Statistics of Reconnection-driven Turbulence

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Received 2016 October 27; revised 2017 February 6; accepted 2017 February 8; published 2017 March 29

Abstract

Magnetic reconnection is a process that changes magnetic field topology in highly conducting fluids. Within the standard Sweet–Parker model, this process would be too slow to explain observations (e.g., solar flares). In reality, the process must be ubiquitous as astrophysical fluids are magnetized and motions of fluid elements necessarily entail crossing of magnetic frozen-in field lines and magnetic reconnection. In the presence of turbulence, the reconnection is independent of microscopic plasma properties and may be much faster than previously thought, as proposed in Lazarian & Vishniac and tested in Kowal et al. However, the considered turbulence in the Lazarian–Vishniac model was imposed externally. In this work, we consider reconnection-driven magnetized turbulence in realistic three-dimensional geometry initiated by stochastic noise. We demonstrate through numerical simulations that the stochastic reconnection is able to self-generate turbulence through interactions between the reconnection outflows. We analyze the statistical properties of velocity fluctuations using power spectra and anisotropy scaling in the local reference frame, which demonstrates that the reconnection produces Kolmogorov-like turbulence, compatible with the Goldreich & Sridhar model. Anisotropy statistics are, however, strongly affected by the dynamics of flows generated by the reconnection process. Once the broad turbulent region is formed, the typical anisotropy scaling $l_\perp \propto t^{2/3}$ is formed, especially for high resolution models, where the broader range of scales is available. The decay of reconnection outflows to turbulent-like fluctuations, characterized by different anisotropy scalings, strongly depends on the $\beta$ plasma parameter. Moreover, the estimated reconnection rates are weakly dependent on the model resolution, suggesting that no external processes are required to make reconnection fast.

Key words: magnetic reconnection – magnetohydrodynamics (MHD) – methods: numerical – turbulence

1. Introduction

Magnetic reconnection is a key problem for the magnetohydrodynamic (MHD) theory. It describes the interaction of magnetic flux tubes, and in particular it describes what happens when magnetic flux tubes cross each other. It is impossible to fully predict how magnetic fields evolve in turbulent environments, ubiquitous in astrophysical fluids, without a solution to this question.

The theory of magnetic reconnection can be traced back to the classical Sweet–Parker scheme (Parker 1957; Sweet 1958), for which the magnetic field of opposite polarity are brought into contact over an extended region parallel to the flux and characterized by the scale $\delta$, while experiencing diffusion and annihilation over the region of thickness $\delta$. This $\delta$ determines the scale of the mass outflow, with the Alfvén speed $V_A$. Using mass conservation one can estimate the speed of crossing flux as $V_{rec,SP} \approx V_A \delta/\lambda$. With $\delta$ determined by microphysics and $\lambda$ being an astrophysical scale, it is evident that the Sweet–Parker reconnection rate is negligibly small for most astrophysical conditions. Indeed, taking into account that $V_{rec,SP}$ is determined by Ohmic diffusion, i.e., $V_{rec,SP} \approx \eta/\delta$, where $\eta$ is the resistivity, one gets the classical Sweet–Parker rate $V_{rec,SP} \approx V_A \eta/\delta$ with $S = LV_A/\eta$ being the Lundquist number, where $L$ is a scale of the flow. Given the highly conductive nature of astrophysical plasmas, Sweet–Parker reconnection predicts that reconnection occurs indeed very slowly.

For most astrophysical environments the Sweet–Parker reconnection essentially means virtually no reconnection at finite timescales, which contradicts the observations of, e.g., solar flares, for which the most natural explanation is based on fast magnetic reconnection. “Fast” in this situation means that the magnetic reconnection rate, $V_{rec}$, must not depend on, or at least exhibits very weak, e.g., logarithmic, dependence on, the Lundquist number $S$.

Different suggestions have been made to solve the problem of fast reconnection. They include solutions that involve making $\lambda$ small, i.e., similar to $\delta$ by bending magnetic field toward the reconnection point at a sharp angle, an ingenious suggestion by Petschek (1964). The corresponding magnetic configuration was termed “X-point reconnection.” A list of other solutions where the magnetic reconnection can achieve large speeds due to a particular configuration of magnetic fluxes can be found in the book by Priest & Forbes (2007).

The main limitation of these models is that they deal with the reconnection in rather special environments and do not address the problem of magnetic reconnection in generic astrophysical/turbulent circumstances. Could turbulence make reconnection fast? This issue has been discussed in a number of papers. While Jacobson & Moses (1984) dealt with the effects of turbulence on Ohmic resistivity, therefore decreasing $S$ by some factor, Matthaeus & Lamkin (1985, 1986) performed 2D numerical simulations of MHD turbulence and claimed that the formation of X-points in 2D made the reconnection fast. Due to significant differences of MHD turbulence nature in 2D and 3D, as well as the ambiguity of reconnection rate measurements within the 2D distribution of turbulent magnetic flux,
these interesting approaches could not be directly compared to real astrophysical scenarios.

A following model that related the modern theory of 3D MHD turbulence, i.e., the Goldreich & Sridhar (1995) one, and magnetic reconnection was proposed in Lazarian & Vishniac (1999, henceforth LV99). There the outflow region D is not determined by microscopic diffusive processes, but by the wandering of magnetic field lines. The prediction of the LV99 theory was that the reconnection changes with the level of turbulence. This quantitative prediction was successfully tested in the numerical studies of Kowal et al. (2009, 2012), investigating the effects of resolution, explicit resistivity, ways turbulence is driven, and strengths of guide field. More recently, another confirmation of LV99 came from relativistic MHD simulations by Takamoto et al. (2015). Observational testing of the theory are discussed, e.g., in recent reviews by Lazarian et al. (2015, 2016).

The LV99 predicts that reconnection is fast, i.e., does not depend on S, in generic astrophysical conditions. It predicts that reconnection happens not only at particular places where the magnetic field lines happen to undergo a special configuration, but through the entire turbulent volume. In fact, the theory makes the MHD turbulence theory by Goldreich & Sridhar (1995) self-consistent and predicts the violation of the classical magnetic flux freezing Alfvén (1942). The violation of flux freezing in turbulent fluids entails many vital astrophysical consequences (see Lazarian 2005; Santos-Lima et al. 2010; Lazarian et al. 2012), and it has been explored theoretically (Eynik 2011; Eynik et al. 2011) and confirmed numerically (Eynik et al. 2013).

