Radiative kinetic simulations of steady-state relativistic plasmoid magnetic reconnection

José Ortuño-Macías & Krzysztof Nalewajko
Nicolaus Copernicus Astronomical Center, Polish Academy of Sciences, Bartycka 18, 00-716 Warsaw, Poland
jortuno@camk.edu.pl

ABSTRACT
We present the results of 2D particle-in-cell (PIC) simulations of relativistic magnetic reconnection (RMR) in electron-positron plasma, including the dynamical influence of the synchrotron radiation process, and integrating the observable emission signatures. The simulations are initiated with a single Harris current layer with a central gap that triggers the RMR process. We achieve a steady-state reconnection with unrestricted outflows by means of open boundary conditions. The radiative cooling efficiency is regulated by the choice of initial plasma temperature $\Theta$. We explore different values of $\Theta$ and of the background magnetisation $\sigma_0$. Throughout the simulations, plasmoids are generated in the central region of the layer, and they evolve at different rates, achieving a wide range of sizes. The gaps between plasmoids are filled by smooth relativistic outflows called minijets, whose contribution to the observed radiation is very limited due to their low particle densities. Small-sized plasmoids are rapidly accelerated, however, they have lower contributions to the observed emission, despite stronger relativistic beaming. Large-sized plasmoids are slow, but produce most of the observed synchrotron emission, with major part of their radiation produced within the central cores, the density of which is enhanced by radiative cooling. Synchrotron lightcurves show rapid bright flares that can be identified as originating from tail-on mergers between small/fast plasmoids and large/slow targets.

Key words: acceleration of particles – magnetic reconnection – methods: numerical – plasmas – relativistic processes

1 INTRODUCTION
Signatures of non-thermal particle energy distributions are commonly observed in the high-energy astrophysical phenomena such as gamma-ray bursts (e.g., Piran 2004), pulsar wind nebulae (e.g. Buehler et al. 2012) and blazars (e.g., Madejski & Sikora 2016). What these diverse sources have in common is very broad photon spectra indicating broad non-thermal energy distributions of the radiating particles.

Blazars are a subclass of active galactic nuclei (AGN), in which a relativistic jet emerging from the supermassive black hole (SMBH) is aligned with our line of sight (Urry & Padovani 1995). They are extremely luminous sources of radiation that spans the entire electromagnetic spectrum, from radio up to very high energy $\gamma$-rays ($\sim$ TeV). The extreme broadness and power-law appearance of the blazar spectra are signatures of an efficient non-thermal particle acceleration (NTPA) process. Two major spectral components are typically observed with regular characteristics (Fossati et al. 1998; Ghisellini et al. 2017), with the low-energy one (radio to optical/UV/X-rays) interpreted as synchrotron emission, and the high-energy one (mainly $\gamma$-rays) due to either leptonic (inverse Compton) or hadronic processes (e.g., Sikora et al. 2009; Böttcher et al. 2013).

The emission of blazars is also characterised by strong variability on timescales ranging from decades (e.g., Ahnen et al. 2016; Goyal et al. 2017) to minutes (e.g., Aharonian et al. 2007; Albert et al. 2007; Meyer et al. 2019). The shortest variability timescales are much shorter than the light crossing time ($\sim$ hours) of the gravitational radius of the SMBH contained in these type of AGN ($M_{BH} \sim 10^9 M_\odot$). Such short variability time scales and high luminosities require very compact emitting regions within the jet (Levinson 2007; Begelman et al. 2008; Katarzyński et al. 2008).

Compact regions containing relativistic particles and magnetic fields have been often invoked to explain the multi-wavelength emission of blazars (e.g., Kirk et al. 1998; Chiaberge & Ghisellini 1999). Theoretical models for the formation of collimated jets with relativistic velocities require relativistically magnetised environments\(^1\) (Blandford & Znajek 1977). In general, magnetisation parameter is defined as $\sigma = \frac{B^2}{8\pi n_e c}$.

\(^1\) In general, magnetisation parameter is defined as $\sigma =$...
Magnetic reconnection is a dissipation mechanism that is most efficient in strongly magnetised plasmas. During the reconnection, free magnetic energy is transferred to the thermal heating and bulk acceleration of plasma, as well as to the NTPA process (Zweibel & Yamada 2009). Numerous arguments point to the RMR as the process responsible for the multi-wavelength and multi-timescale variability emission of blazars (Sironi et al. 2015). A basic requirement for reconnection to occur in relativistic jets is to have local inversions of the magnetic field lines, which may be caused by internal jet instabilities, in particular the current-driven kink modes (Begelman 1998; Giannios & Spruit 2006; Alves et al. 2018; Bromberg et al. 2019), or by global magnetic polarity reversals (Lovecchio et al. 1997; Giannios & Uzdensky 2019).

Traditional models of magnetic reconnection (Sweet 1958; Parker 1957; Petscheck 1964) have been adapted to the relativistic regime by Lutynik & Uzdensky (2003) and Lyubarsky (2005). The latter work became a basis for the minijets model of relativistic bulk outflows driven by the RMR (Giannios et al. 2009; Nalewajko et al. 2011). Other recent works emphasised the fundamental role of plasmoids (or magnetic flux tubes) that arise spontaneously in sufficiently relativistic jet instabilities, in particular the current-driven kink modes (Begelman 1998; Giannios & Uzdensky 2019).

The NTPA mechanism in the case of RMR has been studied extensively by means of kinetic particle-in-cell (PIC) numerical simulations. In the most basic case of electron-positron pair plasma, it has been established that NTPA produces power-law energy distributions of particles \( N(E) \propto E^{-p} \) with localised power-law indices as hard as \( p \approx 1 \) in the limit of \( \gamma \gg 1 \) (Zenitani & Hoshino 2001; Sironi & Spitkovsky 2014; Guo et al. 2014; Werner et al. 2016), converging however to \( p \approx 2 \) in the long term (Petropoulou & Sironi 2018).

Most PIC simulations of the RMR are initiated from Harris-type current layers contained in periodic domains. In the relatively small domains, this leads to the cancellation of reconnection outflows momenta and to the formation of artificially large single plasmoids. An alternative approach is to use open boundaries that allow outgoing particles to escape freely and absorb the outgoing Poynting flux (Daughton et al. 2006). This approach allows to investigate magnetic reconnection as a sustained steady-state process with unimpeded outflows approaching the terminal Alven velocity and continuous generation and evolution of plasmoids (Daughton & Karimabadi 2007).

The open-boundary approach has been first applied to the case of RMR by Sironi et al. (2016), who investigated the statistical properties of plasmoids, determined their size distribution as a power-law, and demonstrated an anti-correlation between plasmoid growth and bulk acceleration. The corresponding predictions for the variability of non-thermal radiation were applied to explain the observational characteristics of blazar flares (Petropoulou et al. 2016). Based on the plasmoid scaling laws, a stochastic model for the evolution of plasmoid chains was developed by Petropoulou et al. (2018). The results of these PIC simulations were also post-processed to calculate light curves of synchrotron and inverse Compton emission as would be observed if the reconnection region was located in a relativistic jet (Christie et al. 2019a,b). A semi-analytical model of broad-band emission from reconnection plasmoids in the context of blazar flares was also developed by Morris et al. (2019).

