, the damping of Alfvénic perturbations is limited by the perpendicular viscosity discussed in Section 3. Thus the expected Re will be larger than $10^{10}$, which much exceeds the threshold for the fluid getting turbulent.

Recent numerical simulations support theoretical conclusions that magnetic reconnection can be a self-regulating process and the initial level of magnetic stochasticity may matter little as the reconnection proceeds. For instance, three-dimensional MHD simulations in T. Hoang et al. (2009, in preparation) showed that magnetic fluxes of different directions in low-$\beta$ plasma brought in the contact at $t = 0$ develop turbulence and dissipate the flux within several crossing Alfvénic times (see also Bettarini & Lapenta 2009).

It is very advantageous that the reconnection in LV99 depends only on the Alfvénic component of MHD turbulence, as the compressible components are subject to more damping (see Brunetti & Lazarian 2007) and also more fancy channels of cascading (see Chandran 2005). The possibility of segregating of the Alfvénic component from the rest of the MHD cascade is suggested theoretically in GS95 and proven in direct three-dimensional compressible MHD numerical simulations (Cho & Lazarian 2002, 2003). In the fully ionized plasma of the heliosheath, the Alfvénic component cascades to the proton Larmor radius (see the discussion of the cascade in a partially ionized gas in Lazarian et al. 2004).

We discuss in the Appendix that the possible uncertainties of the spectral index of the Alfvén turbulence that are consistent with the observations, in situ spacecraft measurements, and numerical simulations do not change the LV99 conclusion that the reconnection is fast. Thus, the model suggested in the paper does not depend on the outcome of the ongoing debates on the exact spectral index of Alfvénic turbulence. Similarly, in the presence of the backreaction of accelerated particles the turbulence spectrum is expected to be modified, but the rates of reconnection and therefore rates of ACR accelerations are not expected to change appreciably.

5.3. Predictions and Limitations of the Model

In the above in this paper, we proposed that the acceleration of particles arising from magnetic reconnection could explain the origin of the ACRs. In particular we predict that the source of the ACRs can be deeper in the heliosheath, close to the heliopause. We maintain that the acceleration of energetic particles by magnetic reconnection is an unavoidable process. Our idea can be tested as the ACR energy spectra of H, He, N, Ne, and O in the heliosheath will slowly unroll as V1 and V2 make their way into the heliosheath.

The actual predictions of V1 and V2 measurements require detailed numerical modeling of the energetic particle propagation in the heliosheath, which is a problem far from being handled reliably. Qualitatively speaking, we expect to see the variations of the anisotropy of the ACRs as Voyagers approach and cross reconnection regions. The turbulent reconnection regions are characterized by substantial changes in the direction of the magnetic field, presence of magnetic flux at substantial angles to the magnetic field in the adjacent regions, and magnetic field reversals of this unusually aligned flux (see Figure 5). We also expect to see an increased level of velocity fluctuations as the reconnection itself drives turbulence. At the same time, the events at scales less than $\lambda_1$ given by Equation (5) can show signatures of collisionless reconnection. Detailed predictions concerning ACRs will become available as the theory of acceleration in turbulent reconnection region matures (see the discussion of problems below).

Difficulties associated with the understanding of astrophysical acceleration processes are well known. Even for the well-recognized process of particle acceleration in shocks (Axford et al. 1977; Krymsky 1977; Bell 1978; Blandford & Ostriker 1978), the details of acceleration are unclear (see Malkov & Diamond 2009). It is not surprising that for the first-order Fermi acceleration within a turbulent reconnection zone, we grope for basic facts. Obviously, the model we presented in Section 4.1 is simplistic. It does not prescribe the energetic particle diffusion coefficients, does not account for the compressibility of the medium, disregards the possibility of the interaction between the adjacent reconnection sites, etc. Thus, it is natural that the uncertainties of the expected particle spectrum that we get are much higher and we have to be satisfied with rough consistency of the theoretical expectations to the observed spectra of ACRs. Below, we discuss the potential uncertainties of the models. We expect that future numerical simulations will resolve the outstanding issues that we outline.

