

ABSTRACT

Title of Dissertation: **RADIATIVE PLASMAS IN
PULSAR MAGNETOSPHERES**

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Pulsars are highly magnetized rotating neutron stars known for their periodic bursts of radio emission. Decades of astronomical observations revealed that pulsars produce non-thermal radiation in all energy bands, from radio to gamma rays, covering more than 20 decades in photon energy. Modern theories consider strongly magnetized relativistic electron-positron plasmas to be the source of the observed emission. In my Thesis, I investigate physical processes that can be responsible for plasma production and the observed high-energy emission in the wide range of photon energies, from eV to TeV.

In the first Chapter of my Thesis, I investigate relativistic magnetic reconnection with strong synchrotron cooling using three-dimensional particle-in-cell kinetic plasma simulations. I characterize the spectrum of accelerated particles and emitted synchrotron photons for varying strengths of synchrotron cooling. I show that the cutoff energy of the synchrotron spectrum can significantly exceed the theoretical limit of 16 MeV if the plasma magnetization parameter exceeds the radiation reaction limit. Additionally, I demonstrate that a small fraction of ions present

in the current sheet can be accelerated to the highest energies, making relativistic radiative reconnection a promising mechanism for the acceleration of high-energy cosmic rays.

In the second Chapter, I present the first multi-dimensional simulations of the QED pair production discharge that occurs in the polar region of the neutron star. This process is believed to be the primary source of the pair plasma in pulsar magnetospheres and also the source of the radio emission. In this work, I focus on the self-consistently emerging synchronization of the discharges in different parts of the polar region. I find that pair discharges on neighboring magnetic field lines synchronize on a scale comparable to the height of the pair production region. I also demonstrate that the popular “spark” model of pair discharges is incompatible with the universally adopted force-free magnetospheric model: intermittent discharges fill the entire polar region that allows pair production, leaving no space for discharge-free regions. My findings disprove the key assumption of the spark model about the existence of distinct discharge columns.

In the third Chapter, I demonstrate how the key findings of two previous chapters can provide a self-consistent explanation of the recently discovered very-high-energy, reaching 20 TeV, pulsed emission in Vela pulsar. Motivated by the results of recent global simulations of pulsar magnetospheres, I propose that this radiation is produced in the magnetospheric current sheet undergoing radiative relativistic reconnection. I show that high-energy synchrotron photons emitted by reconnection-accelerated particles efficiently produce electron-positron pairs. The density of secondary pairs exceeds the supply from the polar cap and results in a self-regulated plasma magnetization parameter of $\sim 10^7$. Electrons and positrons accelerate to Lorentz factors comparable to $\sim 10^7$ and emit the observed GeV radiation via the synchrotron process and 10 TeV photons by Compton scattering of the soft synchrotron photons emitted by secondary pairs. My model self-consistently accounts for the ratio of the gamma-ray and TeV luminosities.

RADIATIVE PLASMAS IN PULSAR MAGNETOSPHERES

by

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Preface

Chapter 2 of this thesis is based on work in collaboration with Prof. Alexander Philippov and Hayk Hakobyan and published in the *Astrophysical Journal* as [Chernoglazov et al. \(2023\)](#). Chapter 3 is based on a paper in collaboration with Prof. Andrey Timokhin and Prof. Alexander Philippov and submitted for publication in the *Astrophysical Journal Letters*. Chapter 4 presents the preliminary results to be submitted to the *Astrophysical Journal Letters*.

During my time in graduate school, I had a privilege to work on numerous additional projects. A significant portion of my time in graduate school was spent working on studying the dynamics of binary compact objects, and I significantly contributed to a publication [Foucart et al. \(2021\)](#). During my internship at the Center for Computational Astrophysics, I performed an extensive study of the relativistic highly-magnetized turbulence, both in strong ([Chernoglazov et al., 2021](#)) and weak ([Ripperda et al., 2021](#); [TenBarge et al., 2021](#)) limits, in the fluid approximation. The results and methods that I developed during my work on relativistic reconnection received further development in studies of neutrinos from coronae of Active Galactic Nuclei ([Mbarek et al., 2024](#)).

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Table of Contents

Preface	ii
Acknowledgements	iii
Table of Contents	iv
List of Tables	vi
List of Figures	vii
List of Abbreviations	ix
Chapter 1: Introduction	1
1.1 Radiation from Pulsar Magnetospheres	1
1.2 Magnetic Reconnection	7
1.3 Particle-In-Cell Method	11
Chapter 2: High-Energy Radiation and Ion Acceleration in Three-dimensional Relativistic Magnetic Reconnection with Synchrotron Cooling	17
2.1 Abstract	17
2.2 Introduction	18
2.3 Numerical Methods	21
2.3.1 Initial configuration and boundary conditions	21
2.3.2 Radiation	25
2.4 Structure of the Current Sheet	27
2.5 Particle acceleration	34
2.5.1 Ion acceleration	38
2.5.2 Pair acceleration	43
2.6 Radiation from the current sheet	52
2.7 Discussion	57
2.7.1 Acceleration of ions in pulsar magnetospheres	58
2.7.2 High-energy emission from young pulsars and supermassive black holes	63
Chapter 3: Coherence of Multi-Dimensional Pair Production Discharges in Polar Caps of Pulsars	66
3.1 Abstract	66
3.2 Introduction	67

