Relativistic jets shine through shocks or magnetic reconnection?

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ABSTRACT
Observations of gamma-ray bursts and jets from active galactic nuclei reveal that the jet flow is characterized by a high radiative efficiency and that the dissipative mechanism must be a powerful accelerator of non-thermal particles. Shocks and magnetic reconnection have long been considered as possible candidates for powering the jet emission. Recent progress via fully-kinetic particle-in-cell simulations allows us to revisit this issue on firm physical grounds. We show that shock models are unlikely to account for the jet emission. In fact, when shocks are efficient at dissipating energy, they typically do not accelerate particles far beyond the thermal energy, and vice versa. In contrast, we show that magnetic reconnection can deposit more than 50% of the dissipated energy into non-thermal leptons as long as the energy density of the magnetic field in the bulk flow is larger than the rest mass energy density. The emitting region, i.e., the reconnection downstream, is characterized by a rough energy equipartition between magnetic fields and radiating particles, which naturally accounts for a commonly observed property of blazar jets.

Key words: acceleration of particles — galaxies: jets — gamma-ray burst: general — magnetic reconnection — radiation mechanisms: non-thermal — shock waves

1 INTRODUCTION
Relativistic jets are ubiquitous in the Universe, with gamma-ray bursts (GRBs) and active galactic nuclei (AGNs) being two representative examples of high-energy sources powered by jets. Despite decades of research, the issue of what powers relativistic astrophysical jets is still unresolved. On the one hand, there is a strong theoretical motivation to consider them as magnetically-dominated objects at their base (Blandford & Znajek 1977, Blandford & Payne 1982). The strong magnetic fields threading a rotating compact object or the associated accretion disk serve to convert the rotational energy of the central engine into the power of the outflow. Part of the magnetic energy is used to accelerate the flow to relativistic speeds but, generally, magnetohydrodynamic (MHD) models predict that the largest fraction of the energy remains locked in the magnetic field, i.e., the jet arrives Poynting-dominated at the dissipation distance (Tchekhovskoy et al 2009, Komissarov et al 2009, Lyubarsky 2009). On the other hand, leptonic models of the radiative signature of blazars – a subclass of AGNs whose jets point along our line of sight – indicate that the energy densities of magnetic field $U_B$ and radiating particles $U_e$ at the emission region do not differ by more than one order of magnitude (Readhead 1992, Celotti & Ghisellini 2008, Ghisellini et al 2014).

Our ignorance on the source of jet power allows for two very different scenarios regarding the dissipative mechanism behind the jet emission. Hydrodynamical flows can dissipate energy and efficiently accelerate particles at shock fronts (e.g., Heavens & Drury 1988, Rees & Meszaros 1994), whereas in strongly magnetized flows the efficiency of shocks is greatly reduced (Kennel & Coroniti 1984b). In this case, a more likely candidate for powering the jet emission is the process of magnetic reconnection (Spruit et al 2001, Drenkhahn & Spruit 2002), where the annihilation of field lines of opposite polarity transfers the field energy to the particles.

Observations of relativistic jets can be used to infer the properties of the emitting region. Yet, these do not directly represent the typical conditions in the bulk flow. Emitting regions are, by definition, special places where intense energy dissipation has occurred. Both the shock downstream and the reconnection outflow qualify as such regions. On the contrary, the upstream medium (in the case of shock dissipation) or the inflow region (for the case of reconnection) describe more closely the bulk of the jet. In this sense, the observational inference on the ratio $U_B/U_e$, applies directly to the emitting region and not to the bulk of the jet flow. Hereafter, we identify the (shock or reconnection) upstream with the large-scale jet and the downstream with the emitting region.

The properties of the emitting region are related to those of the bulk of the jet in a way that depends on the dissipative process. In shocks the magnetic field is generally compressed and its strength is, thus, increased: the magnetic energy per particle is larger in the

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downstream than in the upstream. In contrast to shocks, magnetic reconnection dissipates magnetic energy: the downstream (i.e., the outflow) is less magnetized than the upstream (i.e., the inflow). In fact, the more efficient is reconnection in accelerating particles, the more particle-dominated the downstream region becomes.