While turbulence is really ubiquitous in astrophysical environments (see Armstrong et al. 1995; Padoan et al. 2009; Chepurnov & Lazarian 2010; Chepurnov et al. 2015), a question arises of whether reconnection itself could induce turbulence that would make it fast. This possibility was mentioned in LV99 and quantified in Lazarian & Vishniac (2009) within the model of flares of magnetic reconnection in high β-plasma environments. The first numerical simulations of magnetic reconnection induced by turbulence that is generated by reconnection were performed in Beresnyak (2013) under incompressible approximation, and in Oishi et al. (2015) and Huang & Bhattacharjee (2016) using compressible codes. However, these works did not consider a few of the studies that we address here. For instance, Beresnyak (2013) & Oishi et al. (2015) did not perform studies of turbulence properties and anisotropy statistics. Although Huang & Bhattacharjee (2016) included these studies, their setup was significantly different, and the statistics were calculated at relatively earlier time \( t = 3.5 \Delta t \). Moreover, none of these works addressed the question of turbulence development under different \( \beta \)-plasma parameters.

In this paper we aim to understand the reconnection-driven turbulence in a compressible MHD framework, focusing on long time evolution and turbulence statistics changes during the whole simulation. For that, we present a number of high resolution numerical experiments, for which we analyze and interpret the properties of obtained turbulent velocity fluctuations. We demonstrate the dependencies on the resolution and \( \beta \)-plasma. In Section 2, we present our methodology and preformed numerical models. In Section 3, we describe obtained results from the analysis of temporal evolution and statistical properties of generated turbulent fluctuations. In Section 4, we discuss our results and compare them to the previously done models. Finally, in Section 5 we draw our conclusions.

2. Methodology and Modeling

We use a high-order shock-capturing Godunov-type code AMUN\(^{1}\) based on the adaptive mesh. The code integrates the set of isothermal compressible MHD equations

\[
\frac{\partial \rho}{\partial t} + \nabla \cdot (\rho \mathbf{v}) = 0, \quad (1)
\]

\[
\frac{\partial \rho \mathbf{v}}{\partial t} + \nabla \cdot \left[ \rho \mathbf{v} \mathbf{v} + \left( a^2 \rho + \frac{B^2}{8\pi} \right) \mathbf{I} - \frac{1}{4\pi} \mathbf{B} \mathbf{B} \right] = \nu \nabla^2 (\rho \mathbf{v}), \quad (2)
\]

\[
\frac{\partial \mathbf{B}}{\partial t} + \nabla \times \mathbf{E} = 0, \quad (3)
\]

where \( \rho \) and \( \mathbf{v} \) are plasma density and velocity, respectively, \( \mathbf{B} \) is the magnetic field, \( \mathbf{E} = -\nabla \times \mathbf{B} + \eta \mathbf{J} \) is the electric field, \( \mathbf{J} = \nabla \times \mathbf{B} \) is the current density, \( \alpha \) is the isothermal speed of sound, and \( \nu \) and \( \eta \) are the viscosity and resistivity coefficients, respectively.

We integrated the governing equations using the 3rd order 4-stage Strong Stability Preserving Runge–Kutta method (Ruuth 2006), a 3D spatial Gaussian process-based reconstruction limited near extrema (G. Kowal 2017, in preparation), and a multi-state Harten–Lax–van Leer approximate Riemann solver Mignone (2007). In order to keep \( \nabla \cdot \mathbf{B} \) negligible, we solve the induction equation (Equation 3) using hyperbolic divergence cleaning based on a generalized Lagrange multiplier method (Dedner et al. 2002). This set of numerical algorithms, together with the stability CFL-coefficient being set to 0.3 for all models, results in stable numerically modeled data of high quality.

We performed a number of numerical simulations within a 3D domain with physical dimensions \( 1.0 \times 4.0 \times 0.5 \) with its center placed at Cartesian coordinates \((0, 0, 0)\). Our code makes use of adaptive mesh, for which the refinement criterion is based on the magnitude of the local vorticity \( |\mathbf{\Omega}| = |\nabla \times \mathbf{V}| \) and current density \( |\mathbf{J}| \). The thresholds determining the refinement and derefinement of the mesh are set to 0.1 and 0.01, respectively, indicating that the mesh is quickly refined in the turbulent and reconnection regions. The base domain is divided into \( 2 \times 8 \times 1 \) blocks, each with \( 32^3 \) resolution, from which we refine the mesh using the above mentioned criterion up to 3, 4, and 5 refinement levels with the resulting effective grid sizes \( \Delta x = \Delta y = \Delta z = h = 1/256, 1/512, \) and 1/1024, respectively. For the latter, the effective resolution is \( 1024 \times 4096 \times 512 \), therefore. For models with different \( \beta \)-plasma parameters, indicated with \( \delta \) in the figures, we use box which is squared in the \( XZ \)-plane, i.e., its physical dimensions are \( 1.0 \times 4.0 \times 1.0 \), with the maximum refinement level equal to 4, and the effective resolution \( 512 \times 2048 \times 512 \).

The initial magnetic field configuration is antiparallel along the \( X \) direction with the magnitude equal to 1.0 and a discontinuity placed at the \( XZ \) plane \( y = 0 \). An additional

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\(^{1}\) The code is freely available at http://amuncode.org.

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\( ^{2} \) The violation of flux freezing is implicit in LV99, but is explicitly treated in the subsequent proceedings paper Vishniac & Lazarian (1999).

\( ^{3} \) Plasma \( \beta \) is the ratio of the thermal to magnetic pressure.
3. Analysis and Results

3.1. Turbulence Evolution

We started our models with weak random velocity fluctuations of the order of 1% of the Alfvén velocity, which launches the magnetic reconnection process within the initial current sheet. This process normally converts magnetic energy into kinetic one and heat. Since we used the isothermal equation of state, the setup corresponds to a situation that the heat is quickly dissipated.

The first question we want to answer is whether the process of magnetic reconnection can sustain the injection of kinetic energy to the system or if the initial velocity fluctuations simply decay due to the dissipation. In order to answer it, we analyze our result first in terms of global quantities, such as the kinetic energy evolution, the evolution of velocity component contribution to the kinetic energy, the evolution of total vorticity generated by the reconnection, and the change of the volume fraction characterized by a high vorticity value.