The efficiency of particle acceleration in fast reconnection has not been resolved yet. It is clear that the competition will be between the channeling of magnetic energy into the kinetic energy of thermal plasma and energetic particles. The solar flare reconnection events can provide us with a hint. Indeed, the observations solar mass ejections indicate that most of the energy goes into energetic particles. Similarly, the analogy between the shock acceleration and the first-order Fermi acceleration in the reconnection regions is also suggestive of higher percentage of energy going into the acceleration of energetic particles. Future research should clarify this issue. The self-consistent model of the generation of the ACRs in the heliosheath should include the reconnection layers acting as a source term for solving the transport equation for energetic particles (see Schlickeiser 2003). Unfortunately, with the complex and yet unknown structure of the heliosheath, it may be currently impossible to reliably model the energetic particle injection and propagation.

The fact that magnetic reconnection may be strongly affected by the backreaction of the energetic particles does not necessarily ensure that the magnetic field and energetic particles get into equipartition throughout the entire heliosheath region. Indeed, particles may escape efficiently with the reconnected flux and the parts of heliosheath not subject to reconnection may have the energy density of particles well below the equipartition. At the same time, the accumulation of particles above the equipartition value is prohibited due to insufficient magnetic tension for constraining them. If the energy of energetic particles gets locally larger than the magnetic energy one expects to observe local expanding magnetic bubbles, which will burst creating open magnetic field lines (similar to the open magnetic field lines of the solar wind) unless constrained by the pressure of the ambient regions.

In our calculations in Section 4, we assumed that energetic particles entering the unconnected flux (e.g., upper flux tube in Figure 6) get scattered and randomized before they leave it to collide with the oppositely moving magnetic flux (e.g., lower flux tube in Figure 6). This may be not the only process involved. For instance, energetic particles can be reflected by magnetic mirrors, which will not change the adiabatic invariant of the particle. In this case, only the parallel component of
the particle momentum will grow. However, as was argued in Lazarian & Beresnyak (2006) such a distribution is unstable with respect to gyroresonance instability. This instability is known to generate circularly polarized Alfvénic perturbations with $k \parallel B$ and these will scatter and randomize particles, thus decreasing the anisotropy. Therefore, we do not expect a change of the acceleration efficiency.

At the same time, the damping of turbulence at the Larmor radius scale of protons makes the acceleration of minor ions preferable. Indeed, scattering of energetic particles is an important component of the acceleration process. This scattering depends on the presence of magnetic turbulence at the Larmor radius scales. As this scale is larger for the minority ions, their acceleration may proceed more efficiently, which is also the case for the acceleration in shocks (see Mewaldt et al. 1996).

### 5.4. Role of Reconnection Microphysics

Let us show how the assumed microphysics of reconnection affects the energetic particle acceleration. As we discussed in Section 3, the adopted LV99 model of magnetic reconnection in the heliosheath includes a small-scale event mediated by collisionless effects and large-scale reconnection limited by magnetic field wandering according to the LV99 model. While in Section 3 we argue that collisionless effects are irrelevant in terms of overall rates of reconnection, their presence may be important for the acceleration of electrons. One might also expect that the dynamics of the magnetic loops at scales smaller than the proton gyroradius may be different from the magnetic loops on larger scales as a result of the difference in the properties of whistler (electron MHD) and Alfvénic turbulence (Cho & Lazarian 2004, 2009).

For our discussion of the acceleration in a reconnection layer, we adopted the collisionless model of reconnection on the small scales $\lambda_1$. A number of issues related to the reconnection in this regime are still a subject of debate. For instance, the role of the Hall term is challenged in a number of papers (see Karimabadi et al. 2004; Bessho & Bhattacharjee 2005; Daughton & Karimabadi 2007). Whether the collisionless reconnection is fast has also been questioned (see Wang et al. 2001; Fitzpatrick 2004; Smith et al. 2004). In addition, small-scale events may proceed at different rates due to the external forcing preventing the formation of small-scale X-points within a turbulent reconnection region.