3.3	Overview of the physical model	70
3.3.1	Dynamics of the electromagnetic fields	70
3.3.2	Magnetospheric current distribution	72
3.3.3	QED pair production	74
3.3.4	Atmosphere	77
3.3.5	Initial plasma state	78
3.3.6	Numerical details	79
3.4	Results	80
3.4.1	Space-Charge-Limited Flow	81
3.4.2	Ruderman-Sutherland Gap	90
3.4.3	3D Gap, Space-Charge-Limited Flow	92
3.4.4	Evolution of the Twist of Magnetic Field Lines	95
3.5	Conclusions and Discussion	98
3.5.1	Non-existence of sparks	98
3.5.2	Implication for discharges in old pulsars	101
3.5.3	On repetition rate of the cascades	101
Chapter 4:	Origin of the Pulsed 20 TeV Emission in the Vela Pulsar	103
4.1	Abstract	103
4.2	Introduction	103
4.3	Magnetospheric Current Sheet	106
4.3.1	Plasma Feeding from the Polar Cap	106
4.3.2	Particle Acceleration and Synchrotron Radiation from the Current Sheet	108
4.4	Model of the HE and VHE Emission	111
4.4.1	Pair Production	113
4.4.2	Steady-state Pair Spectrum	116
4.4.3	Secondary Pair Synchrotron Spectrum	118
4.4.4	Inverse Compton Signal	119
4.5	Conclusion and Discussion	121
Chapter 5:	Conclusion and Future Work	124
5.1	Summary	124
5.2	Future Work	125
Appendix A:	Appendices for Chapter 2	128
A.1	Varying the fraction of ions and the mass ratio	128
A.2	Varying the resolution	130
Appendix B:	Appendices for Chapter 3	132
B.1	Additional Details on radiation and QED physics	132

List of Tables

2.1	Summary of Simulation Parameters for Current Sheets	25
3.1	Summary of Simulation Parameters for Discharge Simulations.	77

List of Figures

1.1	Schematic of the pulsar magnetosphere	3
1.2	Schematic of the magnetic reconnection	8
1.3	Schematic of the PIC algorithm	14
2.1	Volume Rendering of the Current Sheet	28
2.2	Distribution of Perpendicular Sizes of the Current Sheet	29
2.3	Time Evolution of the Reconnection Rate	31
2.4	The Internal Structure of a Plasmoid	33
2.5	Distribution Functions of Pairs and Ions	36
2.6	Acceleration of Particles in X-points	37
2.7	Trajectories and Energy Gain of Ions	39
2.8	Time Evolution of Sample of Ions	41
2.9	Time Evolution of a Typical High-Energy Electron	45
2.10	Close-up View of an Electron Acceleration Dynamics	47
2.11	Distribution of the Effective Perpendicular Magnetic Field	50
2.12	Distribution of the Lifetimes of Strongly Cooled Pairs	51
2.13	Radiation Spectra from Current Sheets	53
2.14	Energy-dependent Angular Anisotropy of the Radiation	55
2.15	Spatial Distribution of Radiating Regions	55
2.16	Estimates for Burnoff Limit of Protons in Pulsar Magnetospheres	61
3.1	Distribution of the magnetospheric currents	73
3.2	SCLF discharge with small gap	83
3.3	SCLF discharge, different gap models	87
3.4	RS discharge, different gap sizes	91
3.5	3D SCLF discharge	94
4.1	Plasma supply to the current sheet	107
4.2	Particle and photon spectra from a current sheet	110
4.3	Particle and photon spectra in Vela	117
5.1	Vlasov simulation of the polar discharge	127
A.1	Energy Spectra for Different Fractions of Ions	129
A.2	Comparison of Different Resolutions for Current Sheets	131

B.1 1D model of SCLF discharge	136
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List of Abbreviations

CSS	Current Sheet Synchrotron
FFE	Force-Free Electrodynamics
GJ	Goldreich-Julian
HE	High Energy
H.E.S.S.	High Energy Stereoscopic System
IC	Inverse Compton
LC	Light Cylinder
NS	Neutron star
PC	Polar Cap
PIC	Particle-In-Cell
QED	Quantum Electrodynamics
RS	Ruderman-Sutherland
SCLF	Space-Charge-Limited Flow
SPS	Secondary Pairs Synchrotron
VHE	Very High Energy

Chapter 1: Introduction

1.1 Radiation from Pulsar Magnetospheres

Pulsars are lighthouses of the Universe. Being initially discovered as sources of extremely stable periodic bursts of bright radio emission ([Hewish et al., 1968](#)), they were quickly linked to rapidly rotating neutron stars: compact astrophysical objects with characteristic radii of 12 km and masses of $1.4 - 2.2M_{\odot}$, and typical rotational periods $P \approx 10 \text{ ms} - 10 \text{ sec}$. The next $\gtrsim 55$ years of observations show that pulsars are broad-band emitters of non-thermal radiation, covering about 20 orders of magnitude in photon energy (see, e.g., [Philippov & Kramer, 2022](#)). The studies of pulsar magnetospheres were revolutionized by the detection of their γ -ray emission using *Fermi*, with more than 200 pulsars detected until now ([Abdollahi et al., 2020](#)): for the first time, observations opened a window into an energetically significant process operating in pulsar magnetospheres.

Early theoretical models of pulsar magnetospheres were based on a model of a vacuum magnetic dipole, μ , rotating with angular velocity Ω ([Pacini, 1967](#)), which radiates its rotational energy through the Poynting flux equivalent to a luminosity $L = (2/3)(\mu^2\Omega^4)/(c^3) \sin^2 \chi$ ([Landau & Lifshitz, 1975](#)), where χ is an angle between the rotational axis and the magnetic dipole, c is the speed of light. Poynting flux is supported by the rotation-induced electric field, which is quadrupolar for the dipolar magnetic field. Such an electric field has a parallel to the magnetic

field component that is able to extract charged particles from the stellar surface and accelerate them to enormous energies. The fatal flaw of the vacuum model is its inconsistency with observations, requiring charged particles to produce the observed emission. The presence of plasma leads to modifications of the vacuum fields; in particular, it leads to a screening of the parallel component of the electric field. Under the assumption of a plasma-filled magnetosphere, magnetic field lines loaded with plasma cannot corotate with the neutron star beyond a certain radius. This distance is called the light cylinder radius and is determined as a distance where the corotational speed is equal to the speed of light, $R_{\text{LC}} = c/\Omega$. While the magnetic field lines that do not cross the light cylinder remain mostly not distorted, the ones that cross have to change their geometry. This condition separates the pulsar magnetosphere into two regions: the zone of the closed field lines and the zone of the open field lines. These two zones and the light cylinder are shown in the sketch 1.1.