Figure 1 shows the evolution of kinetic energy density integrated over the whole domain, up to the final simulation time $t = 20.0$. The initial value of the energy for all models is about $5 \cdot 10^{-7}$. We remind the reader that the turbulence is generated only in the vicinity of current sheet by the action of magnetic reconnection. Therefore, it can be dynamically important there. In most of the domain volume, however, the kinetic energy remains negligible. The figure demonstrates a quick growth of $E_{\text{kin}}$ to values above $10^{-3}$ at the relatively short time of one Alfvén time unit. This indicates an increase of the kinetic energy by almost three orders of magnitude. We should note that the initial velocity perturbations were generated from a random noise; therefore, they are strongly dissipated by the numerical dissipation during an initial period of about $0.02 \tau_A$, causing a small drop of kinetic energy, and reducing the strength of fluctuations interacting with the current density, as seen in Figure 1. This drop depends on the used grid size $h$ and sound speed $a$. For example, it is bigger and reaches the minimum at slightly later times (the minimum appears around $0.03 \tau_A$) for a model with the sound speed equal to $4.0$ (black line), compared to model with $a = 1.0$. Even though these fluctuations are weakened, they are still strong enough to sufficiently deform the current sheet locally, which results in the development of instabilities. For the scope of this paper, we do not identify the type of instability, and just focus on the characteristics turbulence produced as a result. After around $1.0 \tau_A$, the levels of kinetic energy are comparable for models with different sound speeds $a$ (or $\beta$) and grid sizes $h$.

Figure 2. Evolution of standard deviation of the velocity and its components: $X$, $Y$, $Z$, and parallel and perpendicular to the local magnetic field for model with sound speed $a = 1.0$ and the grid size $h = 1/512$.
Y, and Z-components of velocity are equal, but soon the X component, which is in the direction of the antiparallel magnetic field $B_x$, starts to dominate. Compared to the kinetic energy evolution, we see that this component carries most of the energy. The Z-component is slightly stronger than the Y one, since we imposed a non-zero guide field in the system.

Comparing components parallel (dashed) and perpendicular (dotted) to the local field, we see that the parallel one is developing to values a few times larger than the perpendicular one. This indicates that the strongest motion generation takes place along the mean magnetic field and is associated with the reconnection process. A similar behavior is observed for two components, $V_x$ and $\rho V_y$.

The evolution of kinetic energy and velocity components described above clearly demonstrates that the reconnection is able to generate significant plasma motions, which could dominate the local dynamics. Below, we investigate more closely what type of motions we observe in our models: turbulent or laminar.

The simplest way of detecting turbulent eddies in a system is by calculating the vorticity, $\omega = \nabla \times v$, which directly measures the rotational motions. In the first analysis, we calculate the magnitude of vorticity and average it over the computation domain at each time step. We also analyze the reconnection rate estimated using two different methods. Later, we analyze their spectral properties and anisotropy.

The top left plot of Figure 3 shows the average vorticity magnitude $|\omega|$ as a function of time for all models. The amount of turbulent motion generated due to reconnection is growing until it reaches saturation at around $t = 5.0$. The vorticity growth rate seems to gradually decrease with time until saturation. Since the reconnection works at the smallest scales, the vorticity growth depends on the resolution, which can be seen by different maximum values of averaged vorticity, $|\omega| \approx 4.0, 5.0$, and 8.0 for models with grid sizes $h = 1/256, 1/512, \text{and } 1/1024$, respectively. This dependence with the grid size is not linear and may indicate that, for smaller grid sizes, even higher vorticities can be generated. After reaching the maximum, the vorticity generation slowly decays for models with medium and large grid sizes. It increases for the model with the highest resolution, however. In three models with different sound speeds ($\alpha = 0.5, 1.0, \text{and } 4.0$, respectively), the total vorticity grows through the whole simulation, but its value is $\alpha$-dependent, i.e., it decays with $\alpha$ (or $\beta$).

The top right plot of the same figure shows the filling factor for vorticity $|\omega| > 10.0$. It simply shows a volume where vorticity is equal to or larger than the threshold, and indicates...
how the volume of strong rotational motions changes. The chosen threshold can be interpreted as a change of velocity by 0.1$V_A$, within a distance 0.01, which corresponds to relatively strong rotation or shear. In our simulations, we observe peak values of vorticity as high as a few hundreds, related to very strong reconnection events (change of the order of $V_A$ at the grid size scale). The top right plot in Figure 3 demonstrates the initial, nearly linear growth of the volume occupied by relatively strong rotational motions, almost independent of the used resolution or sound speed, until about $t = 6.0$. This moment, the point at which the growth rate of the filling factor suddenly changes, is also independent on the effective resolution $h$. The maximum value, however, is resolution sensitive, with the highest volume observed in the model with the smallest grid size. The evolution of filling factor after this moment also has differing behavior for different resolutions. For the model with the largest grid size, the filling factor starts to decay after $t \approx 8.0$, for the middle grid size it stays nearly constant until $t \approx 16.0$, and for the smaller grid size it seems to continue growing. This indicates that, in order to resolve the turbulence well, we have to use the highest resolution possible. The filling factor evolution seems to be very weakly dependent on the plasma $\beta$ parameter. Comparing models with different $\alpha$ and the same grid size $h = 1/512$, we see that the filling factor follows similar evolution growing in all cases, but reaching different filling factor values for different sound speeds. Comparing the model with $\alpha = 0.5$ (dashed line) to models with $\alpha = 1.0$ and $\alpha = 4.0$, the filling factor reaches values 5%–20% and 20%–30% lower, respectively, after $t = 5.0$.

In the two bottom plots of Figure 3, we show two ways of estimating the reconnection rate. In the right plot of Figure 3, we present the evolution of the estimated thickness of the turbulent layer $\Delta$. The thickness was calculated from the region where the magnitude of vorticity is above a threshold, in our case $|\omega| \geq 10.0$. As we see, initially the thickness continues to grow linearly, almost independently of the grid size, and after $t \approx 6.0$ its growth slows down. For the lowest resolution model ($h = 1/256$), it continues to grow up to $t = 9.0$ and then saturates. In the case of the medium resolution model ($h = 1/512$), $\Delta$ continues to increase until $t \approx 16.0$, and for the highest resolution model ($h = 1/1024$) it increases with the same rate until the end of simulation. For models with different sound speeds, the turbulent region thickness grows during the whole simulation. In the subplot, we show reconnection rates for all models estimated from the growth rate of the turbulent region thickness $V_{\text{rec}} = d\Delta/dt$ within the initial period $t \leq 5.0$, in the same way as was done by Beresnyak (2013). The x points correspond to models with different $\beta$, while the + points are for models with different resolutions and the same $\alpha = 1.0$. We see that the reconnection rate depends slightly on $\beta$ plasma, decreasing with its value. We estimated $V_{\text{rec}}$ to be about 0.0327 $\pm$ 0.0003, 0.0302 $\pm$ 0.0002, and 0.0255 $\pm$ 0.0003 for $\beta \approx 0.5$, 2.0, and 32.0, respectively. From these points we estimated an empirical dependence $V_{\text{rec}}(\beta) = -0.0017 \log \beta + 0.0314$, with both coefficients rounded to the fitting error. This weak dependence could indicate the importance of supersonic motions on bending magnetic field lines within the current sheet. There the reconnection outflows can approach the Alfvén speed, which means that for the low $\beta$ plasma, the outflows become supersonic. Also, because the magnetic field strength is reduced, it is easier to bend and effectively reconnect its lines.