In both shocks and reconnection, the partition of energy between magnetic fields and particles can be described self-consistently via fully-kinetic particle-in-cell (PIC) simulations. Recent PIC progress has been remarkable. First-order Fermi acceleration in relativistic shocks has been demonstrated from first principles (Santos et al. 2008a b; Martins et al. 2009; Haugbølle et al. 2011; Sironi et al. 2013). However, as expected from theoretical arguments (Begelman & Kirk 1990), no particle acceleration takes place when relativistic shocks are quasi-perpendicular, i.e., where the field lines are nearly orthogonal to the shock direction (Gallant et al. 1992; Sironi & Spitkovsky 2009). Recent PIC simulations have also demonstrated that magnetic reconnection in magnetically-dominated plasmas is a fast, unsteady process that leads to efficient particle acceleration (e.g., Sironi & Spitkovsky 2014; Kagan et al. 2013, for a review). In both shocks and reconnection, leptons pick up a large fraction of the dissipated energy, so electrons are the most likely candidates for powering the observed emission (i.e., supporting the so-called leptonic scenario).

In this work, we use jet observations to probe the nature of the dissipative mechanism and the properties of the bulk of the jet. For this, we exploit the latest PIC findings about shocks and magnetic reconnection as well as three basic facts about jet emission: (i) jets are efficient emitters, (ii) the radiating particles have extended, non-thermal distributions, and (iii) the energy densities stored in the magnetic field and in radiating particles do not differ more than one order of magnitude. The paper is structured as follows. In Section 2 we summarize these basic observational constraints. In Section 3 we apply them to two alternative dissipative scenarios: shocks and magnetic reconnection. We conclude in Section 4.

2 INFERENCES FROM JET OBSERVATIONS

There is a consensus that relativistic jets are efficient emitters. GRBs have efficiency $f \sim 0.1 - 1$ in converting jet power into gamma-ray (prompt) emission (e.g., Berger et al. 2000; Panaitescu & Kumar 2002; Yost et al. 2003; Granot et al. 2004; Nava et al. 2013). Blazars have similarly high radiative efficiencies. Recently, Ghisellini et al. (2014) modelled the spectral energy distribution (SED) of a large sample of blazars and, assuming one proton per radiating electron, they inferred a radiative efficiency $f \sim 0.1$ (see also Celotti & Ghisellini 2008). The efficiency can be even higher, if blazar jets are pair-loaded, namely they contain more leptons than protons (Ghisellini et al. 2014). It follows that the dissipative mechanism behind the jet emission must be able to transfer a large amount of energy to the radiating particles, achieving a high radiative efficiency ($f \gtrsim 0.1$). Given that $f$ is already high, the dissipation process must involve a large fraction of the jet fluid, in order to explain the observed photon luminosities without pushing the jet energetics to extreme values.

Multi-wavelength jet observations typically suggest that the radiating particles are accelerated into a non-thermal power-law distribution. The blazar SED has a characteristic double-humped shape (e.g., Uebl et al. 1997; Fossati et al. 1998), with a broad low-energy component extending from the radio up to the UV band, and in some extreme cases up to $\gtrsim 1$ keV X-rays (Costamante et al. 2001). The high-energy component extends across the X-ray and y-ray bands, with peak energy around 0.1 TeV, although this is not always clear (see, e.g., Abdo et al. 2011 for Mrk 421). The low-energy hump is believed to result from synchrotron emission from relativistic electrons. Its spectral shape requires, in general, a power-law (or broken power-law) particle spectrum (e.g., Celotti & Ghisellini 2008), which implies non-thermal particle acceleration. Furthermore, the high-energy hump of blazars, which in leptonic scenarios is explained as inverse Compton emission, extends in many sources up to 100 GeV or higher (e.g., see Aleksić 2012, for PG 1553+113). This also points to electron acceleration up to Lorentz factors $\gamma_e \gtrsim 10^5$, or even higher for TeV-emitting blazars. The same holds true in synchrotron models for the prompt emission of GRBs (Daigne et al. 2011, and references therein). In summary, the dissipative mechanism that operates in jets has to be able to accelerate particles well beyond their thermal energies.