Other numerical studies of the RMR were performed in the context of \( \gamma \)-ray flares from the Crab Nebula, taking into account the effect of radiation reaction on individual particles (e.g., Cerutti et al. 2013, 2014). Recent PIC simulations which included the effects of synchrotron and inverse Compton cooling have shown that the particle acceleration can be decreased in comparison with previous non-radiative results (Nalewajko 2018; Schoeffler et al. 2019; Hakobyan et al. 2019; Sironi & Beloborodov 2019). We may therefore expect that evolution of plasmoids will be affected by radiative cooling.

In this work, we present the results of PIC simulations of the steady state RMR process allowed by the open boundaries under strong radiative energy losses due to the synchrotron radiation reaction. We investigate in detail the evolution of individual plasmoids under different levels of radiative cooling, and calculate accurate lightcurves of synchrotron radiation. We demonstrate a connection between rapid radiation flares and tail-on mergers of small/fast and large/slow plasmoids.

Plan of this work: §2 simulations setup; §3 analysis methods; §4.1 initial sequence; §4.2 evolved reconnection layer; §4.3 spacetime diagrams of the current layer; §4.4 particle density and mean energy distributions; §4.5 analysis of individual plasmoids; §4.6 acceleration and cooling of selected individual particles; §4.7 synchrotron light curves; §4.8 global energy conservation; §4.9 energy distributions of particles and photons; §5 discussion; §6 conclusions.

2 SIMULATION SETUP

We make use of a custom version of the Particle-in-Cell (PIC) code Zeltron (Cerutti et al. 2013). We perform 2D simulations of relativistically magnetised pair plasma in a fixed tall domain of dimensions \( L_x \times L_y \) with \( L_y = 4L_x \), within the coordinate ranges \( 0 < x < L_x \) and \( -L_y/2 < y < L_y/2 \).

Our simulations are initiated from an equilibrium configuration \( B(t = 0) = B_{ini} \) and \( E(t = 0) = E_{ini} = 0 \) that involves a single Harris-type current layer placed in the middle of the computational domain:

\[
B_{ini,x} = -B_0 \tanh \left( \frac{y}{\delta} \right),
\]

where \( B_0 \) is the characteristic value of magnetic field strength, and \( \delta \) is the Harris layer half-thickness. We include no initial guide field, hence \( B_{ini,x} = 0 \). The magnetic
field gradient across the Harris layer is supported by the electric current and pressure provided by the drifting population of particles that is characterised by the Maxwell-Jüttner energy distribution with dimensionless temperature $\Theta_d = k_B T_d/(mc^2)$, Lorentz-boosted with the drift velocity $\beta_d = u_d/c = 0.3$, and number density
\[ n_{a0}(y) = n_{a0} \cosh^{-2}(y/\delta), \]
where $n_{a0} = \gamma_a B_0^2/(4\pi \Theta_d m c^2)$, and $\gamma_a = (1 - \beta_d^2)^{-1/2}$ is the drift Lorentz factor. The Harris layer thickness is related to the nominal plasma gyroradius $\rho_0 = \Theta_d m c^2/(e B_0)$ as $\delta = 2\rho_0/(\gamma_a \beta_d)$ (Kirk & Skjæraasen 2003).

The current layer is immersed in a background plasma the Maxwell-Jüttner energy distribution with dimensionless temperature $\Theta_b = k_B T_b/(mc^2) = \Theta_d$, and number density $n_b$ that is determined by the assumed value of the cold magnetisation $s_0 = B_0^2/(4\pi n_m m c^2) \gg 1$. The corresponding initial density contrast can be expressed as $n_{a0}/n_b = \gamma_a s_0/(2\Theta_b)$.

Our choice of drift velocity means that the current layer is relatively thick, and hence stable to the tearing modes. We experimented with higher values $\beta_d \lesssim 0.5$, which results in spontaneous formation of slowly evolving plasmoids. In order to speed up the evolution of the current layer, we follow Sironi et al. (2016) in triggering fast magnetic reconnection by placing a narrow gap ($\Delta x = 8\delta$) in the middle of the current layer. Within the trigger gap, the distribution of drifting particles is initialised with temperature reduced by factor 10, and with no drift velocity.

The dimensions of our computational domain are up to $N_x \times N_y = 4608 \times 18432$ cells with equal spatial resolution for both dimensions $dx = dy = \rho_0/3$, which results in the physical scales of $L_x \times L_y = (1536 \times 6144)\rho_0$. The temporal resolution is $dt = 0.9 dt_{CFL}$, where $dt_{CFL} = [1/(dx)^2 + 1/(dy)^2]^{-1/2}/c$ is the Courant-Friedrichs-Lewy time step. The initial distributions of background and drifting electrons/positrons are represented on average by 64 macroparticles per cell per species.

Our computational domain is open at left/right boundaries, where outgoing particles are removed, and fresh particles are injected at every timestep at the fixed rate calculated to maintain the initial number density of both background and drifting particles. Just within the left/right boundaries we place the field-absorbing layers of thickness $\Delta_{abs} = 30dx$. In the absorbing layers, in addition to the standard time advance of magnetic and electric fields, at every time step we perform the following operation:

\[ B(x) \rightarrow B(x) + \lambda(x) \left[ B_{ini}(x) - B(x) \right], \]
\[ E(x) \rightarrow E(x) + \lambda(x) \left[ E_{ini}(x) - E(x) \right], \]
where $\lambda(x) = 0.5 |x - x_{abs}|/\Delta_{abs}^3$, and $x_{abs}$ is the position of the absorbing layer inner edge. For the top/bottom boundaries, we apply periodic conditions for the particles and fields, while only the $z$-component of the electric field is absorbed using the above rule.

We apply the synchrotron radiation reaction to every particle at every timestep\(^3\), following the prescription of Cerutti et al. (2013):

\[ \frac{d\bar{u}}{dt} = \frac{P_{syn}}{\gamma m c^2}, \quad P_{syn} = \frac{\sigma_T c^2}{4\pi} [\langle \gamma E + u \times B \rangle^2 - \langle E \cdot u \rangle^2], \]

where $\sigma_T$ is the Thomson cross section, and $u = \gamma \beta c$ is the dimensionless momentum of a particle with dimensionless velocity $\beta = \bar{u}/c$ and Lorentz factor $\gamma = (1 - \beta^2)^{-1/2} = (1 + \bar{u}^2)^{1/2}$. Noting that for $\gamma \gg 1$ we have $u \approx \gamma$, the radiative cooling rate can be estimated as:

\[ \frac{\partial \gamma}{\partial t} = \frac{u}{\gamma} \frac{\partial u}{\partial t} \approx -\frac{P_{syn}}{\gamma m c^2}. \]