The reconnection rate $V_{\text{rec}}$ also weakly depends on the used resolution, i.e., it grows with the grid size $h$. The resolution dependence could be related to the numerical resistivity, which is proportional to $h$, causing the reconnection efficiency to be higher for models with bigger $h$. For the model with the highest resolution ($h = 1/1024$), the estimated reconnection rate is about 0.0287 $\pm$ 0.0003, nearly twice as much as the value estimated by Beresnyak (2013). Considering that Beresnyak (2013) used a much less dissipative spectral code and smaller values of the resistivity, our result should be compatible with theirs.

As we indicated, the reconnection rates presented in the bottom left plot were estimated for times $t \leq 5.0$. At later times, the growth rate of turbulent regime depends on the grid size $h$. Once the turbulence is fully developed, we should expect the reconnection to perform more efficiently, as predicted by LV99. The observed tendency is that for $t \geq 5.0$ the reconnection rate estimated using the thickness of turbulent region decreases with $h$, and in the case of the lowest resolution model ($h = 1/256$) it even drops down to zero. For the highest resolution model, the turbulent region thickness $\Delta$ always increases, yet its growth rate is not as high as at earlier periods $t \leq 5.0$. In real systems, however, where the resolution is virtually infinite, we should expect that the thickness of turbulent region grows constantly at the same rate for later times too. In models with different sound speeds, we also see a decrease of the growth rate of the turbulent region thickness after $t \approx 5.0$; however, it is roughly independent of the used $\beta$.

In the bottom right plot of Figure 3, we show the mean inflow speed $V_{\text{in}}$ measured at the upper and lower boundaries at $y = \pm 2.0$. Since these boundaries are open, the magnetic flux can be freely transported through them. The positive values of $V_{\text{in}}$ indicate that the velocity at the boundary is directed toward the interior of the domain, therefore, it brings fresh magnetic flux to the system. In the standard Sweet–Parker model we interpret this quantity as a measure of the reconnection process efficiency. The only way magnetic field can be changed in such a system is by its conversion into other forms of energy, dissipation, or transport through the open boundaries. Since our inflow boundaries are far from the initial current sheet, we observe zero inflow during the initial period, followed by a very quick burst of the inflow and then a decay. The maximum value of $V_{\text{in}}$ and its following decay seem to be nearly independent of the grid size $h$. This indicates that the reconnection process under the same physical conditions is independent of numerical dissipation, even if its rate reaches values that are only small factors of $V_A$. These small values of $V_{\text{in}}$ are mostly due to the fact that our boundaries along the $X$ direction are periodic, so the reconnected flux is not removed from the system. Comparing models with different sound speeds $\alpha$ (or plasma $\beta$ parameters), we see that the inflow speed $V_{\text{in}}$ is very sensitive to $\alpha$. The moment at which $V_{\text{in}}$ starts to quickly grow happens at different times for models with different $\alpha$, and $V_{\text{in}}$ reaches significantly different values at later times. We verified that, within the turbulent region, the maximum sonic Mach number reaches around 2.0 for the model with $\alpha = 0.5$, and around 1.0 or below 0.25 for models with $\alpha = 1.0$ and 4.0, respectively, signifying that supersonic turbulence can be generated in models with $\alpha < 1.0$. This means that the compression may be “squeezing” more magnetic field within the turbulent region and/or it is more efficiently dissipated in local shocks. This process would
saturate in the presence of open boundary conditions along the current sheet, allowing for removal of the reconnected magnetic field. However, with the periodic boundaries the flux is accumulated and therefore the inflow speed may not be a reliable measure of the reconnection rate in these systems.

Figure 4 shows slices of total vorticity along the three main mid-planes at three different moments, \( t = 1.0, 5.0, \) and \( 20.0 \). All plots have the same color range for ease of comparison and are obtained from the model with the smallest grid size \( h = 1/1024 \). We see that the stochastic reconnection develops complex filamentary structures near the current sheet, increasing their volume with time, which supports the filling factor evolution shown in the right plot of Figure 3. It is important to recognize that most of the turbulent motions are developed in the ZY-plane perpendicular to the reconnecting field. Vorticity shows places in which the velocity changes quickly. As we see in the middle column (the ZY-plane), the high vorticity filaments have arbitrary orientations. In the right plots (the XY-plane), however, the filaments tend to align with the \( X \) direction, especially at later times (central and bottom plots).

This is explained well by the fact that the motions can mix reconnecting field lines more easily in the perpendicular plane than in the parallel direction due to the magnetic tension, even if these motions are produced by the reconnection of the same lines. These visualizations demonstrate how a simple current sheet with a weak velocity noise can create complex turbulent structure within a broad vicinity.

### 3.2. Turbulence Properties: Power Spectra and Anisotropy

In the previous subsection we studied the evolution of turbulent motions produced by stochastic reconnection and initiated from a random velocity noise. We analyzed how the kinetic energy and vorticity evolves with time, the growth rate of the turbulent region thickness, and how the estimated reconnection rate depends on the used grid size or \( \beta \). Here we take a closer look into the statistical properties of turbulent motions by studying their power spectra and anisotropy. The fundamental question we want to answer here is if the turbulence generated by reconnection can be characterized a by a mixture of strong turbulence, described by the Goldreich & Sridhar (1995) model, and structures produced by the reconnection ejection regions.