We believe that while some details of the acceleration may change, for example, the acceleration of electrons may be modified in the absence of the Hall term, the major elements of the model of the acceleration of energetic particles in the heliosheath, that we discuss in the paper will stay. First of all, even if collisionless reconnection is not fast, this should not change the overall speed of the reconnection of the weakly turbulent magnetic field (LV99). Then, the contracting open magnetic loops accelerating electrons are still likely to emerge at scales smaller than the proton Larmor radius due to the turbulent motions protruding at the electron MHD scale. Future numerical calculations should clarify these issues.

### 5.5. Role of Turbulent Acceleration of ACRs

While the alternative idea of blunt shocks stays a viable one, we have doubts about the idea of second-order Fermi acceleration of energetic particles observed by Voyagers. First of all, the estimates in Jokipii (1992) suggest that the time for the second-order Fermi acceleration by turbulence is prohibitively long. Then, below we claim that the interacting large-scale turbulent eddies can produce reconnection regions inducing the first-order Fermi acceleration, which sometimes may be more efficient than the traditional second-order Fermi acceleration.

It is well known that apart from accelerating the energetic particles in the heliospheric current sheets, magnetic reconnection can induce turbulence that can accelerate energetic particles (see Petrovian et al. 2006 and references therein). Indeed, in any scheme of fast reconnection only a small fraction of energy is consumed through direct plasma heating. While further research is necessary at this point, it is reasonable to assume that the partition of energy that is released directly in the reconnection zone and is available in the form of magnetic tension after magnetic field lines leave the reconnection zone depends on the initial configuration of the reconnecting magnetic fluxes. If most of the energy is accumulated in the form of magnetic field lines bended outside the reconnection region, as this is the case in some models of solar flares (see Tsuneta 1996), then turbulence may absorb most energy released in the event. Although we do not believe that this is the case for the magnetic fields in the heliosheath, even in this unlikely hypothetical situation we expect to see bursts of turbulent reconnection accelerating energetic particles.

Interestingly enough, we may argue that the interaction of energetic particles with reconnection regions naturally arising in magnetic turbulence may result in the first-order Fermi acceleration. It is easy to see that reconnection events in MHD turbulence should happen through every eddy turnover (see LV99). For small scales magnetic field lines are nearly parallel (see the Appendix) and, when they intersect, the pressure gradient is not $V^2_A/\lambda_1$ but rather $(\lambda_1^2/\lambda_\perp^2)V^2_A$, since only the energy of the component of the magnetic field that is not shared is available to drive the outflow. On the other hand, the characteristic length contraction of a given field line due to reconnection between adjacent eddies is $\lambda_\perp^2/\lambda_1$. This gives an effective ejection rate of $V_A/\lambda_1$. Since the width of the diffusion layer over the length $\lambda_1$ is $\lambda_\perp$, the Equation (3) should be replaced by $V_R \approx V_A(\lambda_\perp^2/\lambda_1)$, which provides the reconnection rate $V_A/\lambda_1$, which is just the nonlinear cascade rate on the scale $\lambda_3$. This ensures self-consistency of the critical balance for strong Alfvénic turbulence in highly conducting fluids (LV99). Indeed, if not for fast turbulent reconnection the buildup of unresolved magnetic knots is unavoidable, flattening the turbulence spectrum compared to the theoretical predictions. The latter contradicts both to solar wind measurements and to numerical calculations.

The energy of reconnect magnetic field during the eddy turnover is comparable with the energy of the eddy. In the absence of cosmic ray acceleration the energy liberated in reconnection goes into motions comparable to the dimensions of the reconnecting eddies, so this energy release will not short circuit the turbulent energy cascade. In the presence of cosmic ray acceleration a substantial part of the turbulent energy may go into the energetic particle acceleration and therefore the second-order Fermi acceleration in turbulence may play

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15 It is interesting to note that some of the in situ measurements of the parameters of the current sheet may be consistent with the predictions of the collisionless reconnection, while the actual reconnection goes according the LV99 scheme.