In the limit of zero electric field and isotropic particle distribution, Eq. (5) reduces to $P_{syn} = (4/3) \pi e c^3 \bar{u}^2 B_{0,0}$, where $B_{0,0} = B_0^2/8\pi$ is the initial background magnetic energy density. Taking into account that for the Maxwell-Jüttner distribution we have $\langle \gamma \rangle \approx 3\Theta$ and $\langle \gamma^2 \rangle \approx 12\Theta^2$, we define the nominal cooling length as (Nalewajko et al. 2018):

\[ l_{cool} = c t_{cool} = \langle \gamma \rangle / \langle |d\gamma/dt| \rangle \approx \langle \gamma \rangle \langle 3m_e c^2 \rangle / \langle 4\pi T \rangle B_{0,0} \approx (3\pi/2)e \sigma_T \Theta^2 B_0. \]

Provided that the temperature $\Theta$ is high enough, the cooling length scale will be contained within the domain size, $l_{cool} < L_x$, which allows us to study the dynamical effects of radiation. The total synchrotron power emitted in all directions per volume element of the initial background plasma is given by:

\[ P_{syn, b, 0} = \frac{dE_{syn}}{dx dy dt} = 16\pi e c^2 \gamma_0^2 B_{0,b,0}. \]

We also collect the synchrotron radiation spectra that would be measured by observers placed at the left and right side of the simulation domain. The total spectrum is calculated as the sum of contributions from all macroparticles present in the domain (see Nalewajko et al. 2018, and references therein).

\[ L_{syn}(\nu) = \sqrt{\frac{3c}{\pi}} \sum_{\epsilon^+ \epsilon^-} N_{e,1} F(\xi) \Omega_1, \]

where $F(\xi) = \int_0^\infty x K_{2\nu}(2x)dx$, $\xi = (4\pi\nu)/(3\gamma^2 \Omega_1)$, $\Omega_1 = (e/m_e c) \langle |E + n \times B| \times n \rangle$, $n = \bar{u}/|\bar{u}|$, $N_{e,1}$ is the number of charged particles (either electrons or positrons) represented by a single macroparticle, and $K_{5/3}$ is the modified second-kind Bessel function of order 5/3. The nominal synchrotron frequency of the initial background plasma is derived from the relation $\nu = \nu_{syn,b} = 1$, substituting $\Omega_1 \sim c e B_0/m_e c \equiv \omega_0 = \epsilon/\rho_0$ and $\gamma^2 = 12\Theta^2$:

\[ \nu_{syn,b} = (9/\pi) \Theta^3 \omega_0. \]

We performed several large simulations for different values of background magnetisation $s_0$ and initial particle distribution temperature $\Theta_b = \Theta_d$. In Table 1, we report these parameters and the corresponding cooling lengths $l_{cool}$.

### 3 ANALYSIS METHODS

#### 3.1 Plasmoids identification

We describe here an automated procedure that we use for identifying individual plasmoids along the reconnection layer. During the simulations, we frequently save one-

\[^3\text{We apply a restriction such that particles cannot lose more than half of their energy or reverse their momentum over a single time step.}\]
The nominal cooling length is calculated from Eq. (7).

| name  | σ0  | Θβ0 = Θγ | lcool/p0 |
|-------|-----|----------|---------|
| s10T1 | 10  | 2 · 10^5 | 2837    |
| s10Tm | 10  | 5 · 10^5 | 454     |
| s10Th | 10  | 1.25 · 10^6 | 73     |
| s50Tm | 50  | 5 · 10^5 | 454     |

3.2 Local reference frames

We convert certain parameters to their values in the local reference frames. As discussed in Werner et al. (2018), in relativistically hot fluid there are two alternative ways to define a ‘co-moving’ reference frame: (1) zero-current (Eckart) frame \( O' \) of velocity \( \beta_1 = ⟨β⟩ \); and (2) zero-momentum (Landau) frame \( O'' \) of velocity \( \beta_2 = ⟨γβ⟩/⟨γ⟩ \). These averages are calculated locally for all particles (electrons and positrons) in a given grid cell as \( ⟨a⟩ = (\sum_i a_i n_i)/(\sum_i n_i) \), where \( n_i \) are the multiplicities of individual macroparticles such that \( \sum n_i ≡ n \) is the local particle number density. For example, we calculate the electric field component in the Eckart frame as:

\[ E'_z = Γ_1 (E_z + ⟨β_2⟩ B_y - ⟨β_y⟩ B_z) \]

where \( Γ_1 = 1 - (⟨β_x⟩^2 - ⟨β_y⟩^2)^{-1/2} \) is the Eckart bulk Lorentz factor. On the other hand, we calculate the mean particle energy in the Landau frame as \( ⟨γ''⟩ = ⟨γ⟩/Γ_2 \), where \( Γ_2 = (1 - β_{z,2}^2 - β_{y,2}^2)^{-1/2} \) is the Landau bulk Lorentz factor.

4 RESULTS

4.1 Initial sequence

Figure 2 shows the initial evolution of the reconnection layer in our reference simulation s10Tm. Each panel shows a particle number density map and magnetic field lines for a central region of our computational domain (only a 1/64 fraction of its total vertical extent), in which the current layer is contained. A centrally located gap in the current layer expands rapidly sideways towards the left/right boundaries, the initial drifting plasma component is pulled towards the boundaries by the tension of closed magnetic field lines. Magnetic reconnection is triggered in the central low-density region, sustained by a thin current layer that is unstable to tearing modes which lead to the emergence of plasmoids.

At \( ct ≃ L_x \), the swept-up fronts of the initial drifting plasma leave the left/right boundaries. From that moment on, we can describe the simulated reconnection process as steady-state, with newer plasmoids continuously generated around the centre of the layer, and older plasmoids escaping through the left/right boundaries. The upper panels of Figure 2 demonstrate how our implementation of open boundaries works when large dense structures move across them.

We also find that the initial trigger gap induces a quasi-circular wave that propagates in all directions. The
Radiative simulations of steady-state plasmoid reconnection

Figure 2. Initial sequence of simulation $s10Tm$, presenting the logarithm of particle number density $n/n_b$ (see the top panel of Figure 3 for the colour scale) and the magnetic field lines (solid white lines). The dashed magenta lines indicate the limits of the field-absorbing boundary layers.

Horizontal wave fronts are effectively absorbed by the left/right boundaries, while the vertical wave fronts reach the top/bottom boundaries by $ct \approx 2L_x$ and are partially absorbed there, some very weak reflections return to the reconnection layer by about $ct \approx 4L_x$ without a significant effect.

4.2 Evolved reconnection layer

Figure 3 presents a selected evolved state of our reference simulation $s10Tm$ illustrated with multiple parameters. In this snapshot centred at the reconnection midplane, we observe diverse substructures. The main current layer is observed at $0.34 < x/L_x < 0.49$, characterised by high values of the specific electric current $|j_z|/j_{\text{max}}$ (electric current density $j_z$ close to its maximum value $j_{\text{max}} = cen$) despite low particle number density $n$, low bulk velocity $v_x$, and strong electric field $E'_z$ (except for the two minor plasmoids located at $x \approx 0.39L_x$ and $x \approx 0.425L_x$).