In Figure 5 we show power spectra of velocity (left) and vorticity (right) for different moments during the turbulence development. The power spectra were calculated using 3D wavelets in Fourier space (see Kirby 2005). For comparison, we show Kolmogorov slope \( k^{-5/3} \) for velocity, and corresponding \( k^{1/3} \) for vorticity) using a gray dashed line. The initial power spectrum \( (t = 0.0) \) is shown using the yellow line. The velocity fluctuations are spread over the smallest scales with very small amplitudes \( (P(k) < 10^{-7} \) for velocity, see the left plot). Fluctuations of both quantities, velocity and vorticity, quickly develop from small to large scales. At \( t = 0.5 \), they are already spread over large range of scales (up to \( k \approx 10 \)). After \( t > 2.0 \), power spectra start to align with the Kolmogorov slope forming the inertial range. At \( t = 5.0 \) the inertial range is well developed and it extends from \( k \approx 10 \) down to nearly \( k = 200 \) for models with the effective grid size \( h = 1/1024 \). At the final moment of simulation, \( t = 20.0 \), the power spectrum is fully developed and stationary (compare its change for \( t \geq 10 \)), both for velocity and vorticity, and is characterized by a slope close to the Kolmogorov one.

The development of a broad power spectrum from small to large scales cannot be simply explained by reconnection, which usually operates at small scales across the current sheet. If considered as one reconnection event, the ejection occurs along a thin slab determined by the local current sheet thickness. However, within the local ejection region, the magnetic field is oriented perpendicularly to the current sheet plane, since the reconnected component is removed along the slab. On the other hand, a typical fluctuation scale along the current sheet is determined by the local separation of the reconnection events. Initially, this separation scale is very short in our case due to densely packed reconnection events, and later it increases with time due to the interactions between ejections acting at larger and larger scales. This picture is justified by plots in the left column of Figure 4, where a typical scale of vorticity structures seems to increase with time (compare ZX-plane cuts for different time moments). This also indicates a development of velocity anisotropy dominated by initially small and nearly isotropic structures and transformed into anisotropic ones. This picture can be confirmed if we estimate the anisotropy with respect to the mean direction of reconnecting components, i.e., along the \( X \) direction. However, what is the anisotropy with respect to the local mean field and the reconnecting one? Should we expect two different scalings, i.e., Goldreich & Sridhar-like at large scales and a different one produced by reconnection outflow regions at smaller scales? Up to which scales can the outflows affect the statistics? How is the anisotropy developed for different \( \beta \)-plasma parameters?

In Figure 6 we show the second-order structure functions of velocity calculated in the local reference frame, i.e., with respect to the local mean magnetic field at three different time moments: \( t = 0.5 \), \( 1.0 \), and \( 10.0 \) (left, center, and right columns, respectively) for models with \( a = 0.5 , 1.0 \), and \( 4.0 \) (upper, middle and lower rows, respectively). We notice the change of anisotropy degree \( \beta_i/\beta_l \) between different moments. For all three models with different \( a \) at time \( t = 0.5 \) the velocity fluctuations are elongated with the local magnetic field. We clearly see that the anisotropy decreases with increasing \( a \) (or \( \beta \)). However, once the reconnection increases its efficiency, the velocity fluctuations instead start to be elongated in the direction perpendicular to the local field (see the central column). At \( t = 1.0 \), the velocity fluctuations are relatively isotropic for model with \( a = 0.5 \) (for \( l < 0.1 \)), while for other models they are elongated perpendicular to the local field.

Why this sudden change of orientation of anisotropic velocity structures? Our understanding is that the Alfvén waves initially dominate the structure of velocity, but the local reconnection events quickly become stronger, taking over the statistics. A local reconnection event can be characterized by a weak and broad inflow of the magnetic flux, and a strong and thin outflow region (see the typical Sweet–Parker picture). Within the outflow region, the velocity is oriented in the direction perpendicular to the local field (i.e., with respect to the reconnected flux component and the guide field), therefore the flow is perpendicular to this field. This is clearly observed in the right top plot of Figure 6 showing the local structure function at \( t = 10.0 \), where its value increases quickly along the parallel direction, determining the averaged thickness of the current sheet \( \delta < 0.05 \), while in the perpendicular direction it increases slowly up to a fraction of the length unit. This anisotropy continues to dominate in the perpendicular direction during the later times of our simulation (see the right plots of
Figure 4. Vorticity amplitude slices along three main planes, ZX (left), ZY (center), and XY (right) for three different moments in time, $t = 1.0, 5.0, \text{ and } 20.0$ (top, middle, and bottom, respectively) for a model with sound speed $a = 1.0$ and grid size $h = 1/1024$. The turbulent region develops near the initial current sheet and expands in the $Y$ direction. The plots show subregions limited to $y \leq 0.5$ from the whole domain extended up to $y = 2.0$. 

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Figure 5. Evolution of power spectra of velocity (left) and vorticity (right) for the model with the grid size \( h = 1/1024 \). We plot several power spectra calculated at different times.

Figure 6 for \( a = 0.5 \) and 1.0). At later time \( t = 10.0 \) for the model with the highest sound speed \( a = 4.0 \), the anisotropy degree indicates that fluctuations change their orientation and are aligned with the local magnetic field, in contrast to models with \( a \leq 1.0 \). This indicates that the compressibility could be an important factor in determining the small scale anisotropy in reconnection-driven turbulence, since it increases with decreasing sound speed.

As we have discussed above, the reconnection generates strong anisotropy in the direction perpendicular to the local field, and during the evolution the anisotropy degree changes significantly. The dominance of the perpendicularly oriented fluctuations is observed in the whole evolution of models with \( a \leq 1.0 \) and during the initial few Alfvén time units in the model with \( a = 4.0 \). The difference comes from the fact that the reconnection outflow interactions decay to turbulent fluctuations, and the timescale for this process seems to depend on the fast magnetosonic speed, since it is faster for higher sound speeds. Of course, this also results in the same timescale for anisotropy changes. The perpendicular anisotropy dominance is in contrast to the Goldreich–Sridhar turbulence, where velocity eddies are elongated with the magnetic field. We have explained that the difference comes from the strong reconnection outflows dominating the structure function, especially at small scales. However, even though the fluctuations are not oriented along the local field, maybe at least the scaling of anisotropy reveals some compatibility with the Goldreich–Sridhar model?