Open field lines originate in a small region on the stellar surface called the polar cap. Its characteristic size is $R_{\text{PC}} \approx R_{\star} \sqrt{R_{\star}/R_{\text{LC}}}$, where $R_{\star}$ is a stellar radius. The exact size and shape of the polar cap depends on the mutual orientation of the magnetic $\boldsymbol{\mu}$ and rotational $\boldsymbol{\Omega}$ axes. The field-aligned flow of charged particles extracted from the polar cap produces an electric current, I . It generates the twist of the magnetic field, resulting in a toroidal component of the magnetic field, $B_{\varphi} = 2I(r_{\perp})/cr_{\perp}$. This toroidal magnetic field together with corotational electric field $\boldsymbol{E}_{\text{cor}}$ create an outward component of the Poynting flux, $(c/4\pi)(\boldsymbol{E}_{\text{cor}} \times \boldsymbol{B}_{\varphi})$. The corotational electric field, $\boldsymbol{E}_{\text{cor}} = -\boldsymbol{\Omega} \times \boldsymbol{r} \times \boldsymbol{B}/c$, is realized by a minimally required charge density, $\nabla \cdot \boldsymbol{E}_{\text{cor}} \approx 4\pi\rho_{\text{GJ}}$, where, $\rho_{\text{GJ}} = -\boldsymbol{\Omega} \cdot \boldsymbol{B}/(2\pi c)$ is the Goldreich-Julian charge density (Goldreich & Julian, 1969). Integration of the Poynting flux through the light cylinder results in an estimate of the energy

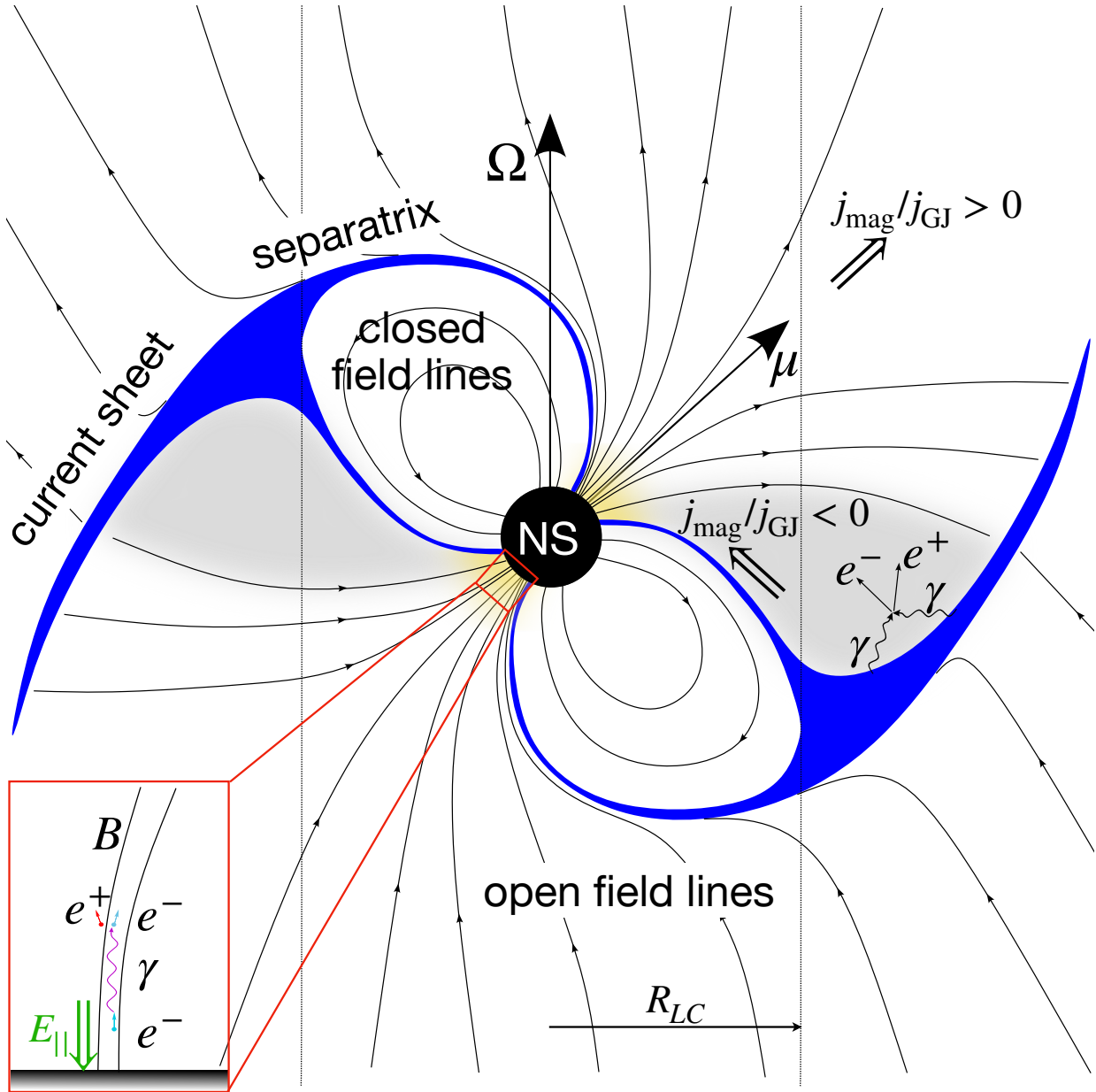


Figure 1.1: Sketch of the magnetosphere of the 45°-inclined pulsar. The main elements of the magnetosphere are shown: zones of open and closed field lines, magnetospheric current sheet, and separatrix. The shadowed region of the open field lines zone highlights the field lines carrying volume return current. The zoom-in shows the scheme of the polar discharge.