To the left of the main current layer, behind a medium-sized plasmoid at $x \approx 0.3L_x$, we find a relativistically fast reconnection outflow (strongly negative $v_x$ for $x < 0.25L_x$; $\langle \gamma \rangle \gg (\gamma''$) due to Lorentz transformation), also characterised by moderate particle density and intrinsic temperature $\langle \gamma'' \rangle /\Theta$, specific electric current $|j_z|/j_{\text{max}}$ decreasing systematically with distance, very weak intrinsic electric field $E'_z$, and weak synchrotron emission. This is a structure that has all characteristics of minijets – regular reconnection outflows of relativistic bulk velocity (Giannios et al. 2009; Nalewajko et al. 2011). We can easily recognise the conical geometry of the outflow region with parallel outflow velocity field (very low values of $v_y$), and oblique magnetic field lines crossing the outflow boundaries, as has been described by an analytical model of relativistic Petschek-type reconnection by Lyubarsky (2005). There is one qualitative difference from that model – the magnetic field lines in the outflow region are not vertical, the magnetic field gradient $\partial B_y/\partial y$ does not vanish in that region and it is supported by the non-zero electric current density $j_z$. We note that there is a roughly uniform vertical inflow of background plasma into the minijet region with velocity $v_y$ (reconnection rate) of the same order as that of the inflow into the main current layer.

To the right of the main current layer, we find a group of several plasmoids of various sizes, all propagating to the right at different velocities $v_x > 0$. The largest plasmoid can be seen centred at $x \approx 0.725L_x$, it is clearly slower than its smaller neighbour centred at $x \approx 0.6L_x$. The smaller plasmoid is in the process of merging with the large one, even though they both propagate in the same direction. This is a natural consequence of the inverse relation between the growth and bulk acceleration of plasmoids that has been first noticed by Sironi et al. (2016).

4.3 Spacetime diagrams

Most of the information on evolution of current layers, and especially on the plasmoids, is contained along the reconnection midplane, this information can be presented very efficiently in the form of spacetime diagrams. Following the practice of Nalewajko et al. (2015), the one-dimensional parameter $x$-profiles described in Section 3.1 are combined into spacetime diagrams of high time resolution. Figure 4 shows...
Figure 3. Selected evolved state of simulation s10Tm (the same as shown in Figure 1). From the top, the panels show: (1) particle number density $n/n_b$ in log scale, (2) out-of-plane electric current per particle number density $j_z/\langle \text{cen} \rangle$, (3) Landau velocity component along the current layer $v_x/c$, (4) Landau velocity component across the current layer $v_y/c$, (5) mean particle Lorentz factor measured in the simulation frame $\langle \gamma \rangle /\Theta$, (6) mean particle Lorentz factor measured in the Landau co-moving frame $\langle \gamma' \rangle /\Theta$, (7) out-of-plane electric field component measured in the Eckart co-moving frame $E_z'$, (8) total synchrotron power $E_{\text{syn}}$ in log scale. The solid white lines show the magnetic field lines, and the dashed magenta lines indicate the limits of the field-absorbing boundary layers.
spacetime diagrams for several parameters for our reference simulation s10Th.

The spacetime diagrams reveal a sustained bifurcating outflow along the $x$ coordinate (regions of negative and positive velocity component $\beta_x$) and a variety of plasmoids (indicated by enhanced plasma density and sharp positive gradients of magnetic field component $dB_y/dx > 0$), most of which are generated along the $\beta_x \approx 0$ line. There are a few large plasmoids that evolve very slowly, their bulk velocities are non-relativistic, acceleration time scale is long, and hence they spend more than the light-crossing time $L_x/c$ before leaving the simulation domain. On the other hand, there are many small plasmoids that accelerate quickly to relativistic bulk velocities, and they spend less than $L_x/c$ in the simulation domain. Small plasmoids may either escape through the boundaries, or merge with a larger plasmoid. As the velocity field is generally divergent, plasmoid mergers are relatively rare, typically they involve plasmoids of different sizes moving in the same direction (tail-on mergers). A large plasmoid can attract nearby small plasmoids, and can even reverse their motion, e.g., Figure 3 shows that at $ct/L_x = 3.36$ a small plasmoid centred at $x \approx 0.85L_x$ has $\beta_x > 0$, however, the spacetime diagrams reveal that it will merge with the large plasmoid located on its left side by $ct/L_x \approx 3.6$.

The regions between plasmoids are characterised by roughly uniform electric field component $E_x \approx 0.1B_0$ and by two-value magnetic field component $B_y \approx \pm 0.15B_0$ (positive where $\beta_x < 0$ and negative where $\beta_x > 0$). The uniformity of electric field measured in the simulation frame is consistent with the uniform reconnection rate indicated in the $\beta_y$ panel of Figure 3. It is remarkable given the non-uniformity of electric field measured in local Eckart frames $E'_x$, as shown in another panel of Figure 3. This indicates a smooth connection between the minijet outflows (where $E'_x$ is close to zero) and the proper magnetic diffusion areas (where $E'_x$ is strong).

High density regions of the plasmoids – essentially the plasmoid cores – are initially characterised by enhanced temperatures (mean particle energies measured in the local Landau frames reaching values of $\langle \gamma'' \rangle \approx 10\Theta$) due to the heating of particles in the reconnection process. Our spacetime diagram of $\langle \gamma'' \rangle / \Theta$ demonstrates that the cores of large plasmoids cool down gradually.

In Figure 5, we compare the spacetime diagrams of $\langle \gamma'' \rangle / \Theta$ for three simulations with different nominal cooling lengths $l_{\text{cool}}$. In the simulation s10TL1, in which $l_{\text{cool}} > L_x$, we find that plasmoid temperatures are the highest, showing no signs of cooling over the light-crossing time scale $L_x/c$. On the other hand, in the simulation s10TLh, the plasmoid temperatures are the lowest, they eventually become even lower than for the plasma between the plasmoids.

4.4 Plasma density and temperature distributions

The left panel of Figure 6 compares the distributions of particle number density $n/n_{\text{th}}$ based on the spacetime diagrams of 3 simulations with $\sigma_0 = 10$ and with different efficiencies of radiative cooling. The distributions form power-law tails with the slope of $-2.3$, independent of the cooling efficiency. However, the highest values of the particle density increase with stronger cooling, from $n_{\text{max}} \approx 150n_{\text{th}}$ for the simulation s10TL1 to $n_{\text{max}} \approx 1500n_{\text{th}}$ for the simulation s10Th. The high particle densities are realised in the cores of large plasmoids. Strong radiative cooling removes the gas pressure support, allowing the plasmoid cores to contract further.

The right panel of Figure 6 compares the distributions of mean particle energy $\langle \gamma'' \rangle / \Theta$ measured in the local Landau frames, also extracted from the spacetime diagrams of simulations with $\sigma_0 = 10$. This shows that the distributions are strongly peaked at values that decrease slightly with increasing cooling efficiency, from $\langle \gamma'' \rangle_{\text{peak}} \approx 5.49$ for the simulation s10TL1 to $\langle \gamma'' \rangle_{\text{peak}} \approx 4.8\Theta$ for the simulation s10Th.

4.5 Individual plasmoids

Figures 7, 8 and 9 present the histories of individual plasmoids compared for the four large simulations listed in Table 1.