16 Jokipii (1999) pointed out that the coexistence of two big power laws in the sky, the one by cosmic rays and one by turbulence, shows the fundamental inter-relation of turbulence and cosmic rays. One may speculate that the process of first-order Fermi acceleration in the turbulent eddies may be at the core of this inter-relation.
a subdominant role even for the part of the magnetic energy that was released beyond the reconnection region and induced magnetic turbulence.

Consider the problem from the point of view of energetics. For driving with turbulent injection velocity of the order of Alfvén velocity, around 80% of energy go into Alfvénic incompressible modes (Cho & Lazarian 2003). This result obtained with solenoidal driving of turbulence corresponds well to the results for hydrodynamic turbulence where, for arbitrary driving, a substantial part of energy goes into solenoidal motions (Biskamp 2003; see also Federrath et al. 2009 for the MHD case). At the same time, Alfvénic modes arising from large scale driving are shown to be very inefficient for accelerating energetic particles due to the high anisotropy of Alfvénic modes (see Cho & Lazarian 2006). The acceleration by fast modes (Yan & Lazarian 2002, 2004, 2008; Brunetti & Lazarian 2007) is limited\(^\text{17}\) by fast modes’ collisionless damping. This can be viewed as an additional argument against the second-order Fermi acceleration of ACRs.

All in all, while turbulence is essential for driving magnetic reconnection, the role of turbulence in the second-order Fermi acceleration of ACRs may be subdominant. We argue that turbulence itself may be associated with first-order Fermi acceleration related to the local reconnection of the magnetic field of adjacent eddies. If proven by further research, this may substantially increase the potential of turbulence in accelerating energetic particles in the situations beyond the one we considered above.

5.6. Astrophysical Implications of the Model

The first-order Fermi acceleration by the reconnection in a turbulent magnetized fluid may have important astrophysical consequences beyond the space physics. For instance, ISM observations indicate the existence of the power law from dozens of parsecs to sub-AU scales (see Crovisier & Dickey 1983; O’Dell & Castaneda 1987; Green 1993; Armstrong et al. 1995; Elmegreen & Scalo 2004; McKee & Ostriker 2007; Lazarian 2009). Thus, it is not inconceivable that magnetic reconnection can play a significant role in the acceleration of cosmic rays on the galactic scale and under other circumstances. We defer a discussion of this interesting possibility for future papers.

Now that the plasma and field data are both available from the Voyagers at the heliosheath, reconnection sites can be directly probed. If it is confirmed that energetic particles observed by Voyagers are accelerated by reconnection, this work could induce further efforts in identifying situations where first-order Fermi acceleration arising from the reconnection of the weakly stochastic magnetic field is important. The natural place to look is solar flares. While second-order Fermi acceleration is frequently involved to explain energetic particles arising during flares (see La Rosa et al. 2006; Petsionis et al. 2006; Yan et al. 2008), in view of fast reconnection it looks promising that the acceleration is driven by the mechanism we discuss in this paper. The origin of fast ions in solar wind (see Fisk & Gloeckler 2006), relativistic electrons in the galaxy clusters (see Fusco-Femiano et al. 2004; Rephaeli et al. 2006; Brunetti & Lazarian 2007) may be also related to the acceleration within reconnection regions.

We also note that magnetic field reversals and reconnection are an intrinsic part of the magnetic field dynamics in accretion disks. For some of these disks, for example, circumstellar disks, the issues of ionization are important (see Shu et al. 2007). If energetic particles are accelerated in accretion disks, this make magnetic activities there self-sustained. Further research should quantify this and related issues. The theoretical calculations of the reconnection rate in a weakly turbulent partially ionized gas are provided in Lazarian et al. (2004).

5.7. Summary

We proposed that the magnetic reconnection could accelerate energetic particles in the heliosheath and especially near the heliopause. This can explain the fact that Voyagers failed to detect the signatures of shock acceleration. Our predictions include the localization of the source of the energetic particles close to the heliopause.

