Nonthermal Particle Acceleration in 3D Relativistic Magnetic Reconnection in Pair Plasma

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Received 2017 May 15; accepted 2017 June 9; published 2017 July 10

Abstract

As a fundamental process converting magnetic to plasma energy in high-energy astrophysical plasmas, relativistic magnetic reconnection is a leading explanation for the acceleration of particles to the ultrarelativistic energies that are necessary to power nonthermal emission (especially X-rays and gamma-rays) in pulsar magnetospheres and pulsar wind nebulae, coronae and jets of accreting black holes, and gamma-ray bursts. An important objective of plasma astrophysics is therefore the characterization of nonthermal particle acceleration (NTPA) effected by reconnection. Reconnection-powered NTPA has been demonstrated over a wide range of physical conditions using large 2D kinetic simulations. However, its robustness in realistic 3D reconnection—in particular, whether the 3D relativistic drift-kink instability (RDKI) disrupts NTPA—has not been systematically investigated, although pioneering 3D simulations have observed NTPA in isolated cases. Here, we present the first comprehensive study of NTPA in 3D relativistic reconnection in collisionless electron–positron plasmas, characterizing NTPA as the strength of 3D effects is varied systematically via the length in the third dimension and the strength of the guide magnetic field. We find that, while the RDKI prominently perturbs 3D reconnecting current sheets, it does not suppress particle acceleration, even for zero guide field; fully 3D reconnection robustly and efficiently produces nonthermal power-law particle spectra closely resembling those obtained in 2D. This finding provides strong support for reconnection as the key mechanism powering high-energy flares in various astrophysical systems. We also show that strong guide fields significantly inhibit NTPA, slowing reconnection and limiting the energy available for plasma energization, yielding steeper and shorter power-law spectra.

Key words: acceleration of particles – galaxies: jets – gamma-ray burst: general – magnetic reconnection – pulsars: general – relativistic processes

1. Introduction

A long-standing puzzle in high-energy plasma astrophysics is the mechanism behind nonthermal particle acceleration (NTPA) that produces power-law X-ray and \( \gamma \)-ray spectra observed in pulsar wind nebulae (PWNe), coronae and jets of accreting black holes (BHs) in X-ray binaries (XRBs) and active galactic nuclei (AGNs) including blazars, and gamma-ray bursts (GRBs). A leading candidate is magnetic reconnection, a basic plasma process that rapidly converts magnetic into particle kinetic energy through magnetic field rearrangement and relaxation. In the high-energy universe, magnetic reconnection is often relativistic: the energy density of the reconnecting magnetic field \( B_0 \) exceeds that of the ambient plasma (including rest-mass), heating plasma to relativistic temperatures, driving relativistic flows, and accelerating particles to ultrarelativistic energies. Relativistic reconnection has been invoked in electron–ion or mixed-composition plasmas in accreting BH coronae (Beloborodov 2017), GRBs (Drenkhahn & Spruit 2002; Giannios 2008; McKinney & Uzdensky 2012), and blazar jets (Giannios et al. 2009; Nalewajko et al. 2011), but also, especially, in electron–positron pair plasmas in pulsar magnetospheres (Lyubarsky 1996; Uzdensky & Spitkovsky 2014; Cerutti et al. 2015), pulsar winds (Coroniti 1990; Arka & Dubus 2013), PWN (Uzdensky et al. 2011; Cerutti et al. 2012a, 2013, 2014a) and magnetar flares (Lyutikov 2003, 2006; Uzdensky 2011). Beyond its astrophysical applications, relativistic pair reconnection is important as the simplest reconnection scenario, a reference case for studying effects of plasma composition, collisions, radiation, etc.