We present separately the total plasmoid widths $w$ (dominated by the widths of plasmoid layers) and the widths of their cores $w_c$, noting that there is no correlation between them. In all simulations, the plasmoid core widths, after a brief formation phase, show a decreasing trend. On the other hand, the total plasmoid widths grow systematically in time for every simulation. The widths of large plasmoids show occasional dips that correspond to their mergers with smaller plasmoids, during which the positions of the minima of magnetic potential $A_z$ (that define the plasmoid boundaries) shift temporarily inwards.

We find that the densities of plasmoid cores $n_c$ grow systematically in time for all simulations. The larger the plasmoid, the denser its core becomes. For simulations with $\sigma_0 = 10$, higher core densities are reached for higher radiative efficiencies. In the case of $\sigma_0 = 50$, the core densities are even higher, roughly in proportion to $\sigma_0$.

The histories of plasmoid core velocities $\beta_{x,c}$ confirm the picture discussed before of small plasmoids being accelerated very rapidly to relativistic velocities and of large plasmoids being accelerated slowly only to mildly relativistic velocities.

The mean particle energies of plasmoid cores, measured in the simulation frame, are somewhat higher for small plasmoids. However, when measured in the Landau frame of the core, they are instead higher for large plasmoids until the radiative cooling effects become significant. The difference is mainly due to relativistic bulk motions of the small plasmoids. In particular, in the simulations s10TL1 and s10Th characterised by weak/moderate cooling, small plasmoids have $\langle \gamma'' \rangle_c \approx 6\Theta$, roughly constant in time, while large plasmoids can reach $\langle \gamma'' \rangle_c \approx 11\Theta$ in the case s10TL1. In the simulation s10Th characterised by strong cooling, the intrinsic mean particle energies of the core are reduced down to $\langle \gamma'' \rangle_c \approx 2\Theta$ by $L_x/c$. In the simulation s50Th, the cores of large plasmoids are heated up to $\langle \gamma'' \rangle_c \approx 23\Theta$ before cooling down to $\langle \gamma'' \rangle_c \lesssim 10\Theta$.

In contrast, the plasmoid layers are characterised by similar and stable intrinsic temperatures $\langle \gamma'' \rangle_c \approx (4-6)\Theta$ for all simulations with $\sigma_0 = 10$ (and $\langle \gamma'' \rangle_c \approx 15\Theta$ for $\sigma_0 = 50$), showing virtually no signs of radiative cooling.

Most of the plasmoids show their synchrotron emissivity building up in time, both for the core and for the layer.
Figure 4. Spacetime diagrams extracted along the reconnection midplane for the simulation s10Tm. From left to right we plot: (1) the magnetic field component \( B_y / B_0 \), (2) the electric field component \( E_z / B_0 \), (3) the Landau-type bulk velocity component \( \beta_x \), (4) logarithm of the cold magnetisation parameter \( \log(\sigma / \sigma_0) \), (5) mean particle Lorentz factor \( \langle \gamma' \rangle / \Theta \) measured in the Landau frame, and (6) logarithm of the particle number density \( \log(n/n_b) \). In the first five panels, we show the particle number density contours \( n = 7 n_b \) with green solid lines. We indicate the simulation time \( ct = 3.36 L_x \) corresponding to the simulation state shown in Figures 1 and 3 with the magenta dashed horizontal lines. The black dotted horizontal lines in the last panel indicate the initial simulation states shown in Figure 2.

Some of the smaller plasmoids show a very slow decline, some plasmoids show a rapid decline towards the end of their histories. The cores of large plasmoids produce significantly more (2-3 orders of magnitude) synchrotron emission per volume element than the cores of small plasmoids, and also about 2 orders of magnitude higher emission density than the large plasmoid layers. Our basic finding is that synchrotron emission of the plasmoids is sustained over a long term, irrespective of the cooling efficiency. For the cores of large plasmoids that undergo the most efficient radiative cooling in our high-temperature simulations, it appears that systematically increasing core density, as well as systematically increasing magnetic field strength are able to offset the reduction in particle mean energy. In the case of \( \sigma_0 = 50 \), we find a systematic increase of the emission from the cores of large plasmoids to the levels up to \( 10^7 P_{\text{syn}, b, 0} \).

### 4.6 Individual particles

Here we describe the behaviour of four selected energetic positrons (denoted as particles \#1 — \#4), for which a detailed history has been recorded. The spacetime tracks of these particles are indicated in Figure 10. From this one can see that particles \#1 and \#2 become trapped in different plasmoids, while particles \#3 and \#4 become accelerated in the low-density regions between plasmoids.

Figure 11 presents the detailed acceleration histories of these four particles, all in the same time window. Particle \#1 becomes accelerated to Lorentz factor \( \gamma \simeq 22 \Theta \) in a single episode. The beginning of acceleration episode coincides with the particle becoming trapped in the reconnection midplane (\( |y| < 5 n_0 \)). We can also see that the acceleration episode coincides with a formation of a new plasmoid that traps the particle also in the \( x \) coordinate. The particle experiences mainly the electric field component \( E_z \sim 0.1 B_0 \), and its momentum gain is at first mainly in the \( z \) direction.
Radiative simulations of steady-state plasmoid reconnection

Figure 5. Spacetime diagrams of average particle Lorentz factor \( \langle \gamma'' \rangle / \Theta \) measured in the Landau frame compared for three simulations characterised by different radiative cooling efficiencies.

Figure 6. Distributions of particle density \( n/n_b \) (left panel) and mean particle energy \( \langle \gamma'' \rangle / \Theta \) calculated in the Landau co-moving frame (right panel) extracted from the spacetime diagrams probing narrow regions along the reconnection layers for the 3 simulations with \( \sigma_0 = 10 \) that differ by initial gas temperature \( \Theta \), and hence by the radiative cooling efficiency (higher \( \Theta \) corresponds to more efficient cooling).

later also in the \( x \) direction. After the acceleration episode, the particle oscillates around the plasmoid centre, which results in oscillations of its Lorentz factor \( \gamma \) measured in the simulation frame. However, its Lorentz factor \( \gamma' \) measured in the plasmoid frame shows a slow gradual decline in time, which we attribute to the radiative cooling.

Particle \#2 shows a similar behaviour to particle \#1, the acceleration episode also coincides with the formation of a new plasmoid. In this case, the acceleration episode is shorter and the energy gain is also lower, with acceleration proceeding in similar electric fields. The radiative cooling is less efficient, and it should be noted that the co-moving perpendicular magnetic field is about twice weaker. At \( ct \simeq 2.1 L_x \), the plasmoid in which particle \#2 is trapped merges into a larger plasmoid on its right side. We see that this results in particle \#2 losing most of its energy measured in the simulation frame. It appears that this particle did not experience direct deceleration by strong electric
field, instead it just happened to be at its lower energy level in the oscillation cycle at the moment of plasmoids merger. It has also been kicked out from the reconnection midplane, settling briefly at $y \sim 15\rho_0$, trailing behind the merged plasmoid before it exits the right boundary of the domain.