In Figure 7 we show anisotropy scaling of velocity fluctuations at several moments, \( t = 0.5, 2.0, 5.0, 10.0, \) and 20.0 for the model with \( a = 1.0 \) (\( h = 1/1024 \), left plot), and \( t = 0.5, 1.0, 2.0, 5.0, \) and 10.0 for the model with \( a = 4.0 \) (\( h = 1/512 \), right plot). These plots confirm that the anisotropy degree varies significantly during the system evolution, especially at the beginning of the simulation. Initially isotropic fluctuations develop strong anisotropy at all scales along the perpendicular direction \( l_\perp \) (shown in the horizontal axis) with scaling close to \( l_\perp \propto l_x \), but after about \( t = 5.0 \) for the model with \( a = 1.0 \) and \( t = 2.0 \) for the model with \( a = 4.0 \) the anisotropy degree tends to decrease. Even though the anisotropy changes significantly, its scaling seems to be relatively consistent up to \( t = 10.0 \) for the model with \( a = 1.0 \). We observe that for scales below 0.1 it continues to follow \( l_\parallel \propto l_x \), while only at larger scales it starts to manifest the Goldreich–Sridhar scaling \( l_\parallel \propto l_x^{2/3} \). On the right, the anisotropy scaling becomes compatible with \( l_\parallel \propto l_x^{2/3} \) earlier. We see it is formed already at \( t = 5.0 \) from \( l_x \approx 0.1 \) down to scales as small as 0.01. At this time, the turbulent region is relatively thick, as estimated in the bottom left plot in Figure 3. The developed turbulence have enough space to create inertial range and structure compatible with strong turbulence. At the same time reconnection events spread sparsely over the same turbulent region, affecting only the local velocity structures and relaxing to turbulent motions at shorter scales for the model with higher sound speed. This indicates the presence of strong turbulence described by the Goldreich–Sridhar theory, substantially disturbed by the injection from the reconnection outflows at smaller scales, characterized by different properties.

3.3. Decaying Turbulence Without Global Field Reversal

In order to remove the strong contribution of the reconnection outflows observed in the anisotropy scaling of velocity fluctuations, which is responsible for deviation from Goldreich & Sridhar (1995) scaling in the model with sound speed \( a = 1.0 \), we performed a run with the grid size \( h = 1/512 \) in which we removed the global reversal of magnetic field by setting \( |B_y| \) everywhere and leaving other two components, \( B_x \) and \( B_z \), unchanged. From this initial state, we restarted our simulation and let the fluctuations evolve without energy injection from the reconnection process. Due to the lack of global reversed field, the reconnection should be strongly suppressed. We should note that the reconnection can still produce weak energy injection from the unchanged components of magnetic field. Nevertheless, the main source of injection should be removed, and previously produced velocity and magnetic field fluctuations should start to decay due to the dissipation. We expect that, due to the lack of reconnection outflows, the fluctuations should quickly produce scaling compatible with the Goldreich–Sridhar model. We ran this model for a few Alfvén times, and then calculated local structure functions in order to analyze the velocity anisotropy scaling.

In Figure 8 we compare anisotropy scalings for driven and decaying models at the same moment \( t = 12.0 \), i.e., the reconnection-driven turbulence shown with \( \times \)-points, and the
decaying turbulence in which the global magnetic field reversal was removed (open and closed circles for \(t = 11.0\) and \(t = 12.0\), respectively). We clearly see the difference in the anisotropy scaling. In the reconnection-driven case, the outflows still strongly affect the velocity statistics and a scaling closer to \(l_{\|} \propto l_{\perp}^{1/3}\) is observed. In the case of decaying turbulence, the reconnection outflows do not affect the statistics anymore and after one Alfvén time the scaling is already aligned with \(l_{\|} \propto l_{\perp}^{2/3}\) at larger scales. The same is observed at later time \(t = 12.0\). This test indicates that the reconnection outflows are responsible for the change of anisotropy scaling.

4. Discussion
4.1. Comparison to Previous Results

We performed numerical modeling of reconnection-driven turbulence and studied its properties. We should note that we did not impose initial Sweet–Parker or Harris current configurations (as in Lapenta 2008; Oishi et al. 2015; Huang & Bhattacharjee 2016), but we let a weak noise affect the current sheet along which magnetic field is discontinuous initially. Moreover, in this work we applied periodic boundary conditions along the current sheet, but we left the vertical
boundaries open. With the exception of the vertical boundary type, our setup closely resembles the one studied by Beresnyak (2013). The choice of different vertical boundaries has some important consequences. Periodicity in all three directions with two separated current sheets imposed, as used in Beresnyak (2013), allows for the possibility of horizontal large scale interactions through the deformations of both current sheets. These interactions, especially at later times, can generate large scale motions in the perpendicular direction to the current sheet providing additional to the reconnection energy input (see Drake et al. 2010; Kowal et al. 2011, for a similar setup with many current sheets, in which the current sheet deformations are well manifested). In our case, such large-scale interactions are not allowed, therefore the only energy input comes from the small-scale reconnection events. The reconnection rate values obtained from our models are compatible to those in Beresnyak (2013). We estimated the reconnection rate to be about $-V_0^{0.025}$, while the value estimated in Beresnyak (2013) is around twice smaller. This can be justified by less diffusive spectral code together with smaller resistivity coefficient (of the order of $4 \cdot 10^{-5}$, as estimated from the Lundquist number for their lowest resolution run) used in their work. The spectral slope of the turbulence generated in their models was estimated to be around Kolmogorov slope, which is comparable with the one we show in Figure 5. Beresnyak (2013) did not demonstrate any results related to the properties of velocity fluctuations, however, and since they focused only on incompressible MHD regime, they did not study $\beta$-dependence.

As we mentioned already, Oishi et al. (2015) used a different numerical setup starting from the Sweet–Parker configuration and periodic box along the current sheet. The Sweet–Parker configuration requires open boundary conditions to allow inflow and outflow of the magnetic flux in order to maintain steady state. Since the periodic boundary conditions were applied, the growth of the fluctuations at large scales may be attributed to the presence of the configuration of magnetic field which is not in initial equilibrium. It is also important to notice that these authors have misleadingly claimed to perform very long simulations up to $1000t_A$. They assumed an uncommon definition of Alfvén time $t_A = \delta_0/V_A$, where $\delta_0 = 0.02$ was the current sheet thickness. Typically, the timescale is defined by $t_A = L/V_A$, where $L$ is the size of the system, which results in

![Figure 7](image1.png)

**Figure 7.** Anisotropy scalings for total velocity for models with different sound speeds ($a = 1$ on the left and $a = 4$ on the right) obtained at different moments. For comparison, we show isotropic and strong turbulence scalings (denoted by gray thin lines). The vertical dashed line in the right plot shows the effective grid size $h = 1/512$. In the left one, the grid size is below the left axis.