Recently, first-principles particle-in-cell (PIC) kinetic simulations have significantly advanced understanding of NTPA in relativistic pair-plasma reconnection, mostly in two dimensions (Zenitani & Hoshino 2001, 2005, 2007; Jaroschek et al. 2004; Bessho & Bhattacharjee 2007; Lyubarsky & Liverts 2008; Liu et al. 2011, 2015; Cerutti et al. 2012b, 2013; Guo et al. 2014, 2015; Sironi & Spitkovsky 2014; Sironi et al. 2015, 2016; Werner et al. 2016), with a few 3D studies (Jaroschek et al. 2004; Zenitani & Hoshino 2008; Kagan et al. 2013; Cerutti et al. 2014b; Guo et al. 2014, 2015; Sironi & Spitkovsky 2014). Such studies are very challenging because even the qualitative character of reconnection dynamics, hence also NTPA, depends on the scale separation between the global system size \( L \) and the microscopic plasma scales, e.g., the average electron gyroradius in the reconnection layer, \( \bar{r}_e = \frac{\bar{v}_e e^2}{e B_0} \), where \( \bar{v}_e \) is the average relativistic energy of electrons energized by reconnection. When \( L/\bar{r}_e \) is large, as in astrophysical systems, reconnection is highly dynamic, with many secondary magnetic islands (plasmoids, or in 3D, flux ropes) and inter-plasmoid current sheets (Bhattacharjee et al. 2009; Uzdensky et al. 2010; Loureiro et al. 2012). The relevant large-system plasmoid-dominated regime, in which reconnection characteristics like the particle energy distribution \( f(\gamma) \) no longer depend on \( L \), is realized above a critical size \( L_c \sim 10^2 \bar{r}_e \) (Werner et al. 2016).

Only recently have 2D PIC simulations managed to probe this large-system regime systematically, confirming reconnection-powered NTPA (e.g., Guo et al. 2014, 2015; Sironi &...
and positron density $n_p$. Conforming to periodic boundary conditions, the magnetic field undergoes two reversals (supported by two initial current layers of half-thickness $\delta$), $B_4(y) = \pm B_0 \tanh \left[ (y - y_1)/\delta \right]$, where $y_1 = L_z/4$ and $y_2 = 3L_z/4$ are the midplanes of the two layers. A uniform guide field $B_z$ is added in the $z$-direction; we explore $B_z/B_0 \in \{0, 0.25, 0.5, 1, 2\}$.

To study magnetically dominated reconnection, we utilize the upstream plasma with the largest “hot” magnetization resolving the Debye length (described below), $c_0 = B_0^2/(4\pi\epsilon_0) = 25$ (where $h = 4\pi n_b m_e c^2$ in the relativistic limit; cf. Melzani et al. 2014). The “cold” magnetization, which sets the overall simulation scale and could be anything consistent with $c_0 = 25$ and $\theta_b \gg 1$, is $\sigma = B_0^2/(4\pi n_0 m_e c^2) = 46 n_0 c_0 = 2.75 \times 10^4$, as in Werner et al. (2017).

The two current layers are relativistic Harris sheets (Kirk & Skjæraasen 2003) with simulation-frame density $n_d(y) = n_{d0} \cosh^{-2}\left[ (y - y_1)/\delta \right]$. Within each sheet, electrons and positrons drift with opposite average velocities $\pm \beta^{\text{drift}} c^2$ to generate the current supporting the magnetic field reversal, i.e., satisfying Ampère’s law with $\delta = B_0/(4\pi \epsilon_0 \sigma^{\text{drift}})$. The Harris sheets also provide pressure balance, with rest-frame temperature $\theta_b = (1/2) \gamma^{\text{drift}} (n_b/n_{d0}) \sigma = 50 \gamma^{\text{drift}} (n_b/n_{d0}) \theta_d$, where $\gamma^{\text{drift}} = (1 - \beta^{\text{drift}})^{-1/2}$. Here, we use $\theta_b = 275$, $n_{d0}/n_b = 5$, and $\beta^{\text{drift}} = 0.3$; hence, $\gamma^{\text{drift}} \approx 1.05$, and $\theta_b = 2890$.