In blazars both the synchrotron and the Compton humps are visible and their relative brightness constrains the ratio of energy densities $U_B/U_e$ at the emitting region (e.g., Sikora et al. 2009). The fact that in BL Lacs the two bumps have comparable luminosity indicates that, in the synchrotron self-Compton (SSC) model, $U_B/U_e \sim 1$ (see also Mastichiadis & Kirk 1997). However, for gamma-ray luminous blazars, such as flat spectrum radio quasars (FSRQs), where the high-energy component is much more luminous than the low-energy one, the SSC scenario would predict that $U_B/U_e \ll 1$, i.e., far from equipartition. Yet, a rough equipartition between the magnetic field and radiating electrons is still inferred in such blazars, in the external Compton (EC) interpretation of their gamma-ray emission (Sikora et al. 2009). In summary, detailed modeling of many blazars within the SSC or EC scenarios points to the fact that, quite generally, $U_B/U_e \sim 1$, but some emitting regions may be moderately magnetically-dominated, i.e., $U_B/U_e \sim 3$ (Ghisellini et al. 2014). The rough equipartition between emitting particles and magnetic field in the blazar emitting region is also suggested by the observed surface brightness temperatures at radio wavelengths (Readhead 1994; Lahéenmäki et al. 1999; Homan et al. 2006; Hovatta et al. 2013). In this work, we adopt the range $0.3 \leq U_B/U_e \leq 3$ as representative of the physical conditions in the emission region of blazars (Coppi & Aharonian 1999; Celotti & Ghisellini 2008; Ghisellini et al. 2014). In the case of GRBs, no high-energy bump in the SED can be securely associated with the prompt emission, making the $U_B/U_e$ ratio at the emitting region harder to quantify.

In summary, any model for the jet emission has to account for the fact that both blazars and GRBs are characterized by (i) a large fraction of the jet fluid, (ii) relativistic particles (typically leptons), and (iii) a non-thermal particle distribution. In this work, we apply them to two alternative dissipative scenarios: shocks and magnetic reconnection.
The dissipative efficiency, which, under our assumptions, is equal to the radiative one, is then

\[ f_{\text{sh}} = \frac{U_e}{e + pc^2} = \xi_e \left(1 - \frac{1}{\gamma_{\text{rel}}}\right) \tag{1} \]

which in the ultra-relativistic regime \((\gamma_{\text{rel}} \gg 1)\) reduces to \(f_{\text{sh}} \approx \xi_e\). For our physically-motivated choice of \(\xi_e = 0.5 - 1\), the above relation yields an efficiency in the range of the observed values.

In general, the upstream region is magnetized, and in this case the efficiency drops, unless \(\gamma_{\text{rel}}\) increases accordingly. Thus, shocks formed in a flow with any appreciable magnetization must be mildly relativistic or relativistic, in order to satisfy the efficiency requirement \(f_{\text{sh}} \approx 0.1\) (see, e.g., Mimica et al. 2009). This will be shown with detailed numerical examples at the end of this section. Similar to the hydrodynamic case, the downstream and upstream quantities can be related through the MHD jump conditions (for a treatment of the MHD jump conditions in relativistic shocks see, e.g., Kennel & Coroniti 1984; Apel & Camenzind 1988). In the following, we introduce the notation \(\xi_i\), with \(i = 1\) for upstream and \(i = 2\) for downstream, to refer to quantities of region \(i\) measured in the respective rest frame. The quantities to be determined and tested against the observational constraints are the efficiency \(f_{\text{sh}}\) and the ratio of the magnetic to particle energy densities \((U_B/U_e)\) as a function of the upstream magnetization \(\sigma\). From this point on, we will refer to \(\sigma_1\) simply as \(\sigma\), in order to use a uniform notation throughout the text. The efficiency is defined as

\[ f_{\text{sh}} = \frac{U_{e,2}}{e_2 + \rho_2 c^2 + U_{B,2}} \tag{2} \]

where \(U_{e,2} = \xi_e e_2\), \(U_{B,2} = B_2^2/8\pi\) and \(\rho_2\) is the proper density of the downstream region. The ratio of magnetic to electron energy is

\[ \frac{U_{B,2}}{U_{e,2}} = \frac{\Gamma - 1}{\xi_e \beta_2} \tag{3} \]

where \(e_2 = P_2/(\Gamma - 1)\), \(\Gamma\) is the adiabatic index ranging from \(5/3\) (non-relativistic ideal gas) to \(4/3\) (ultra-relativistic gas), and \(\beta_2 = P_2/U_{e,2}\) is the plasma beta parameter for the downstream region, while \(P_2\) is the gas pressure.