Particle #3 was accelerated in a long acceleration episode under constant electric field $E_z \sim 0.12B_0$ after becoming trapped in the $y$ coordinate to the reconnection midplane. The $y$ and $B_x$ data reveal a typical Speiser orbit with gradually increasing period. The acceleration episode is interrupted at $ct \approx 2.1L_x$, when the particle begins to interact with a large plasmoid, which forces the particle away from the reconnection midplane. The particle bounces twice off the trailing edge of the plasmoid before it exits the left boundary of the domain. Particle #4 shows another example of Speiser-orbit acceleration in the low-density region of reconnection midplane that is enabled by trapping the particle in the $y$ coordinate.

### 4.7 Synchrotron emissivity and lightcurves

Figure 12 presents the spacetime distribution of the total synchrotron power emitted from the reconnection midplane in all directions for simulation $s10Tm$, and lightcurves that would be received by two distant observers placed at the opposite sides of the simulation domain, one on the left ($-x$ axis), one on the right ($+x$ axis). While the spacetime diagram of synchrotron power is based on the $x$-profiles integrated over narrow stripes of $|y| < \delta/2$, the lightcurves are calculated from a much wider region of $|y| < L_x/2$ to ensure contribution from the whole plasmoids and their broader surroundings.

The synchrotron emission is strongly concentrated along the plasmoid trajectories. As we have noted in Section 4.5, large plasmoids produce significantly more synchrotron emission. The emissivity of the large plasmoid cores exceeds the emissivity of the background plasma by four orders of magnitude.

Lightcurves received by either observer are completely different, indicating a high level of anisotropy. The most conspicuous features of the lightcurves are very sharp and bright.
flares. We can identify the origin of these flares by drawing the corresponding lightcones on the spacetime diagram to the events of small plasmoids approaching the observer with relativistic velocity and merging with a large target plasmoid that is also approaching the observer. Zooming up on these flares, we find their characteristic time scales of $\tau \sim \rho_0/c$, so they are basically unresolved\(^4\). Because of their very short duration, the contribution of these flares to the overall radiation fluence is rather insignificant. The lightcurves also contain smooth structures characterised by relatively long rise and short decline. The can be attributed to large plasmoids propagating with mildly relativistic velocities, and the sharp declines observed in the lightcurves coincide with these plasmoids exiting the simulation domain. The contributions of these plasmoids to the lightcurves is not complete, since their emission is not contained within the simulation boundaries, as we have already noted in Section 4.5.

Comparing the lightcurves recorded in two frequency bands\(^5\), we find that in general the contribution from large plasmoids is higher in the lower-frequency band (cyan lines), while the flares produced by small plasmoids are stronger in the higher-frequency band (orange lines). Lightcurves obtained from different simulations are qualitatively very similar, in particular the level of radiative cooling efficiency does not clearly affect the lightcurve characteristics.

4.8 Energy conservation

Because of the use of open boundaries with steady injection of fresh particles, our simulations do not conserve energy globally. In order to evaluate the efficiency of energy transformations, we defined a fixed region $R$ centred around the reconnection midplane between the left/right absorbing boundary layers, defined by $2\Delta_{abs} < x < (L_x - 2\Delta_{abs})$ and

---

\(^4\) Contributions to every lightcurve are recorded at every simulation timestep at the temporal resolution of $dt$.

\(^5\) The lightcurves are presented in the same arbitrary units equivalent to the $\nu F_{\nu}$ flux density.
Figure 9. Further histories of individual plasmoids. From the top, the rows present: (1) mean synchrotron emissivity $\langle P_{\text{syn},c} \rangle / P_{\text{syn},b,0}$ of the plasmoid core, (2) mean synchrotron emissivity $\langle P_{\text{syn},l} \rangle / P_{\text{syn},b,0}$ of the plasmoid layer. See Figure 7 for more description, and Eq. (8) for the definition of $P_{\text{syn},b,0}$.

Figure 10. Spacetime diagram of the tracks of selected energetic particles, the acceleration of which is characterised in detail in Figure 11. The line colour indicates the instantaneous particle energy measured in the simulation frame. Particle density contours $n = 7n_0$ are indicated with grey lines.

$-L_x/4 < y < L_x/4$. In addition to the instantaneous energy contained in $R$ in the form of magnetic and electric fields, as well as in the particles, we also calculate the cumulative energy emitted by all particles in the synchrotron process, and the fluxes of particles and electromagnetic fields (i.e., the Poynting flux) inflowing/outflowing across the $R$ boundaries.

In Figure 13, we present the time evolution of different forms of energy contained in the region $R$ for the simulation s10Tm. At the beginning of the simulation, the region $R$ is dominated by magnetic energy ($E_B,0 \approx 0.6E_{\text{tot}}$). The initial ($ct/L_x < 0.6$) energization of particles at the cost of magnetic energy is due to the trigger mechanism. This is followed by the somewhat erratic variation of the particle energy, which reflects systematic heating by magnetic reconnection and episodic escapes of large plasmoids. Over the course of the simulation ($ct/L_x \approx 4.5$), the magnetic energy of the region $R$ decreases by $\approx 40\%$, while the particle energy decreases only by about $\approx 15\%$. At the same time, we measure a large net influx of electromagnetic energy (accumulating to $\approx 1.5E_{B,0}$), mainly through the top/bottom boundaries of the region $R$, and even larger net outflow of particle energy, mainly through the left/right boundaries. The net energy outflow (particle minus electromagnetic) through the region boundaries amounts to $\approx 0.25E_{B,0}$ of the initial magnetic energy, which is slightly less than the particle energy lost to the synchrotron radiation ($\approx 0.3E_{B,0}$). Accounting for all these energy components and flows, the total energy in $R$ is conserved at the $\sim 0.1\%$ level.

4.9 Energy distributions of particles and photons

Figure 14 shows the energy distributions of all particles: electrons and positrons. For each simulation, it is averaged over a period of time that excludes only the initial stage ($ct/L_x < 0.85$). In all studied cases, the particle energy distributions established after the initial period show no significant evolution in time. As energetic particles escape across the open left/right boundaries, other particles are energised across the current layer, and are subject to radiative energy losses within the plasmoids. The balance between these processes is maintained regardless of the efficiency of radiative cooling. In the case of $\sigma = 10$, the high-energy excess ap-
Figure 11. Acceleration histories for selected energetic particles, the \( x \)-positions of which are shown in Figure 10. For each particle we present detailed information as functions of simulation time on five panels, from the left: (1) the \( y \)-position measured from the reconnection midplane, (2) particle Lorentz factor \( \gamma \) (also indicated with a colour scale on panels 1 and 2) normalised to \( \Theta \), (3) three components of the particle 4-velocity \( \mathbf{u} \) (\( x \) - red, \( y \) - green, \( z \) - blue) normalised to \( \Theta \), (4) three components of the local electric field \( \mathbf{E} \) in units of \( B_0 \), (5) three components of the local magnetic field \( \mathbf{B} \) in units of \( B_0 \). For particles #1 and #2, the black lines indicate parameter values (Lorentz factor \( \gamma' \), electric field \( E'_\parallel \) and magnetic field \( B'_\perp \)) measured in the co-moving frame of a plasmoid to which the particle is attached.

pears to be modest, because a large fraction of the simulation volume contains only the cold background particles. The relative number of energetic particles is not sensitive to the efficiency of radiative cooling. In the case of \( \sigma = 50 \), particles can reach energies higher by factor 5, and hence energetic particles form a distinct second component of the distribution that is locally harder than \( p = 2 \).