In nonrelativistic ultrarelativistic pair reconnection, lengths and times scale with $\sigma$; as long as $\theta_b \gg 1$, simulations with the same $c_0$ but different $\sigma$ are equivalent after rescaling. Thus, we define a characteristic length scale $\rho_c = 2 \sigma c_0 = 2m_e c^2/eB_0$, four times the gyroradius of a particle with energy $B_0^2/(8\pi n_0)^3$. The corresponding timescale is $\omega = c/\rho_c$. Convergence studies led us to choose $40$ total macroparticles per cell and grid cell size $\Delta x = \Delta y = \Delta z = \rho_c/24$, sufficient to resolve both the evolved and initial ($\delta = 8\Delta x = \rho_c/3$) current-layer thicknesses, and also the upstream Debye length ($\lambda_D = 1.2\lambda_\text{Debye}$), preventing unphysical finite-grid heating (Langdon 1970). We use the Courant–Friedrichs–Lewy time-step $\Delta t = 0.99\Delta x/(\sqrt{3} c)$ and typically run until $t \approx 5L_z/c$ (i.e., $t \approx 200/\omega$; for $L_z = 40\rho_c$). These parameters yield better than 1% energy conservation.

We show results mostly for the fiducial size $L_z = 40\rho_c = 80\sigma c_0$, with additional $L_z = 20\rho_c$ and $L_z = 64\rho_c$ simulations for $B_z = 0$. All systems have $L_y = 2L_z$ to prevent undesirable interaction between the two layers. By varying $L_z$ from a single cell (2D) up to 4$L_z$, we explore the 3D nature of reconnection.

As in many 2D reconnection studies, the initial magnetic field is perturbed uniformly in $z$, as described by the vector potential, e.g., for the lower layer

$$A_z = \left[ 1 + 0.01 \cos \left( \frac{2\pi x}{L_x} \right) \cos^2 \left( \frac{2\pi (y - y_1)}{L_y} \right) \right] \times B_0 \delta \left[ \ln \cosh \frac{y_2 - y_1}{2\delta} - \ln \cosh \frac{y - y_1}{\delta} \right].$$

In all our 2D simulations, and in 3D simulations with $B_z \geq 0.25B_0$, the presence of a perturbation does not change the reconnection rate or NTPA. In perturbationless 3D runs

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8 This definition agrees with Werner et al. (2017), but differs from $\rho_c \equiv \theta_d \rho_0$ used by Werner et al. (2016).
with $B_{gz} = 0$, however, particles from the dense initial Harris sheet avoid becoming trapped in flux ropes, and instead spread about the layer, increasing the immediately upstream density, effectively lowering $\sigma_h$, and significantly slowing reconnection and hindering NTPA. Leaving details for a future publication, we emphasize that initializing 3D $B_{gz} = 0$ simulations with a perturbation makes them more closely resemble (2D and 3D) $B_{gz} = B_0/4$ simulations.

3. Results

Our main findings are: (1) magnetic energy dissipation and nonthermal particle acceleration (NTPA) are largely unaffected by three-dimensionality, i.e., by $L_z/L_x$; and (2) a strong guide field reduces the rate and amount of magnetic dissipation and inhibits NTPA.

We illustrate these results with plots showing energy dissipation and final particle energy spectra in simulations with $L_z = 40\rho_0 = 80\sigma_h\rho_0$, for $L_z/L_x$ ranging from 0.001 (one cell in $z$) up to 1 (and up to $L_x/L_x = 2$ for $B_{gz} = B_0$ and 4 for $B_{gz} = 0$), and for $B_{gz}/B_0 \in \{0, 0.25, 0.5, 1, 2\}$, all with $\sigma_h = 25$.

In agreement with previous 3D studies (Zenitani & Hoshino 2008; Kagan et al. 2013; Cerutti et al. 2014b), we find that reconnection proceeds through rapid growth of tearing (plasmoid) and RDKI instabilities, quickly disrupting the initial current sheet and generating a complex, dynamic hierarchy of interacting flux ropes. Although RDKI has clearly visible effects (e.g., current-layer kinking in the $xz$ plane; Figure 1, left), magnetic dissipation is similar in 2D and 3D. Figure 1 (right) shows the evolution of transverse magnetic energy, $U_{mag,xy} \equiv \int (B_x^2 + B_y^2)/8\pi \, dV$, for different values of $B_{gz}$ and $L_z/L_x$. This evolution is typical for a closed system: reconnection ramps up gradually and continues until the available free magnetic energy is exhausted. The released $U_{mag,xy}$ is converted mostly into particle kinetic energy, with some into $U_{mag,z} \equiv \int B_z^2/8\pi \, dV$ and a small amount into the electric field (Table 1).