The parameter \(f_{\text{sh}}\) as defined above represents the fraction of post-shock energy that is deposited into electrons (or generally, pairs), and thus available to be radiated. Yet, the observational requirement of efficient dissipation \((f_{\text{sh}} \gtrsim 0.1)\) is not the only constraint imposed on the candidate dissipation mechanism. Unless shocks can accelerate particles to energies well beyond the peak of their thermal distribution, dissipation at shocks is not a promising process for explaining the observed non-thermal emission from jets (see also Section 2). In fact, it is well known that Fermi acceleration cannot operate in the so-called superluminal shocks (e.g., Kirk & Heavens 1989; Begelman & Kirk 1990; Sironi & Spitkovsky 2009, 2013; Sironi et al. 2013). There, particles moving along the magnetic field at the speed of light cannot outrun the shock, whose speed in the pre-shock medium is \(v_1 < c\), for ultra-relativistic flows. Since the Fermi process requires repeated crossings of the shock, this implies that Fermi acceleration in superluminal relativistic shocks is extremely inefficient. For this reason, we also require the shock to be subluminal, or equivalently the condition \(\beta_1 / \cos \theta_1 < 1\) should hold (see, e.g., eq. (1) in Kirk & Heavens 1989), where \(\beta_1 = v_1/c\) and \(\theta_1\) is the angle between the shock normal and the magnetic field as measured in the upstream frame.

We solve the MHD jump conditions for a cold upstream
medium and two representative shock obliquities, i.e. \( \theta_1 = 30^\circ \) and \( 60^\circ \). The results are presented in Fig. and our main conclusions can be summarized as follows:

- For \( \sigma \leq 0.1 \), the efficiency does not strongly depend on the flow magnetization.
- The shock must be mildly or ultra-relativistic in order to be efficient, since \( f_{ah} \geq 0.1 \) only for \( \gamma_{el} \beta_{el} \geq 1 \), which is consistent with the argument by Rees & Meszaros (1994).
- The ratio of magnetic to particle energy densities in the downstream increases with the upstream magnetization. For \( \sigma \approx 0.001 - 0.1 \), there is a small range in \( \gamma_{el} \beta_{el} \) (dependent on \( \sigma \)) leading to \( 0.3 \leq U_{B_{3}}/U_{e_{3}} \leq 3 \). This conclusion does not depend significantly on the magnetic field inclination relative to the shock front.
- For oblique shocks with angles \( \theta_1 \leq 30^\circ \), there is a small parameter space where all of the three constraints are satisfied, namely \( 0.03 \leq f_{ah} \leq 0.3 \), \( 0.3 \leq U_{B_{3}}/U_{e_{3}} \leq 3 \) and the shock is subluminal. This parameter regime involves mildly relativistic shocks with \( \beta_{el} \gamma_{el} \approx 1 \) and moderate magnetization \( \sigma \approx 0.1 \).
- For substantial magnetic field inclination (\( \theta_1 > 45^\circ \)), there is a strong tension between the efficiency constraint, that favors fast shocks, and the subluminality requirement. For \( \theta_1 = 60^\circ \), we find that the efficiency for subluminal shocks is at most \( \sim 3\% \).

The analysis above assumes that the magnetic field in the downstream region is merely the result of shock compression (as opposed to fields generated via plasma instabilities at the shock front) and that the upstream region is sufficiently magnetized for the superluminal constraint to hold, i.e., \( \sigma \geq 10^{-3} \) in electron-positron shocks and \( \sigma \geq 10^{-4} \) in electron-proton shocks (Sironi et al. 2013).

We showed that the observational requirements of \( 0.03 \leq f_{ah} \leq 0.3 \), \( 0.3 \leq U_{B_{3}}/U_{e_{3}} \leq 3 \) and the subluminality constraint for efficient particle acceleration are simultaneously satisfied only for certain shock obliquities and for \( \sigma > 0.01 \). In this case, the downstream field is likely to be dominated by the shock-compressed component, in agreement with our assumptions.

These assumptions, however, may not hold in other environments, such as the external shocks of GRBs. There, the shock propagates into the interstellar medium, whose magnetization is extremely small (\( \sigma \sim 10^{-9} \)). Modelling of the GRB afterglow emission implies that the emitting region is far from equipartition, having \( U_e \gg U_B \) (Kumar & Barniol Duran 2009, 2014, Lemoine 2013, Sironi et al. 2013). Thus, in the context of GRBs, our arguments are applicable only to the “internal” jet emission mechanism and not to the afterglow phase.