losses

$$L \approx \frac{\mu^2 \Omega^4}{c^3} \sim 10^{35} \text{ erg/s}, \quad (1.1)$$

for typical surface magnetic field strength, $B = 10^{12} \text{G}$, and rotational period, $P = 0.1 \text{s}$. This estimate turns out to be close to the vacuum dipolar losses, with the main difference being that the losses are now non-zero for the aligned rotator as well. Since the Poynting flux carries away the rotational energy of the star, it produces a spindown, $\dot{E} = -I\Omega\dot{\Omega} = 4\pi^2 I \dot{P}/P^3$. For a moment of inertia of the star, $I = 10^{45} \text{g} \times \text{cm}^2$, the increase of the rotational period is small, $\dot{P} \sim 10^{-15}$, but observable and provides an important constraint on the strength of the magnetic field. The luminosity of the radio emission, which pulsars are known for, accounts only for a small fraction of the total spindown luminosity: $L_{\text{radio}} \sim 10^{-7} - 10^{-2} L_{\text{spindown}}$ (Lorimer & Kramer, 2005). Its brightness temperature¹, however, is incredibly high, $T_{\text{b}} \sim 10^{25} - 10^{40} \text{K}$ (Hankins & Eilek, 2007), requiring a coherent emission mechanism.

In the Goldreich-Julian model, the accelerating electric field is screened by the plasma extracted from the stellar surface, and no particle acceleration occurs. The model was modified by Sturrock (1971), who assumed that the non-zero accelerating electric field exists in the magnetosphere. It leads to the production of additional pair plasma via quantum electrodynamics (QED) effects. The driver behind this is a huge electric potential drop across the polar cap, $V = E_{\text{rot}} R_{\text{PC}}$. Electric fields lead to efficient particle acceleration along the curved magnetic field lines and subsequent emission of high-energy curvature photons with the energies $\varepsilon_{\text{curv}} = (3/2)\hbar(c/\rho)\gamma^3$, where ρ is the curvature radius of the magnetic field line. These photons, emitted initially along the magnetic field lines, propagate along a straight line, and the angle between the wave vector

¹The brightness temperature is the temperature needed for a blackbody radiator to produce the same intensity as the observed source.

and the magnetic field lines increases. As a result, the probability of the photon decay into an electron-positron pair via a single photon conversion process increases, $\gamma + B \rightarrow e^- + e^+$ (Erber, 1966). The sketch of this process is shown in the inset of Figure 1.1. The produced pairs have a non-zero pitch angle and intensively radiate transverse energy into synchrotron photons which are again converted into new electron-positron pairs. During the acceleration process, a primary charged particle emits multiple energetic photons, and each of them is converted into pairs. As a result, the magnetosphere is filled with a high-density plasma with number density $n = \lambda n_{\text{GJ}}$, which screens the parallel electric field. Here, $n_{\text{GJ}} = \rho_{\text{GJ}}/e$ is the characteristic density of particles extracted from the stellar atmosphere (Beloborodov, 2008; Timokhin & Arons, 2013). The dimensionless parameter $\lambda \gg 1$ is called the multiplicity and characterizes how many secondary particles can be produced by a single primary particle accelerated in the gap – the zone where the unscreened electric field could exist. The multiplicity can be estimated from energy conservation: estimates for Lorentz-factors of the primary particles, $\gamma_{\text{th}} \sim 10^6 - 10^7$, and of the secondary ones, $\gamma_{\text{sec}} \sim 10^2 - 10^3$ (Timokhin & Harding, 2015), result in the multiplicity $\lambda \sim \gamma_{\text{th}}/\gamma_{\text{sec}} \sim 10^4 - 10^5$. The process of particle acceleration with subsequent secondary plasma creation and screening of the electric field is called a discharge. It is considered the primary source of pair plasma in the magnetosphere.

Ruderman & Sutherland (1975) developed these ideas further and produced the first quantitative model of the pair discharge. They demonstrated that the position of the pulsars’ “death line” corresponds to the potential drop across the polar cap, $\Phi \approx 10^{12}\text{V}$. At lower electric potential, primary particles are not accelerated to high enough energy to ignite the pair-production discharge, and the pulsar is no longer able to produce radio emission. Initially, the discharge was considered as a stationary process (e.g., Ruderman & Sutherland, 1975; Arons & Scharlemann,

1979). Later, using first-principles plasma simulations, [Timokhin \(2010\)](#) demonstrated that the discharge proceeds intermittently: intense bursts of pair-production are followed by quiet moments when the QED-produced plasma screens an electric field. After the plasma leaves the gap zone, the unscreened electric field forms again, and the process restarts. The cyclic behavior of the discharge provides a natural explanation for the observed microstructure of individual radio pulses. Recently, the process of cyclic screening of the electric field has been directly connected to the formation of the radio emission (e.g., [Philippov et al., 2020](#); [Melrose et al., 2021a](#)). In this Thesis, I develop the first self-consistent multi-dimensional models of the pair discharge.