Figure 15 shows the spectral energy distributions (SED) \( \nu F_\nu \) of the synchrotron emission produced by all particles in all directions, averaged over the same periods of time as the particle energy distributions presented in Figure 14. In the case of \( \sigma = 10 \), the SED are dominated by the contribution from cold particles, with a high-frequency excess, the level of which increases with decreasing gas temperature \( \Theta \). This means that radiative cooling suppresses the high-frequency radiation component more clearly than it affects the high-energy particle tail. In the case of \( \sigma = 50 \), the SED is strongly dominated by the contribution from energetic particles with the maximum photon energies consistent roughly with an increase by factor 25, as expected from the increase of maximum particle energies: \( \nu_{\text{syn, max}} \propto u_{\text{max}}^2 \).

5 DISCUSSION

Our results are consistent with the basic picture of steady-state relativistic plasmoid reconnection that has been established since the work of Sironi et al. (2016). The open
Figure 12. The left and right panels show synchrotron lightcurves measured by observers placed at left and right sides of the reconnecting layer, respectively. The cyan and orange lines correspond to the two selected frequency bands indicated in Figure 15. The middle panel shows the spacetime distribution of total synchrotron power normalised to the value $P_{\text{syn,b,0}}$ defined in Eq. (8). The particle number density contours $n = 7n_b$ are shown with green solid lines. The black dashed lines represent the lightcones corresponding to selected features in either light curve. The magenta dashed line indicates the simulation time $ct = 3.36L_x$ corresponding to the simulation state shown in Figures 1 and 3.

Figure 13. Top panel: mean energy densities, normalised to initial magnetic energy density $U_{B,0}$, calculated for the analysis region $\mathcal{R}$, defined by $2\Delta_{\text{abs}} < x < L_x - 2\Delta_{\text{abs}}$ and $-L_y/4 < y < L_y/4$, for the simulation s10Tm. We plot the contributions from magnetic fields (green), electric fields (magenta), all particles ($e^+$ and $e^-$; red), and total conserved energy including inflows and outflows (black). The black dotted line shows the total instantaneous energy contained in $\mathcal{R}$, including contributions from the magnetic and electric fields, and all particles. Middle panel: energy flux inflows and outflows: the inflowing Poynting flux (green), outflowing particle energy flux (red), and total synchrotron emission (blue). Bottom panel: conservation of the total energy of region $\mathcal{R}$, including inflows and outflows.

boundaries allow to evacuate the reconnected plasma and develop unimpeded reconnection outflows that reach relativistic bulk velocities of the order of the background Alfvén velocity. In the centrally positioned active magnetic X-point, plasmoids are generated spontaneously over a wide range of sizes. The smallest discernible plasmoids, with the width of order $\sim 20\rho_0$ (deep blue lines in Figures 7, 8 and 9) are rapidly accelerated to relativistic speeds, their lifetimes are about $\sim 0.2L_x/c$. The largest plasmoids, with the width of order $\sim (150 - 200)\rho_0$ (brown lines in Figures 7, 8 and 9) accelerate slowly, reaching mildly relativistic speeds $\sim c/2$ only after $\sim 1.5L_x/c$. This is a confirmation of the anti-correlation between growth and acceleration of plasmoids identified by Sironi et al. (2016).

An important consequence of the relation between plasmoid sizes and their acceleration length scale is that it is not possible to contain the bulk acceleration of large plasmoids within a domain with open boundaries. We have found that as we increase the size of the simulation domain, given pro-
portionally more simulation time, we obtain ever larger plasmoids that require ever more acceleration time. In other words, we are unable to isolate the bulk acceleration of plasmoids from the domain boundaries and obtain a coasting phase of uniformly Alfvénic reconnection outflows.

A similar problem applies to the production of synchrotron light curves. In our simulations, even for the highest plasma temperatures investigated, we were not able to contain synchrotron emission from the largest plasmoids within the domain boundaries. As we show in Figures 8 and 9, even though in the case of short cooling length $l_{\text{cool}} \ll L_x$ the plasmoid cores indeed undergo efficient radiative cooling within the $L_x/c$ time scale, the total synchrotron emission of plasmoid cores and layers is not decreasing significantly.

The nominal cooling time scale cannot be reduced indefinitely by a further increase of the particle distribution temperature. We found that, with unrestricted radiation reaction, already for $\Theta \approx 1.25 \times 10^5$, the particle densities in the cores of the largest plasmoids increase to the level at which the Debye length eventually becomes unresolved $\lambda_D = (\Theta m_e c^2/4\pi n e^2)^{1/2} < dx$, which leads to the development of numerical electrostatic instabilities. Our ad hoc restriction of radiation reaction (see a footnote preceding Eq. 5) helps to avoid developing these instabilities. We note that a similar restriction of radiation physics (pair creation) in the plasmoid cores has been applied by Hakobyan et al. (2019).

Christie et al. (2019a) calculated lightcurves of synchrotron and inverse Compton emission by post-processing the results of non-radiative PIC simulations of relativistic steady-state reconnection of Sironi et al. (2016). Their radiation transfer model is based on several assumptions that can be verified by the results of our radiative PIC simulations. One of their key assumptions is that the intrinsic structure of plasmoids is not important and can be approximated by using their average parameters that in addition are constant in time. Our results suggest that the synchrotron emissivity is strongly concentrated in the central parts of the plasmoids (see the bottom panel of Figure 3), and that in the radiatively efficient regime the plasmoid cores undergo significant time evolution with systematic increase of plasmoid core density and peak magnetic field strength (Figure 7). We suggest that small plasmoids and the cores of large plasmoids are important for understanding the production of rapid radiation flares. Investigation of these structures is also the most challenging from the numerical perspective.

Our study suggests that properly resolving the cores of large plasmoids will be critical for understanding the radiative signatures of plasmoid reconnection. Recent non-radiative PIC simulations of relativistic reconnection demonstrated an important role of large plasmoids in extending the high-energy tail of the particle energy distribution along a power-law of slope $\sim 2$ (Petropoulou & Sironi 2018). However, taking into account radiative cooling, which is expected to be particularly strong in the plasmoid cores, the maximum energy achievable in the plasmoids may be significantly limited.

Plasmoid mergers have been suggested previously to be important for particle acceleration in relativistic reconnection, based on the results of relatively modest PIC simulations of reconnection in periodic boundaries (Nalewajko et al. 2015). Subsequent numerical studies emphasised a more fundamental role of magnetic X-points as a crucial first step for particles that eventually achieve the highest energies (Guo et al. 2019; Ball et al. 2019). Here we would like to point out that even if plasmoid mergers may not dominate particle acceleration in non-radiative reconnection, they are important for the production of transient

---

**Figure 14.** Energy distributions of all particles contained in the simulation domain, compared for the 4 main simulations. For each simulation, the distribution is averaged over simulation time, excluding the initial stage ($ct/L \lesssim 0.85$), in the space of flux logarithm. The dashed lines indicate the corresponding standard deviation values. The distributions are presented in arbitrary units, they are normalised to match the low-energy sections.