### 3.2 Magnetic reconnection

In reconnection, magnetic field energy is transformed into particle energy. Assuming that half of the magnetic energy is converted into internal energy (quite a realistic assumption, as we demonstrate below with PIC simulations), and assuming fast cooling electrons, the expected radiative efficiency is \( f_{rec} \approx 0.5 \xi_c \sigma/(\sigma + 2) \). For \( \xi_c \sim 0.5 \), as found in PIC simulations of electron-positron (yielding \( \xi_c \sim 1 \)) and electron-proton (giving \( \xi_c \sim 0.5 \)) reconnection (e.g., Sironi & Spitkovsky 2013; Melzani et al. 2014), this implies that \( \sigma \approx 1 \) for reconnection to be energetically viable in powering the jet emission. Hence, we deal with relativistic reconnection, in which the mean magnetic energy per particle is larger than the rest mass energy, or equivalently \( \sigma \geq 1 \).

Recent progress in the study of relativistic reconnection via PIC simulations has been remarkable (see Kagan et al. 2015, for a review), demonstrating that reconnection in high-\( \sigma \) plasma is fast and unsteady. The reconnection downstream is composed of plasmoids that contain the reconnected plasma. In two dimensions (2D), they appear as overdense magnetic islands (see the 2D density pattern in the top panel of Fig. connected by thin
X-lines. The plasmoid instability further fragments each X-line into a series of smaller islands, separated by X-points (e.g., see the small islands at $700c/\omega_p < x < 1400c/\omega_p$ in the top panel of Fig. 2 where $c/\omega_p$ is the electron skin depth). At the X-points, the particles are not tied to the magnetic field lines and they get accelerated along the reconnection electric field. The late-time particle spectrum integrated over the whole reconnection region is a power-law, whose slope is $-2$ for $\sigma = 10$ (as commonly assumed in blazar SED modeling), and becomes harder with increasing magnetization (Sironi & Spitkovsky 2014; see also, e.g., Zenitani & Hoshino 2001; Jaroschek et al. 2004; Bessho & Bhattacharjee 2012; Kagan et al. 2013; Cerutti et al. 2014; Guo et al. 2014; Werner et al. 2014). Efficient particle acceleration to non-thermal energies is a generic by-product of the long-term evolution of relativistic reconnection in both two and three dimensions. In three dimensions (3D), the so-called drift-kink mode corrugates the reconnection layer at early times (Daughton 1998; Zenitani & Hoshino 2008), but the long-term evolution is controlled by the plasmoid instability, that facilitates efficient particle acceleration, in analogy to the two-dimensional physics (Sironi & Spitkovsky 2013).

In summary, recent PIC simulations of relativistic reconnection have convincingly demonstrated that for $\sigma \gtrsim 1$ (of interest for relativistic jets), the accelerated particles populate extended non-thermal power-laws. We now investigate whether relativistic reconnection can provide the required efficiency $f_{rec}$ and kinetic-to-magnetic energy ratio $U_e/U_B$ inferred from jet observations.

The dissipated magnetic energy in the reconnection downstream is distributed between particles and magnetic fields. The plasmoids appear as inhomogeneous structures, with the core dominated by particle energy (panel (d) in Fig. 2), whereas the outskirts are magnetically dominated (panel (c) in Fig. 2). Because of the inhomogeneity of the plasmoids, it is meaningful to define the dissipative efficiency as a volume- or surface-averaged quantity in 3D and 2D, respectively, summed over many plasmoids:

$$f_{rec} \equiv \frac{\sum_i \int_{V_i} U_e dV_i}{\sum_i \int_{V_i} (e + p c^2 + U_B) dV_i}$$

where $V_i$ is the volume (or surface) of each plasmoid, $U_e$ is the kinetic energy of the emitting leptons (electrons alone for electron-proton reconnection, or both species for electron-positron reconnection), while the sum runs over all the plasmoids in the reconnection layer (see the colored regions in Fig. 2(b)).