In the plasma-filled magnetosphere, the bundles of the open field lines from two opposite polar caps cross the light cylinder and approach each other. Since their magnetic field lines have opposite polarities, a discontinuity appears in the region of contact of these two fields - a magnetospheric current sheet (shown in blue in Figure 1.1). This current sheet undergoes magnetic reconnection and is the site of the efficient particle acceleration and subsequent γ -ray emission ([Lyubarskii, 1996](#); [Cerutti et al., 2015](#); [Philippov & Spitkovsky, 2018](#)). Other possible sites of the γ -ray emission are possible in the magnetosphere: the escaping high-energy photons from the gap region close to the stellar surface, and the curvature radiation of the energetic particles propagating along the boundary of open and closed field lines, the separatrix, carrying intense currents ([Contopoulos & Kalapotharakos, 2010](#)). Global magnetospheric simulations show that the energy release in the inner magnetosphere is insignificant ([Hu & Beloborodov, 2022](#); [Hakobyan et al., 2023a](#)). Observations rule out high-energy radiation from near the surface region: data show no super-exponential cutoff of the spectrum at high energies expected from single-photon conversion. The mismatch of rotational phases at which the radio and γ -ray light curves peak (e.g., [Kuiper & Hermsen, 2015](#)) also disfavors the inner magnetosphere as the source of the observed

high-energy radiation (the Crab pulsar and some other very energetic pulsars are an exception, [Eilek & Hankins 2016](#)). At the same time, *Fermi* observations demonstrate that the high-energy emission contains $\sim 10\%$ of the total pulsar spindown luminosity ([Abdo et al., 2010a,b, 2013](#); [Ansoldi et al., 2016a](#)), pointing out to the magnetic reconnection as the particle energization process. Emission produced near and beyond the light cylinder accumulates in caustics: all of the photons, emitted by relativistic particles propagating to the observer, travel together with an emitting particle and arrive to an observer simultaneously, leading to the pile-up of the photons in the skymap ([Bai & Spitkovsky, 2010](#); [Kalapotharakos et al., 2014](#)). This caustic results in narrow peaks in γ -ray lightcurves, consistent with observations.

1.2 Magnetic Reconnection

Magnetic reconnection is believed to be the most efficient mechanism of converting magnetic field energy into plasma energy and is accompanied by the change of the magnetic field topology schematically shown in Fig. 1.2a. Magnetic reconnection events were directly observed near the surface of the Sun, in the Solar corona, and also in the Earth magnetotail ([Yamada et al., 2010](#)). It appears plausible that the reconnection can occur in remote astrophysical systems as well, although its presence can not be demonstrated directly. Around compact astrophysical objects, neutron stars, and black holes, magnetic reconnection occurs in highly magnetized environments with $\beta_p \equiv 8\pi P_{\text{th}}/B^2 \ll 1$, and transfer of the magnetic field energy into particles significantly affects their dynamics.

The switch of the field polarity necessitates the presence of a large current confined in the current sheet. The characteristic width of the current sheet in magnetized pair plasma is given by

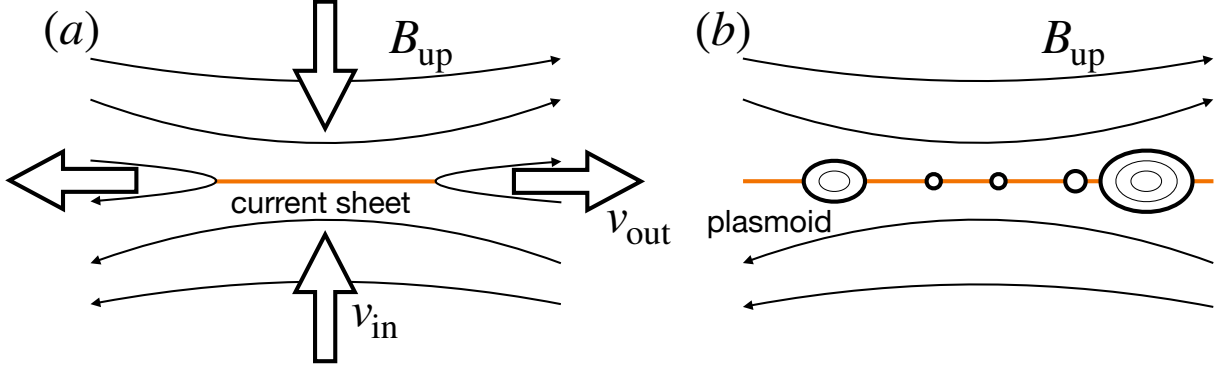


Figure 1.2: Sketch of the current sheet formation. Subplot (a) shows the formation of a short Sweet-Parker-like current sheet. Arrows indicate the directions of the plasma inflow and outflow and the magnetic field orientation. Subplot (b) shows a long plasmoid-unstable current sheet with a developed plasmoid chain.

a skin depth, $d_e = c/\omega_{pl,e}$ where $\omega_{pl,e} = \sqrt{4\pi ne^2/m_e}$ is the plasma frequency. At smaller scales, the magnetic field is no longer frozen into the plasma, and the rearrangement of the magnetic field and consecutive energy release are spontaneously triggered.

The first analytical steady-state model of reconnection was proposed by [Parker \(1957\)](#) and [Sweet \(1958\)](#), the Sweet-Parker model. In their model, two opposite polarities of the magnetic field are advected into the current sheet with a constant speed, v_{in} . In the steady-state, the associated motional electric field, $\mathbf{E} = -\mathbf{v}_{in} \times \mathbf{B}/c$, is constant everywhere, including inside the current sheet, due to the continuity of the tangential electric field at the boundaries of the upstream zones and the current sheet. The inflow and outflow mass fluxes are equal, resulting in $n\delta v_{out} = nLv_{in}$. Here, n is the plasma density, δ is the half-width of the current sheet, and L is the length of the interface of the opposite magnetic field regions. The dimensionless parameter $r = v_{in}/v_{out} = \delta/L$ is called the “reconnection rate” and characterizes the efficiency of advection and dissipation of the magnetic field. The outflow velocity, v_{out} , can be found from energy conservation under the assumption that all of the magnetic energy is converted into the kinetic

energy of plasma, $B^2/8\pi = nm_e v_{\text{out}}^2/2$. Hence, the outflow velocity is equal to the Alfven velocity, $v_{\text{out}} = B/\sqrt{4\pi\rho} \equiv v_A$. For the plasma resistivity η , Ohm's law determines the thickness of the current sheet, $J \approx B/\delta \approx v_{\text{in}}B/(\eta c)$. These results determine the reconnection rate as

$$r = \frac{1}{\sqrt{S}}, \quad S = \frac{v_A L}{\eta} - \text{Lundquist number.} \quad (1.2)$$