Overall, energy dissipation differs little between 2D and $L_z/L_x \gtrsim 1$; however, 3D simulations with weak guide field ($B_{gz} \lesssim B_0/4$) dissipate slightly more energy than their 2D counterparts by accessing final relaxed states that are nonuniform in $z$, e.g., slightly kinked flux ropes with lower magnetic energy than the straight (translationally symmetric in $z$) flux ropes that are the only possibility in 2D. In 2D or nearly 2D runs with $B_{gz} = 0$, $U_{mag,xy}$ fell by $\sim 40\%$ from initial to final state, while decreasing by $46\%$ for $L_z/L_x = 2$ (Figure 1, right). A modest $B_{gz} \gtrsim 0.5$, however, stabilizes the kink (for the explored range of $L_z/L_x \lesssim 2$), resulting in nearly identical 2D and 3D energy budgets.

While $L_z/L_x$ has little effect, we find that $B_{gz}$ significantly influences all aspects of reconnection. As expected, a strong guide field ($B_{gz} \gtrsim B_0$) suppresses RDKI (Zenitani &

Table 1

| $B_{gz}/B_0$ | $\sigma_{eff}$ | $L_z/L_x$ | $\Delta U_{mag,xy}/U_{mag,xy}(0)$ | $\Delta U_{mag,zy}/U_{mag,zy}(0)$ | $\Delta U_{mag,zx}/U_{mag,zx}(0)$ | $p$ |
|-----------|--------------|----------|-------------------------------|-------------------------------|-------------------------------|----|
| 0         | 25           | 2D       | $-39\%$                       | $\lesssim 0.02\%$            | 39\%                         | 1.9 ± 0.1 |
| 0         | 25           | 1        | $-46\%$                       | $\lesssim 0.3\%$            | 46\%                         | 2.0 ± 0.1 |
| 0.25      | 9.8          | 2D       | $-34\%$                       | 4\%                          | 30\%                         | 2.1 ± 0.1 |
| 0.25      | 9.8          | 1        | $-36\%$                       | 5\%                          | 31\%                         | 2.1 ± 0.1 |
| 0.5       | 3.4          | 2D       | $-26\%$                       | 6\%                          | 19\%                         | 2.1 ± 0.1 |
| 0.5       | 3.4          | 1        | $-26\%$                       | 6\%                          | 20\%                         | 2.2 ± 0.1 |
| 1         | 0.96         | 2D       | $-15\%$                       | 4\%                          | 11\%                         | 3.0 ± 0.5 |
| 1         | 0.96         | 1        | $-15\%$                       | 4\%                          | 11\%                         | 2.7 ± 0.2 |
| 2         | 0.25         | 2D       | $-8\%$                        | 1\%                          | 8\%                          | 3.3 ± 0.5 |
| 2         | 0.25         | 1        | $-8\%$                        | 1\%                          | 8\%                          | 3.3 ± 0.3 |

Notes. For each $B_{gz}$ (for both 2D and $L_z/L_x$), we show $\sigma_{eff}$ and the measured change in $U_{mag,xy}, U_{mag,zy},$ and particle kinetic energy (normalized to initial $U_{mag,xy}$), as well as the particle energy spectral index $p$ (the “error” expresses the fitting uncertainty within a single simulation).
Hoshino 2007, 2008). In addition, it slows down reconnection and reduces the total dissipated magnetic energy, equally in 2D and 3D. We attribute these effects to the guide field’s effective inertia and pressure, respectively (see also Dahlin et al. 2016); since the guide field is approximately frozen into the plasma, its magnetic enthalpy $h_{\text{mag},z} \equiv B_{\text{gz}}^2/4\pi$ lowers the effective $\sigma_p$-parameter:

$$\sigma_{p,\text{eff}} \equiv B_0^2/4\pi(h + h_{\text{mag},z}) = B_0^2/(4\pi h + B_{\text{gz}}^2),$$

which governs the relevant transverse Alfvén speed, $V_A = c \sigma_{\text{eff}}^{1/2} / (1 + \sigma_{p,\text{eff}})^{1/2} = c B_0 / (4\pi h + B_0^2 + B_{\text{gz}}^2)^{1/2}$. In the high-$\sigma$ regime, $B_0^2 \gg 4\pi h$, $V_A$ falls to $V_A / c \approx B_0 / (B_0^2 + B_{\text{gz}}^2)^{1/2}$, becoming subrelativistic for strong guide field ($B_{\text{gz}} \gg B_0$), thereby reducing the reconnection rate, $E_{\text{rec}} \sim 0.1B_0V_A/c$, to $\sim 0.1B_0^2 / B_{\text{gz}}$.

In addition, since a strong $B_{\text{gz}}$ makes the plasma less compressible, it also affects the final relaxed state (Uzdensky et al. 1996). In particular, the tension force of reconnected field lines that contracts plasmoids must perform work against the combined plasma and guide-field pressure; a strong $B_{\text{gz}}$ resists compression and makes the plasmoids stiffer, reducing the work that can be done and hence the overall amount of $U_{\text{mag},xy}$ released by reconnection. Moreover, since some of the released $U_{\text{mag},xy}$ is spent compressing $B_r$, the fraction of the energy that goes to particles is also reduced (Table 1).

Most of the released $U_{\text{mag},xy}$ goes to particles, producing a nonthermal power-law spectrum in both 2D and 3D. We will show particle energy spectra $f(\gamma) = dN/d\gamma$ at late times, when reconnection has ceased and the spectra stabilize, as illustrated in Figure 2, which shows nearly identical spectra from times $t \approx 3.2$ and $4.3 L_z/c$ for simulations with $B_{\text{gz}} = 0$ and $L_z = L_x$ (for three different $L_x$). For most of our large-system ($L_z = 40\rho_0$) runs, the displayed time of $t = 143/\omega_c = 3.6L_z/c$ is thus sufficiently late to capture final particle distributions.

Our simulations reach sufficiently large $L_x$ so that the spectral slopes become nearly independent of system size. This is illustrated in Figure 2, which shows $f(\gamma)$ for runs with $L_x/\rho_0 \in \{20, 40, 64\}$, $L_z/L_x = 1$, and $B_{\text{gz}} = 0$. While the power law for $L_x/\rho_0 = 20$ ($p \approx 2.3$) is slightly steeper than for $L_x/\rho_0 = 40$ and 64, the latter two runs have essentially the same $p \approx 2.0$. Similarly, we find for other $B_{\text{gz}}$ and $L_z$ that $p$ converges with $L_x$, within measurement error.

However, the high-energy cutoff $\gamma_c$ of the nonthermal power law still increases with system size (Figure 2). This is expected from our previous cutoff investigation in large 2D systems (Werner et al. 2016) employing a similar, but not identical setup, which showed that, while $p$ converges at modest $L_x$, similar to the present study, the rise of $\gamma_c$ with $L_x$ finally saturates only at $L_x/\sigma\rho_0 \sim 200$, well above the sizes ($L_x/\sigma\rho_0 = 2L_x/\rho_0 = 80 - 128$) accessible to our present 3D simulations. We thus focus here on $p$ and leave determination of $\gamma_c$ to future studies. Nevertheless, we note that our simulations achieve cutoff energies (in both 2D and 3D) around $\gamma_c \sim 4-8\sigma$, comparable (despite the different setup) to the asymptotic, large-$L_x$ cutoff found in Werner et al. (2016), and to those in our 2D study of relativistic electron–ion reconnection (Werner et al. 2017). This indicates that 3D reconnection can indeed accelerate particles to about the same energies as 2D.