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APPENDIX

MODEL OF MHD TURBULENCE ADOPTED

The nature of the Alfvénic cascade is expressed through the critical balance condition in the GS95 model of strong turbulence, namely

\[ \lambda_+^{-1} V_A \sim \lambda_\perp^{-1} v_\perp, \quad (A1) \]

where \( v_\perp \) is the eddy velocity, while \( \lambda_\parallel \) and \( \lambda_\perp \) are, respectively, eddy scales parallel and perpendicular to the local direction of the magnetic field. The critical balance condition states that the parallel size of an eddy is determined by the distance Alfvénic perturbation can propagate during the eddy turnover. The notion of local is important,\(^{18}\) as no universal relations exist if eddies are treated with respect to the global mean magnetic field (LV99; Cho & Vishniac 2000; Maron & Goldreich 2001; Lithwick & Goldreich 2001; Cho et al. 2002). Combining this with the Kolmogorov cascade notion, that is, that the energy transfer rate is \( v_\perp^2/(\lambda_\perp/v_\perp) = \text{const} \), one gets \( \lambda_\parallel \sim \lambda_\perp^{2/3} \). If the turbulence injection scale is \( L_{\text{inj}} \), then \( \lambda_\parallel \sim \lambda_\perp^{2/3} \), which shows that the eddies get very much anisotropic for small \( \lambda_\perp \).

\(^{18}\) To stress the difference between local and global systems, here we do not use the language of \( \kappa \)-vectors. Wavevectors parallel and perpendicular to magnetic fields can be used, if only the wavevectors are understood in terms of a wavelet transform defined with the local reference system rather than the ordinary Fourier transform defined with the mean field system.
The critical balance is the feature of the strong turbulence, which is the case when the turbulent energy is injected at $V_A$. If the energy is injected at velocities lower than $V_A$, the cascade is weak with $\lambda_\perp$ of the eddies increasing while $\lambda$ stays the same (Ng & Bhattacharjee 1996; LV99; Galtier et al. 2002) In other words, as a result of the weak cascade the eddies get thinner, but preserve the same length along the local magnetic field. This decreases $\lambda_\perp$ and eventually makes Equation (A1) satisfied. If the injection velocity is $v_I$ and turbulent injection scale is $l$, the transition to the strong MHD turbulence happens at the scale $l (v_I/V_A)^2$ and the velocity at this scale is $V_{\text{turb}} = V_A (v_I/V_A)^2$ (LV99, Lazarian 2006). Thus the weak turbulence has a limited, that is, $[l, l (v_I/V_A)^2]$, inertial interval and gets strong at smaller scales.

While GS95 assumed that the turbulent energy is injected at $V_A$ at the injection scale $l$, LV99 provided general relations for the turbulent scaling at small scales for the case that the injection velocity $v_I$ is less than or equal to $V_A$, which can be written in terms of $\lambda_\parallel$ and $\lambda_\perp$:

$$\lambda_\parallel \approx l \left(\frac{l_\perp}{l}\right)^{2/3} \left(\frac{V_A}{v_I}\right)^{4/3} \quad \text{(A2)}$$

$$v_I \approx v_\perp l \left(\frac{l}{\lambda}\right)^{1/3} \left(\frac{v_\perp}{V_A}\right)^{1/3} \quad \text{(A3)}$$

The present-day debates of whether the GS95 approach should be augmented by additional concepts such as “dynamical alignment,” “polarization,” “non-locality” (Boldyrev 2005, 2006; Beresnyak & Lazarian 2006, 2009; Gogoberidze 2007) do not change the nature of the reconnection of the weakly turbulent magnetic field as LV99 also considered the modification of reconnection when $\lambda_\parallel \sim \lambda_\perp (V_A/v_I)^m$ and showed that for a choice of $p$ and $m$ which agrees with the present-day simulations the nature of the reconnection does not change.

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