The fact that the reconnection upstream is taken to be cold facilitates the identification of the regions where dissipation of energy has taken place, as we now explain. We select the plasmoid volume $V_i$ where to compute the integrals above (or actually surface, for 2D simulations as in Fig. 2) by selecting all the regions where the mean particle Lorentz factor is $\gamma > 1.1$ (more precisely, in electron-ion reconnection we employ here the mean ion Lorentz...
factor). The inflow region is cold and moves at the reconnection speed, which for \( \sigma \geq 10 \) is \( v_{\text{rec}} \sim 0.15c \) (e.g., Sironi & Spitkovsky 2014), so that the mean Lorentz factor there is only \( \gamma_{\text{rec}} \sim 1.01 \), which is below the threshold we employ. In Fig. 2(b), we only display the regions where the condition \( \gamma > 1.1 \) is met (otherwise, the color is artificially set as black), showing that this criterion provides an excellent identification of the reconnection downstream (i.e., of the magnetic islands). We have employed this criterion across the whole range of magnetizations we have investigated (\( \sigma = 1 \sim 50 \)) and for different strengths of the guide field \( B_g \) orthogonal to the annihilating fields \( B_0 \) (from \( B_g/B_0 = 0 \) up to 3), finding always an excellent spatial correlation with the reconnection downstream. In Fig. 2(c) and (d), we need to define a measure of \( U_e/(U_e + U_B) \) in the downstream region that closely corresponds to what the homogeneous models probe. In the limit of fast cooling particles, we should calculate the volume-average (surface-average in 2D simulations)

\[
\langle U_e/(U_e + U_B) \rangle \equiv \frac{\sum_i \int V_i U_e dV_i}{\sum_i \int V_i dV_i}.
\]

The reason for our choice of the weighting factor \( U_e \) is that regions with low electron energy density cannot contribute much to the emission. Although in regions where \( U_B > U_e \) (e.g., in the outskirts of the plasmoids, compare Fig. 2(c) and (d)), the synchrotron cooling rate is significantly higher than in regions with lower \( U_B \) (e.g., the plasmoid cores), the total radiated power in the fast cooling regime is always limited by the one injected into the emitting particles. As a side note, we remark that the choice to integrate \( U_e/(U_e + U_B) \) rather than \( U_e/U_B \) in Eq. (5) is motivated by the fact that the integral would otherwise be largely dominated by the very center of magnetic islands, where \( U_B \ll U_e \) (see Fig. 2(c) and (d)). Our choice is then appropriate for comparing to observational inferences based on homogeneous emission models.

By integrating over the surface of the magnetic islands, we compute \( f_{\text{rec}} \) and \( \langle U_e/(U_e + U_B) \rangle \) as described above. Our main results are shown in Fig. 3 for the case of electron-positron reconnection and in Fig. 4 for electron-proton reconnection. Our conclusions can be summarized as follows (see also Tables 1 and 2):

- Jet observations do not impose strong constraints on the reconnection model as long as the reconnection is relativistic. For \( \sigma \geq 1 \) and guide fields weaker that the reconnecting field (\( B_g < B_0 \)), all constraints are satisfied. We have checked that this conclusion holds for a broad range of conditions: pair and electron-ion plasma, and upstream magnetizations as large as \( \sigma \sim 50 \).

- In electron-positron plasmas (top panel in Fig. 3), the reconnection efficiency asymptotically approaches \( f_{\text{rec}} \sim 0.5 \) for \( \sigma \geq 10 \) and no guide field. In electron-ion plasmas (colored lines in the top panel of Fig. 4), it is roughly half of that value, since in the numerator of Eq. (3) only electrons contribute (as opposed to both species, for electron-positron plasmas). For lower \( \sigma \), the rest mass

| \( B_g/B_0 \) | 0.0 | 0.1 | 0.3 | 0.6 | 1 | 3 |
| --- | --- | --- | --- | --- | --- | --- |
| \( \sigma = 1 \) \( f_{\text{rec}} \) | 0.26 | 0.25 | 0.21 | 0.15 | 0.11 | 0.03 |
| \( \langle U_e/(U_e + U_B) \rangle \) | 0.75 | 0.70 | 0.57 | 0.36 | 0.22 | 0.04 |
| \( \sigma = 10 \) \( f_{\text{rec}} \) | 0.52 | 0.48 | 0.36 | 0.24 | 0.13 | 0.03 |
| \( \langle U_e/(U_e + U_B) \rangle \) | 0.72 | 0.66 | 0.52 | 0.35 | 0.22 | 0.04 |
| \( \sigma = 50 \) \( f_{\text{rec}} \) | 0.52 | 0.50 | 0.35 | 0.25 | 0.14 | 0.04 |
| \( \langle U_e/(U_e + U_B) \rangle \) | 0.70 | 0.63 | 0.49 | 0.35 | 0.21 | 0.05 |

Table 1. Results from 2D PIC simulations in electron-positron plasmas without guide field, as a function of the flow magnetization.