This fact reveals the main drawback of the Sweet-Parker theory. For astrophysical systems, the Lundquist number is huge, $S \sim 10^{20}$, leading to an extremely low reconnection rate and, hence, inefficient particle acceleration. However, the current sheets of such high Lundquist numbers are always unstable. [Furth et al. \(1963\)](#) showed that the thinning of the current sheet leads to the growth of the tearing instability. [Biskamp \(1986\)](#); [Loureiro et al. \(2007\)](#) demonstrated that long enough current sheets, with $S \gtrsim 10^4$, have to split into smaller ones on the dynamical timescales. As a result of continuous fragmentation into the magnetic islands, or plasmoids, current sheets maintain multiple evolving reconnection sites (a sketch of the plasmoid-unstable current sheet is shown in Fig. 1.2b). These plasmoids were directly observed in flares in the Solar corona ([Lin et al., 2005](#)) and in the Earth magnetotails ([Zong et al., 2004](#)). The fragmentation of a long current sheet marks the “fast reconnection” regime when the reconnection rate becomes independent of the length of the current sheet and is about 10%, $r \approx 0.1$, for collisionless plasmas (e.g., [Zenitani & Hoshino, 2008](#); [Bessho & Bhattacharjee, 2012](#)).

Similar conclusions hold for reconnection in collisionless relativistic plasmas ([Lyubarsky, 2005](#)). Here, the ratio of the supplied magnetic energy and plasma energy is characterized by a

dimensionless magnetization parameter:

$$\sigma = \frac{B^2}{4\pi n m_e c^2}. \quad (1.3)$$

For magnetized environments of black holes and neutron stars, $\sigma \gg 1$, making reconnection proceed in the relativistic regime. For these conditions, the Alfven velocity reaches the speed of light, $v_A \approx c\sqrt{\sigma/(\sigma+1)} \approx c$, and the reconnection rate can be defined as $r \approx \beta_{\text{rec}} \equiv v_{\text{in}}/c$. Following the conversion of the magnetic field energy into plasma energy at reconnection sites, the characteristic energies of the accelerated particles become comparable to the magnetization parameter, $E \sim m_e c^2 \sigma$, and in the relativistic regime particles form a hard non-thermal power-law distribution (e.g., [Guo et al., 2014](#); [Sironi & Spitkovsky, 2014](#)): $f(E) \propto E^{-1}$. Subsequent acceleration to energies exceeding the value of the magnetization parameter happens only for a small fraction of particles, and three-dimensional simulations show steepening of the distribution function, $f(E) \propto E^{-2}$ ([Zhang et al., 2021a](#)).

The formation of current sheets is ubiquitous in astrophysical systems. In some of them, it is determined by the global geometries of the magnetic field, e.g., magnetospheres of pulsars (e.g., [Spitkovsky, 2006](#); [Philippov et al., 2015](#); [Cerutti et al., 2020](#); [Hakobyan et al., 2023a](#)) and accreting black holes (e.g., [Ripperda et al., 2020](#); [Galishnikova et al., 2023](#)). Additionally to global current sheets, localized reconnection events can occur in magnetized relativistic turbulence in both collisional ([Zhdankin et al., 2013](#); [Chernoglazov et al., 2021](#); [Galishnikova et al., 2022](#)) and collisionless (e.g., [Comisso & Sironi, 2018, 2019](#)) regimes, due to the scale-dependent intermittency of turbulence ([Boldyrev, 2006](#); [Mallet et al., 2015](#); [Chandran et al., 2015](#)). In some astrophysical systems, the system-sized current sheet forms intermittently and is believed to drive

powerful flares in erupting magnetospheres of magnetars (e.g., [Beloborodov, 2017a](#); [Mahlmann et al., 2023](#)), kink-unstable relativistic jets (e.g., [Davelaar et al., 2020](#)), black holes surrounded by magnetically-arrested disks ([Ripperda et al., 2022](#)), and in coronae of x-ray binaries and Active Galactic Nuclei ([Beloborodov, 2017b](#)). In this Thesis, I investigate three-dimensional reconnection in conditions appropriate for pulsar magnetospheres.

In the vicinity of compact astrophysical objects, electromagnetic radiation emitted by accelerated particles is so powerful that the radiation reaction force significantly alters the dynamics of particles and, thus, of reconnection itself ([Uzdensky, 2016](#)). In pulsar magnetospheres and, likely, in magnetospheres of black holes, electron-positron pairs created by high-energy photons significantly affect the mass loading. Thus, self-consistent predictions of the high-energy emission from reconnecting current sheets require the inclusion of radiative effects into kinetic simulations of relativistic reconnection. The 2D regime of radiative reconnection has been thoroughly explored in recent years ([Jaroschek & Hoshino, 2009](#); [Uzdensky & McKinney, 2011](#); [Cerutti et al., 2012b,a](#); [Beloborodov, 2017b](#); [Werner et al., 2019](#); [Sironi & Beloborodov, 2020](#); [Mehlhaff et al., 2020](#); [Sridhar et al., 2021](#)). However, little progress has been made in studying three-dimensional reconnection, incorporating at least some of the radiative effects self-consistently, which motivates my explorations in Chapter 2.