![Figure 8](image2.png)

**Figure 8.** Comparison of the anisotropy scalings for velocity in the reconnection-driven turbulence run with the global field reversal at $t = 12.0$ (× points) and at two time moments $t = 11.0$ and 12.0 (open and closed circles, respectively) for a model in which the global reversal of magnetic field was removed at $t = 10.0$ and the simulation was restarted. We see that, while in the reconnection-driven run the reconnection outflows affect the velocity statistics and the scaling is closer to the isotropic one, in the run without strong outflows (no field reversals) the anisotropy scaling tends to the Goldreich & Sridhar (1995) scaling, $l_\parallel \propto l_\perp^{2/3}$, at scales $l_\perp > 2 \cdot 10^{-2}$. The Astrophysical Journal, 838:91 (13pp), 2017 April 1 Kowal et al.
their maximum simulation time of 20\(t_A\), same as in our models. Another important conclusion claimed by Oishi et al. (2015) was that for sufficiently large Lundquist number \(S\), the \(\xi_{SP}\) becomes weakly dependent on \(S\) (see their Figure 1). However, the authors did not provide any estimation of the numerical resistivity \(\eta_{num}\), which we expect to be comparable with the one we estimated since both studies were performed using Godunov-type codes.9 From Table 1 in Oishi et al. (2015) we can see that, for high Lundquist number models, the explicit resistivity \(\eta\) is far below \(\eta_{num}\) we estimated. Therefore, the deviation from Sweet–Parker dependence \(V_{rec} \propto S^{-1/2}\) observed in their simulation could be artificial and result from the numerical effects.

Recently, Huang & Bhattacharjee (2016) analyzed the statistics of velocity fluctuations generated by reconnection. They also used, as a starting point, a global Sweet–Parker configuration, but with non-periodic boundaries along the current sheet. Weak velocity perturbations were injected into such a domain in order to understand their effect on the reconnection rate. If the perturbations are injected too far from the current sheet, they simply propagate out of the domain as Alfvén waves before reaching the reconnection zone. They used a guide field comparable in strength to the reconnecting component, which additionally prevents the interactions of the waves. They calculated windowed 2D power spectra of velocity over \(XZ\)-planes and averaged them over the vertical direction. Their power spectra demonstrated a steeper slope compared to Kolmogorov turbulence. The structure functions were also calculated over \(XZ\)-planes, and the local reference frame was determined by the in-plane magnetic field. Here, we applied the same technique for calculation of local structure function in turbulent simulations, and found that this method results in much reduced anisotropy, when compared to the fully 3D structure function analysis. In Figure 9 we present the comparison between both methods applied to the driven turbulence model without global field reversal simulated in a fully periodic box with grid size \(h = 1/768\). The reduction of the anisotropy scaling due to a different technique is clear. This indicates that these authors’ conclusion of a different type of turbulence (steeper power spectrum and different than Goldreich–Sridhar anisotropy scaling) should be taken carefully. In our numerical experiments we obtain Kolmogorov-like power spectra of generated velocity fluctuations, which also present Goldreich–Sridhar anisotropy scaling at later times in large \(\beta\) models. This indicates that the turbulent statistics are similar to strong MHD turbulence, but are strongly affected (contaminated) by the dynamics of reconnection induced flows. In particular this is clearly visible for low \(\beta\) models, where supersonic reconnection outflows are present. In high \(\beta\) models we observe that the Goldreich–Sridhar anisotropy scaling is visible at much earlier times, since the fluctuations generated by reconnection outflows are propagated faster. We also claim that the Goldreich–Sridhar anisotropy scaling is more apparent for models with higher resolution due to larger dispersion in scales between the turbulent fluctuations and reconnection outflow, which is confirmed by Figure 7. As a consequence, our conclusions differ from those of Huang & Bhattacharjee (2016) with regards to the type of turbulence in magnetic reconnecting plasmas.

In Kowal et al. (2009) we estimated that \(\eta_{num} \approx 10^{-4}\) for models with grid size \(h = 1/512\).

4.2. Validity of Our Approach

In this paper we study magnetic reconnection using MHD description of a plasma (see Section 2). The validity of this description has been broadly discussed in the literature (see, e.g., Kulsrud 1983; Somov 2006a; Eyink et al. 2011; Lazarian et al. 2015), and its justification can be based on either collisionality or strong magnetization. Following Eyink et al. (2011), where we refer the reader for more detailed review of the validity of MHD description, we consider three characteristic length scales of importance: the ion gyroradius \(r_i\), the mean free-path length \(\lambda_{mfp,i}\) due to Coulomb collisions, and the scale \(L\) of large-scale variations of magnetic field and velocity. The fluid picture of a plasma is justified if the plasma is at least somewhat collisional, i.e., \(L \gg \lambda_{mfp,i}\). This regime can be divided into “weakly collisional” \((\lambda_{mfp,i} \gg r_i)\) and “strongly collisional” \((r_i \gg \lambda_{mfp,i})\) plasmas. In the opposite case, i.e., \(\lambda_{mfp,i} \gg L\), the plasma is “collisionless,” and hydrodynamic description does not work anymore. In astrophysical plasmas we can encounter all three regimes of collisionality. For example, plasmas in star interiors and accretion disks are strongly collisional. Very hot and diffuse plasmas, such as a warm ionized interstellar medium, are weakly collisional, while solar wind at magnetosphere or post-coronal mass ejection current sheets are examples of collisionless plasmas (see Table 1 in Eyink et al. 2011). Among the additional assumptions applied to MHD description is the assumption of high conductivity, which is easily fulfilled for most astrophysical plasmas (see, e.g., Somov 2006b).

In this work we are interested in reconnection processes at length scales much larger than \(r_i\). For such scales, plasmas can be more precisely described by “kinetic MHD equations,”
which differ from the standard MHD equations by the isotropic thermal pressure \( p \) replaced with the pressure tensor, which has two components, \( p \parallel \) and \( p \perp \), parallel and perpendicular to the local magnetic field, respectively (see, e.g., Chew et al. 1956). In the presence of anisotropic pressure, plasma can develop kinetic instabilities, such as “firehose” or “mirror” instabilities, which strongly affect the plasma evolution at small scales (see, e.g., Hau & Wang 2007, and references therein). Our focus in the studies here is to prohibit development of any microscopic instabilities (by using the particle-in-cell or kinetic MHD approaches) or enhanced reconnection rates (by using Hall terms or anomalous resistivity), which could influence the generation of turbulence by the reconnection process, even if other frameworks better describe the astrophysical plasma of interest in which the reconnection process takes place.