Table 2. Results from 2D PIC simulations in electron-positron plasmas, as a function of the flow magnetization and the guide field strength.
energy appreciably contributes to the denominator of Eq. (3), which results in a decrease of the reconnection efficiency.

- As a result of the efficient field dissipation, the energy density of the radiating particles is larger than that of the magnetic field at the emitting region by a factor of a few, both in electron-positron plasmas (black line in the bottom panel of Fig. 3) and electron-ion plasmas (colored lines in the bottom panel of Fig. 4). Once again, the difference between the colored lines and the black line in Fig. 4 is due to the fact that, in electron-positron reconnection (black line), both species contribute to the kinetic energy $U_e$ of emitting particles that enters in Eq. 5.

- For weak guide fields ($B_g/B_0 \leq 0.5$), the magnetic energy is very efficiently dissipated ($f_{\text{rec}} \approx 0.2$ in electron-positron plasmas). For progressively stronger guide fields, the radiative efficiency drops (see the colored circles in the top panel of Fig. 3) and the emitting regions turn magnetically dominated (colored circles in the bottom panel). The transition happens gradually around $B_g/B_0 \sim 0.5$, independently of $\sigma$.

- For electron-positron reconnection, we find that the ratio between particle kinetic energy and magnetic energy has an upper limit of $U_e/U_B \sim 3$, see Fig. 3(b) and $U_e/U_B \sim 2.5$ for electron-proton plasmas, see Fig. 4(b)). This is independent of $\sigma$ and is reached in the limit of weak guide fields. When increasing the guide field strength, the emission region gets more magnetized and $U_e/U_B$ can become arbitrarily small. However, we remark that when the emission region is strongly magnetized, the dissipation process is inefficient (i.e., $f_{\text{rec}} \ll 1$), so the resulting observational signature is very weak. It follows that in reconnection-dominated systems there is an observational bias to infer equipartition.

- We can provide a convenient parameterization of the results presented in Fig. 3 and Fig. 4 (see also Tables 1, 2, 3), assuming that a fraction $\sim 55\%$ of the magnetic energy in alternating fields (so, excluding the guide field) is converted into particle kinetic energy, and that the guide field energy is not dissipated. We obtain

$$f_{\text{rec}} = 0.55 \xi_e \frac{\sigma}{\sigma + 1 + (B_g/B_0)^2} + 2$$  \hspace{1cm} (6)

where $\xi_e \sim 1$ in electron-positron plasmas and $\xi_e \sim 0.5$ in electron-proton plasmas. Similarly, the expected equipartition parameter is

$$\frac{U_e}{U_e + U_B} = 1.3 \frac{0.55 \xi_e}{0.55 \xi_e + 0.45 + (B_g/B_0)^2}$$  \hspace{1cm} (7)

where the extra factor of 1.3 accounts for the fact that the integrand in Eq. 5 has been weighted with the electron energy density $U_e$. These two formulæ provide a satisfactory description of the trends observed in Fig. 3 and Fig. 4.

- Our results for electron-ion reconnection are nearly independent of the numerical choice of the mass ratio, which is constrained in PIC simulations to be smaller than the realistic value, for computational convenience. In fact, the three colored lines in Fig. 3 (blue for $m_p/m_e = 6.25$, green for $m_p/m_e = 25$ and red for $m_p/m_e = 100$) nearly overlap. It follows that our results can be confidently applied to the case of a realistic mass ratio.

### 4 DISCUSSION AND CONCLUSIONS

For a long time, theoretical modeling of relativistic astrophysical jets has faced a conundrum: if jets are launched as magnetically-dominated flows, as most theories predict, why do they appear to be in rough equipartition between magnetic fields and radiating particles at the emission region? Does the Poynting flux convert into bulk kinetic energy prior to the jet emission, and then shock waves are responsible for accelerating the radiating particles?