1.3 Particle-In-Cell Method

The dynamics of collisionless plasmas in many astrophysical systems are determined by the self-consistent interplay between the charged particles and electromagnetic fields. This evolution

is described by the Vlasov equation (e.g., [Lifshitz & Pitaevskii, 1981](#))

$$\frac{\partial f_s}{\partial t} + \mathbf{v} \cdot \nabla f_s - \frac{q_s}{m_s} \left(\mathbf{E} + \frac{\mathbf{v} \times \mathbf{B}}{c} \right) \cdot \frac{\partial f_s}{\partial \mathbf{v}} = 0, \quad (1.4)$$

which is a kinetic equation for the 6D distribution function $f(\mathbf{x}, \mathbf{v}, t)$ of charged particles of specie s , with charge q_s and mass m_s . $\mathbf{E}$ and $\mathbf{B}$ are the coarse-grained electric and magnetic fields consistent with the charge density and electric currents created by the charged particles in the system. The evolution of the electromagnetic field is described by Maxwell's equations:

$$\frac{1}{c} \frac{\partial \mathbf{B}}{\partial t} = -\nabla \times \mathbf{E}, \quad \frac{1}{c} \frac{\partial \mathbf{E}}{\partial t} = \nabla \times \mathbf{B} - \frac{4\pi}{c} \mathbf{j}, \quad (1.5)$$

and satisfies the constraints $\nabla \cdot \mathbf{E} = 4\pi\rho$ and $\nabla \cdot \mathbf{B} = 0$. Here, the charge, ρ , and current, $\mathbf{j}$, densities are given by the distribution function f_s as

$$\rho = \sum_s q_s \int f_s d^3\mathbf{v}, \quad \mathbf{j} = \sum_s q_s \int \mathbf{v} f_s d^3\mathbf{v}. \quad (1.6)$$

Equations (1.4-1.5) for a collisionless plasma form a set of hyperbolic evolution equations with a closure (1.6). These equations can rarely be solved analytically in the non-linear regime, requiring the adoption of numerical techniques. The straightforward numerical solution of the kinetic equation (1.4) requires evolution of the distribution function in 6-dimensional space and is extremely computationally expensive. The most widely applied method for solving the Vlasov equations is the Particle-In-Cell (PIC) algorithm, which I employ in this Thesis.

The PIC algorithm ([Birdsall & Langdon, 1991](#)) discretizes the distribution function by sampling it with a finite number of macro-particles. These particles can be considered as a large

collection of real particles with mass m_s and charge q_s . At any given moment t , the distribution function is expressed as a sum over the sampling macro-particles

$$f_s(\mathbf{x}, \mathbf{v}, t) = \sum_i S(\mathbf{x} - \mathbf{x}_i(t))\delta(\mathbf{v} - \mathbf{v}_i(t)), \quad (1.7)$$

where $\mathbf{x}_i(t)$, $\mathbf{v}_i(t)$ are position and velocity of a macro-particle i at the moment of time t . While the macro-particle has a determined velocity $\mathbf{v}_i(t)$, its spatial distribution is extended and is described by a shape-function S . For the relativistic plasmas, further discussed in the Thesis, the equation of motion of a separate particle can be conveniently described by the evolution of the 4-velocity $\mathbf{u}_i = \mathbf{v}_i\gamma_i$. Here, $\gamma_i = 1/\sqrt{1 - v_i^2/c^2}$ is the Lorentz factor of a particle. Thus, the equation of motion of each particle is given by

$$\frac{d\mathbf{x}_{s,i}}{dt} = \frac{\mathbf{u}_{s,i}}{\gamma_{s,i}}(t), \quad (1.8)$$

$$\frac{d\mathbf{u}_{s,i}}{dt} = \frac{1}{m_s}\mathbf{F}_{s,i}. \quad (1.9)$$

Consistency with the Vlasov equation (1.4) requires spatial averaging of the Lorentz force, $\mathbf{F}_{s,i}$, using the same shape function, S :

$$\mathbf{F}_{s,i} = q_s \left(\mathbf{E}(\mathbf{r}_{s,i}) + \frac{\mathbf{u}_{s,i} \times \mathbf{B}(\mathbf{r}_{s,i})}{\gamma_{s,i}c} \right). \quad (1.10)$$

While the macroparticles sample the plasma distribution function, the electric and magnetic fields are sampled on a spatial grid. The scheme of such a grid is shown in Fig.1.3a. The staggered positions of the field components in the cell are chosen to satisfy the divergence-free constraint