**Figure 15.** Isotropic spectra of the synchrotron radiation emitted across the simulation domain, compared for the 4 main simulations. For each simulation, the distribution is averaged over simulation time, excluding the initial stage ($ct/L \lesssim 0.85$), in the space of flux logarithm. The dashed lines indicate the corresponding standard deviation values. The frequencies are normalised to the nominal synchrotron frequency of the background plasma $\nu_{\text{syn, b}}$ defined in Eq. (10). The distributions are presented in arbitrary units, they are normalised to match the low-energy sections.
radiation signals. Even if major head-on collisions of large plasmoids are rare events, tail-on collisions of unequal plasmoids, arguably more frequent events due to the growth-acceleration anti-correlation, can be responsible for most of the sharpest features observed in the resulting lightcurves. Recent results of Christie et al. (2019b) show that large plasmoids can attract many small plasmoids originating on both sides of their trajectory, enhancing the rates of both tail-on and head-on mergers.

Our results also demonstrate the co-existence of plasmoids and minijets in the same reconnection layer. We find minijets persisting in the gaps forming between plasmoids, plasmoids are able to slide along a minijet without causing much disturbance, and a minijet reforms behind a passing plasmoid. The structure of minijets found in our simulations is qualitatively very similar to the analytical model of Lyubarsky (2005). Although the minijets contain some highly energetic particles, their contribution to the observed radiative signatures appears to be very weak. There may be two reasons behind this: (1) the minijets are characterised by much lower particle density than the plasmoids, (2) energetic particles propagating along the minijets show only weak radiative energy losses due to weak perpendicular magnetic field component.  

6 CONCLUSIONS

We presented the results of the first kinetic simulations of relativistic magnetic reconnection (RMR) within open boundaries that enable steady-state plasmoid reconnection and including the synchrotron radiation reaction. We confirm the general picture of steady-state relativistic plasmoid reconnection established by Sironi et al. (2016) and subsequent works. We find that synchrotron emission of plasmoids cannot be contained within open boundaries. The cores of large plasmoids are the main sites of synchrotron emission, their particle densities are significantly enhanced due to radiative pressure losses. Rapid flares of synchrotron radiation can be produced by tail-on mergers between small/fast plasmoids with large/slow targets. The plasmoids are also found to co-exist with the minijets that do not produce a lot of radiation due to their low particle densities.

ACKNOWLEDGMENTS

We acknowledge discussions with Benoît Cerutti, Dimitrios Giannios and Maria Petropoulou. The original version of the Zeltron code was created by Benoît Cerutti and co-developed by Gregory Werner at the University of Colorado Boulder (http://benoit.cerutti.free.fr/Zeltron/). These results are based on numerical simulations performed at the supercomputer Prometheus located at Cyfronet AGH, Poland (PLGrid grant recjose18; PI: K. Nalewajko). This work was supported by the Polish National Science Centre grant 2015/18/E/ST9/00580.

6 Radiative cooling in the minijets might be stronger if the guide field component $B_z$ were included.

REFERENCES

Aharonian, F., Akhperjanian, A. G., Bazer-Bachi, A. R., et al. 2007, ApJ, 664, L71
Aharonian, F., Akhperjanian, A. G., Anton, G., et al. 2009, ApJ, 696, L150
Ahnen, M. L., Ansoldi, S., Antonelli, L. A., et al. 2016, A&A, 593, A91
Albert, J., Aliu, E., Anderhub, H., et al. 2007, ApJ, 669, 862
Alves, E. P., Zrake, J., & Fiuzaa, F. 2018, Phys. Rev. Lett., 121, 245101
Ball, D., Sironi, L., & Ozel, F. 2019, ApJ, 884, 57
Begelman, M. C. 1998, ApJ, 493, 291
Begelman, M. C., Fabian, A. C., & Rees, M. J. 2008, MNRAS, 384, L19
Blandford, R. D., & Znajek, R. L., 1977, MNRAS, 179, 433
Böttcher, M., Reimer, A., Sweeney, K., & Prakash, A. 2013, ApJ, 768, 54
Bromberg, O., Singh, C. B., Davelaar, J., et al. 2019, ApJ, 884, 39
Buehler, R., Scargle, J. D., Blandford, R. D., et al. 2012, ApJ, 749, 24
Cerutti, B., Werner, G. R., Uzdensky, D. A., et al. 2013, ApJ, 770, 147
Cerutti, B., Werner, G. R., Uzdensky, D. A., & Begelman, M. C. 2014, ApJ, 782, 104
Chiaberge, M., & Ghisellini, G. 1999, MNRAS, 306, 551
Christie, I. M., Petropoulou, M., Sironi, L., et al. 2019, MNRAS, 482, 65
Christie, I. M., Petropoulou, M., Sironi, L., et al. 2019, arXiv:1908.02764
Daughton, W., Scudder, J., & Karimabadi, H. 2006, Physics of Plasmas, 13, 072101
Daughton, W., & Karimabadi H. 2007, Physics of Plasmas, 14, 072303
Fossati, G., Maraschi, L., Celotti, A., et al. 1998, MNRAS, 299, 433
Ghisellini, G., Righi, C., Costamante, L., & Tavecchio, F. 2017, MNRAS, 469, 255
Giannios, D., & Spruit, H. C. 2006, A&A, 450, 887
Giannios, D., Uzdensky, D. A., & Begelman, M. C. 2009, MNRAS, 395, L29
Giannios, D. 2013, MNRAS, 431, 355
Giannios, D., & Uzdensky, D. A. 2019, MNRAS, 484, 1378
Goyal, A., Stawarz, L., Ostrowski, M., et al. 2017, ApJ, 837, 127
Guo, F., Li, H., Daughton, W., & Liu, Y.-H. 2014, PhRvL, 113, 155005
Hakobyan, H., Philippov, A., & Spitkovsky, A. 2019, ApJ, 877, 53
Itoh, R., Nalewajko, K., Fukazawa, Y., et al. 2016, ApJ, 833, 77
Katarzyński, K., Lenain, J.-P., Zech, A., Boisson, C., & Sol, H., 2008, MNRAS, 390, 371
Kirk, J. G., Rieger, F. M., & Mastichiadis, A. 1998, A&A, 333, 452
Kirk, J. G., & Skjæraasen, O. 2003, ApJ, 591, 366
Komissarov, S. S., Blandford, R. D., & Romanova, M. M., 2007, MNRAS, 380, 51
Levinson, A., 2007, ApJ, 671, L29
Li, Z.-Y., Chiueh, T., & Begelman, M. C. 1992, ApJ, 394, 459
Lourens, N. F., Schekochihin, A. A., & Cowley, S. C. 2007, Physics of Plasmas, 14, 105703
Lovelace, R. V. E., Newman, W. I., & Romanova, M. M., 1997, ApJ, 484, 628
Lyubarsky, Y. E., 2005, MNRAS, 358, 113
Meyer, M., & Sikora, M. 2016, ARA&A, 54, 725
Meyer, M., Scargle, J. D., & Blandford, R. D. 2019, ApJ, 877, 39

MNRAS 000, 000–000 (0000)