We now characterize the effects of three-dimensionality and guide field on NTPA, comparing simulations with the same size $L_x = 40\rho_0 \approx 80\sigma\rho_0$ but different $L_z/L_x$ and $B_{\text{gz}}$. Our main finding is that, despite fundamental differences in 2D and 3D dynamics (Figure 1, left), particle acceleration is nearly unchanged by 3D effects, regardless of guide field (Figure 3). Importantly, 3D relativistic reconnection with weak or no guide field is an efficient particle accelerator, consistent with single simulations run by Sironi & Spitkovsky (2014) and Guo et al. (2014, 2015), but contradicting Zenitani & Hoshino (2007, 2008). In addition, we also note that a sufficiently strong $B_{\text{gz}}$ hinders NTPA in 2D and 3D, reducing the number and total energy of accelerated particles, and steepening the power law (see Figure 4, left). The effects of three-dimensionality and guide field are summarized in Figure 4 (right), comparing 2D and 3D ($L_z = L_x$) simulations. We conjecture that the lower reconnection rate (hence weaker accelerating electric field) and the smaller available energy budget (Table 1) both contribute to the relative inefficiency and limited range of NTPA for strong guide fields, $B_{\text{gz}} > B_0$. We further propose that $p$ steepens with $B_{\text{gz}}$ because $\sigma_{p,\text{eff}} \sim B_0^2 / B_{\text{gz}}^2$ decreases (Equation (2)), and lowering $\sigma_p$ has been seen to increase $p$ in $B_{\text{gz}} = 0$ 2D studies; Guo et al. 2014; Sironi & Spitkovsky 2014; Werner et al. 2016). Intriguingly, while it does not yield $p \sim 1$ for $\sigma_{p,\text{eff}} \rightarrow \infty$ as those 2D studies suggest, the empirical fit $p \approx 1 + 0.7\sigma_p^{-1/2}$, found for $B_{\text{gz}} = 0$ semirelativistic electron–ion reconnection (Werner et al. 2017), captures the $B_{\text{gz}}$-dependence found here.
4. Conclusions

This Letter presents a systematic first-principles confirmation via PIC simulation that 3D relativistic reconnection in pair plasmas can robustly and efficiently drive nonthermal particle acceleration (NTPA), yielding unambiguous power-law particle distributions despite the presence of RDKI. In addition, the strong influence of guide field on the NTPA power-law index \( p \) is described through the dependence of \( p \) (see Werner et al. 2017) on the effective hot magnetization \( h_{\text{eff}} \) including the enthalpy of the guide field. This study thus resolves the longstanding controversy regarding the effects of 3D RDKI structures and guide field on NTPA (Jaroschek et al. 2004; Zenitani & Hoshino 2007, 2008; Guo et al. 2014; Sironi & Spitkovsky 2014). Importantly, our results show that 2D studies are in fact pertinent to 3D reconnection; intriguingly, however, the poorer particle trapping within 3D flux ropes (Dahlin et al. 2015) may allow 3D NTPA to extend beyond the energy cutoff in 2D reconnection (Werner et al. 2016). Our findings lend strong support to reconnection-based models of, e.g., rapid \( \gamma \)-ray flares in the Crab PWN and AGN/blazar jets, prompt GRB emission, and hard-X-ray emission in accreting BH binaries (cf. Section 1). Furthermore, these results lay the foundation for characterizing NTPA in realistic 3D reconnection as a function of ambient plasma, helping to diagnose plasma conditions in remote astrophysical systems using observed radiation spectra.

We thank Vladimir Zhdankin and Mitch Begelman for helpful discussions. This work was supported by grants DE-SC0008409, DE-SC0008655 (DOE); NNX12AP17G, NNX16AB28G (NASA); and AST-1411879 (NSF). D.A.U. gratefully acknowledges the hospitality of the Institute for Advanced Study and support from the Ambrose Monell Foundation. Simulations were possible thanks to an INCITE award at the Argonne Leadership Computing Facility, a DOE User Facility supported under Contract DE-AC02-06CH11357. Some 2D simulations were run using XSEDE resources (Towns et al. 2014) supported by NSF grant ACI-1548562, as well as resources provided by the NASA High-End Computing Program through the NASA Advanced Supercomputing Division at Ames Research Center.

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