### 4.3. Restrictions Imposed by Periodic Boundary Conditions

As we mentioned in Section 2, the choice of periodic boundary conditions along the \( X \) and \( Z \) directions imposes some restrictions on the applicability of the model. For instance, the fluctuations of developed turbulence in the parallel direction to the initial current sheet (the \( XZ \)-plane) cannot reach sizes larger than the computational domain, therefore after around \( t \approx t_A = V_A^{-1} \), the energy could be accumulated at large scales due to the interactions with the domain boundaries and enforce the inertial range production down to the small scales. This, as a result, may affect the development of anisotropy. On the other hand, dissipation of the cascade-originated waves occurs on similar timescales. One would then expect the energy at small scales to be greatly dominated by the local cascade and less affected by the boundary effects. We performed models with different ratios of the \( Z \) to \( X \) dimensions of the box \( (L_z = 0.5L_x \text{ versus } L_z = L_x) \) and did not see significant changes in the developed spectra and anisotropy statistics. The problem of periodic boundary conditions, however, must be properly addressed and quantified in a follow-up paper.

### 4.4. What Drives the Observed Turbulence?

One of the most important results of this work is the reassurance of self-generation of turbulence in reconnection events. The other is that this turbulence follows standard Kolmogorov and Goldreich–Sridhar statistics. However, a still open issue is related to the driving mechanism of the observed turbulent motions. What drives the turbulence in reconnection events? The literature has suggested, without any quantitative proof, that tearing modes, plasmoid instabilities, and shear-induced instabilities could mediate the energy transfer from coherent to turbulent flows.

In our numerical experiments we do not identify tearing modes, although filamentary plasmoid-like structures are present. The filling factor of these are, however, visually recognized as very small. Sheared flows, on the other hand, are present around and within the whole current sheet. As the field lines reconnect, the \( \mathbf{v} \times \mathbf{B} + \mathbf{E} \) force increases, accelerating the plasma and creating the current sheet. This process is, in three dimensions, patchy and bursty. Therefore, the accelerated flows are strongly sheared. The statistical importance of these burst flows is large, as already shown in this work, as we compared the velocity anisotropy of reconnecting events to that of decaying turbulence without the reversed field. Kelvin–Helmholtz instability due to the sheared velocities in reconnecting layers has already been conjectured as possible origin of turbulence by Beresnyak (2013). In this work, we therefore provide real evidence for the self-generated turbulence driven by the velocity shear. Nevertheless, a proper analysis of the growth-rates of such instabilities must be conducted in future work.

Another consequence of the mechanism responsible for turbulence self-generation is on the statistics of perturbations. Velocity shear is a global process that occurs in regular magnetized and unmagnetized fluids. The nonlinear evolution of related instabilities, such as the Kelvin–Helmholtz instability, is known to be one of the main contributors to the energy transfer rate between wavemodes, i.e., the energy cascade. If the energy cascade in reconnection layers is led by similar mechanisms, it is straightforward to understand why the statistics observed resemble those of Kolmogorov-like turbulence, and Goldreich–Sridhar anisotropy scaling. In other words, our claim is that the turbulent onset and cascade in reconnection is not different to those found in regular MHD and hydrodynamical systems.

### 5. Conclusions

We performed numerical modeling of a realistic setup in which magnetic reconnection can develop turbulence from the initial weak noise of velocity fluctuations. We analyzed the time evolution of several quantities, including kinetic energy, its distribution among the velocity components, mean vorticity, its filling factor, the thickness of turbulent region, and we estimated the reconnection efficiency using two independent methods. Once the turbulence was developed in a broad region, we analyzed its properties using power spectra of velocity and vorticity, and anisotropy scaling of velocity fluctuations with respect to the local mean field.

The conclusions coming from these studies can be summarized in several points:

1. Reconnection is able to develop a substantial amount of turbulence from initially weak noise of velocity fluctuations. We observe growth of the kinetic energy by over three orders of magnitude.

2. The generation of rotational motions measured by vorticity is initially very quick, and reaches saturation after around 5.0\( t_A \). The saturation level of mean vorticity depends on the grid size \( h \) and \( \beta \), being higher for models with smaller effective grid size and smaller \( \beta \). The filling factor of vorticity with a threshold \(|\omega| \geq 10.0\) also depends on \( h \) and \( \beta \); however, for the model with the largest \( h \), the filling factor decays after reaching a maximum value around 7.0\( t_A \), while for the highest resolution model it constantly grows until the end of simulation. This indicates that the filling factor of strong rotational motions should be much higher in realistic systems. In models with \( L_z = L_x \), the filling factor of vorticity grows for all models with different sound speeds, even though the effective grid size was \( h = 1/512 \).

3. The estimated reconnection rates using two independent methods, i.e., the growth rate of the turbulent region thickness \( d\Delta/dt \) and the inflow speed \( V_{in} \), with which the fresh magnetic flux is brought to the system, are independent or weakly dependent on the resolution for similar \( \beta \), indicating that the reconnection could be fast without needing the presence of external processes, such
as driven turbulence, or collisionless plasma effects. The dependence of $V_{in}$ on $\beta$ might be related to additional magnetic field dissipation in supersonic shocks created within the turbulent region.

4. The velocity and vorticity present power spectra that are compatible with Kolmogorov slope. There is a small bump observed at small scales which could be explained by the action of reconnection or sort of bottleneck effect.

5. The anisotropy degree and scaling depend on the $\beta$-plasma parameter, related to the timescale of decay of the reconnection outflow interactions to turbulent fluctuations.

6. The velocity fluctuations generated by reconnection-driven turbulence are compatible with the Goldreich & Sridhar (1995) model, however, their statistics are strongly distorted by reconnection outflows driving the turbulence, especially for low $\beta$ models, where supersonic motions can be generated. In high $\beta$ models, the velocity anisotropy follows the $l_{\parallel} \propto l_{\perp}^{2/3}$ scaling. This scaling is also visible at large scales for the highest resolution model with $\beta \approx 1.0$.

G.K. acknowledges support from FAPESP (grants no. 2013/04073-2 and 2013/18815-0) and PNPD/CAPES (grant no. 1475088) through a Postdoctoral Fellowship at University Cruzeiro do Sul. This work has made use of the computing facilities of the Laboratory of Astrophysics (EACH/USP, Brazil) and the Academic Supercomputing Center in Kraków, Poland (Supercomputer Prometheus at ACK CYFRONET AGH). D.F.G. thanks the Brazilian agencies CNPq (no. 302949/2014-3) and FAPESP (no. 2013/10559-5) for financial support. A.L. acknowledges the NSF grant AST 1212096, NASA grant NNX14AJ53G, and a distinguished visitor PVE/CAPES appointment at the Physics Graduate Program of the Federal University of Rio Grande do Norte, the INCT INEspaço and Physics Graduate Program/UFRN. E.T.V. acknowledges the support of the AAS.

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