The answer is: probably not. For shocks to be efficient in powering the jet emission, we find that they need to be at least mildly relativistic, i.e., $\gamma_{\text{rel}} \approx 30^\circ$. To attain rough equipartition between $U_e$ and $U_B$ in the shock downstream, as inferred from modeling the SED of blazar sources, the magnetization in the shock upstream has to be in the range $\sigma \sim 0.01 - 0.1$. Under these conditions, for most magnetic field orientations, shocks are superluminal. Charged particles are constrained to follow the field lines, whose orientation prohibits repeated crossings of the shock. So, the particles have no chance to undergo Fermi acceleration, which instead is required to explain the broadband non-thermal emission signatures of blazar jets. We find that the shock model can work only in a small region of the parameter space, with $\gamma_{\text{rel}} \sim 1.5$ and $\theta_1 \lesssim 30^\circ$.

If the jet remains magnetically-dominated at the dissipation distance (as most models of MHD jets predict), dissipation through
reconnection is more likely. PIC simulations of magnetic reconnection in high-σ plasmas (see Kagan et al. 2013 for a review) can now describe the evolution of the system for long enough times to probe the dynamics of the downstream region as well as the physics of particle acceleration. The simulations have demonstrated that high-σ reconnection has several attractive features to explain the observations. The reconnection process is fairly fast, because of the fragmentation of the layer into plasmoid chains. The plasmoids (or magnetic islands) contain the reconnecting fluid and qualitatively behave like the emitting regions in this model. For σ ≳ 1, the particles are efficiently accelerated in the course of the reconnection process (e.g. Sironi & Spitkovsky 2014). In agreement with the observations, the particle energy spectrum is a power-law non-thermal distribution, with a harder power-law slope for higher σ. In this work, we have shown that the reconnection process is particularly efficient in depositing magnetic energy into non-thermal particles, and at the same time that the emitting regions are in equipartition between particles and magnetic field.

From a suite of 2D PIC simulations of relativistic reconnection, we have found that, for σ ≳ 1 and zero guide field, the emitting regions (plasmoids) have \( U_e/U_B \sim 3 \) and that a significant fraction of the the upstream energy is deposited into energetic electrons. More precisely, the dissipative efficiency reaches \( f_{\text{rec}} \sim 0.5 \) for electron-proton plasmas (where both species can contribute to the emission) and \( f_{\text{rec}} \sim 0.2 \) in electron-proton plasmas. In general, the presence of a strong guide field makes the process slower and less efficient, and the reconnection downstream becomes more magnetically-dominated. Nevertheless, for a broad range of guide field strengths \( B_g/B_B \lesssim 0.5 \), we find that \( U_e/U_B \gtrsim 0.3 \) and \( f_{\text{rec}} \gtrsim 0.1 \), which is within the range inferred from the observations. This conclusion holds regardless of the seed magnetization, in the regime \( \sigma \gtrsim 1 \) of relativistic reconnection (we tested from \( \sigma = 1 \) up to \( \sigma = 50 \)). In fact, the measured ratio \( U_e/U_B \) might be used to directly probe the strength of the guide field in the reconnection region (see Figs. 3 and 4).

The fact that \( \sigma \gtrsim 1 \) also implies that the emitting regions may be characterized by relativistic motions, with outflow Lorentz factors \( \Gamma_{\text{out}} = \sigma \) measured in the rest frame of the jet (Lyubarsky 2003), which has profound implications for jet variability (e.g., Lyutikov & Blandford 2003; Giannios et al. 2009; Nalewaja et al. 2011; Narayan & Piran 2012; Giannios 2013).

In summary, we have shown that shocks are unlikely to mediate the dissipation of energy in relativistic jets, mainly because the efficiency and subliminality constraints cannot be satisfied simultaneously. If the bulk of the jet has \( \sigma \gtrsim 1 \) and sufficiently small-scale fields prone to be dissipated, it appears that magnetic reconnection can satisfy all the basic conditions for the emission: extended particle distributions, efficient dissipation and rough equipartition between particles and magnetic field in the emitting region. Dissipation via reconnection is a promising process to explain the multi-wavelength non-thermal emission of relativistic jets.

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