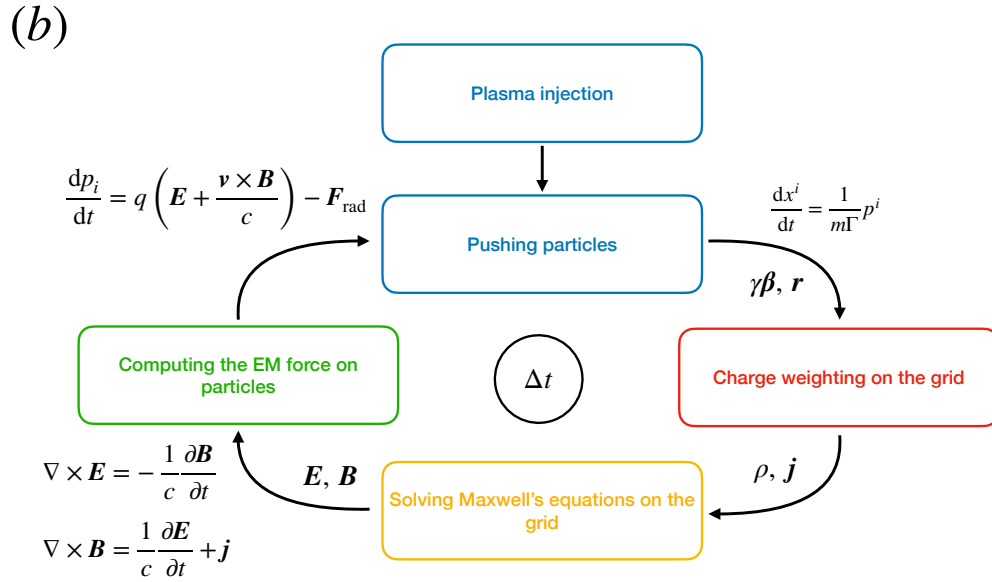
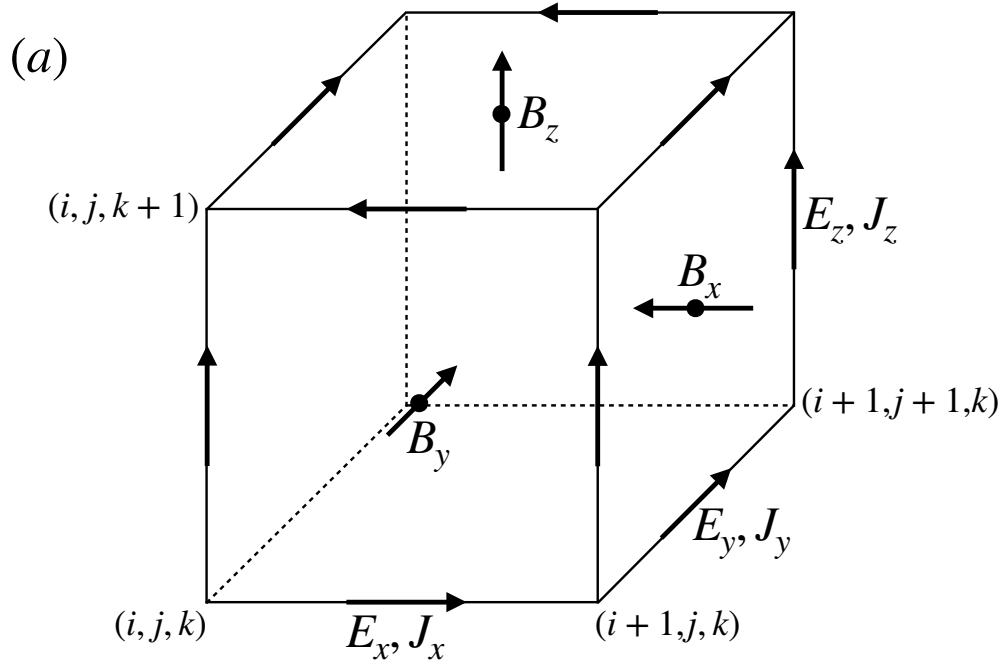


Figure 1.3: Schematic representation of the PIC algorithm: Subplot (a) shows the sampling of the fields and electric currents on the grid (Yee, 1966). Subplot (b) shows the key steps performed by a PIC algorithm at each timestep.

to the machine precision and discretize the spatial derivatives with the second-order numerical accuracy (Yee, 1966). The particle current is collected from macroparticles via the current deposit scheme, which can be done to satisfy the charge-conservation equation to the machine's precision. Further description of the method will be based on PIC code `Tristan v2` (Hakobyan et al., 2023), which I used for simulations in this Thesis. The scheme of the algorithm is presented in Fig. 1.3b.

The main steps of the PIC algorithm are:

1. Computation of the change of particle 4-velocity during the timestep due to the Lorentz force, Eqn. (1.10), and the corresponding coordinate update, using a leap-frog algorithm (Boris, 1970).
2. Deposit of the particle electric currents $\mathbf{j}$ onto the grid. This process needs to preserve the charge conservation constraint, $\nabla \cdot \mathbf{E} = 4\pi\rho$. `Tristan v2` code uses the zigzag method, which conserves charge with machine precision (Umeda et al., 2003).
3. Computation of the time evolution of the electromagnetic field, $\mathbf{E}$, $\mathbf{B}$, governed by Maxwell equations, using the grid-based values of the fields and the electric current. As with most other PIC codes, `Tristan v2` adopts the leap-frog scheme for the time update. In Chapter 3, I adopt higher-order spatial discretization (Blinne et al., 2018) in order to improve the dispersive properties of the field solver.
4. Interpolation of the updated electromagnetic field from the grid into the position of each particle. Here, similar to the current deposit, the spatial non-locality of the particles is used.

Sampling of the continuous distribution function by the finite number of macroparticles leads

to the introduction of the numerical shot noise, which often results in parasitic heating of the plasma (e.g., [Shalaby et al., 2017](#)). There are two main ways to reduce it: to increase the number of computational particles or to use a smoother shape function ([Birdsall & Langdon, 1991](#)). Since the first method is numerically expensive and only slowly reduces the noise, `Tristan v2` employs the second method by applying digital filter passes on the deposited currents with a $(1/4, 1/2, 1/4)$ stencil.

In order to model relativistic radiative plasmas in pulsar magnetospheres, I have developed a synchrotron emission module, which takes into account the synchrotron kernel, and the complete QED pair production scheme that models the single-photon decay process, $\gamma + B \rightarrow e^- + e^+$, into `Tristan v2`. I describe details of these numerical implementations in [Chapters 2 and 3](#) and corresponding [Appendices](#).

Chapter 2: High-Energy Radiation and Ion Acceleration in Three-dimensional Relativistic Magnetic Reconnection with Synchrotron Cooling

2.1 Abstract

We present the results of 3D particle-in-cell (PIC) simulations that explore relativistic magnetic reconnection in pair plasma with strong synchrotron cooling and a small mass fraction of non-radiating ions. Our results demonstrate that the structure of the current sheet is highly sensitive to the dynamic efficiency of radiative cooling. Specifically, stronger cooling leads to more significant compression of the plasma and magnetic field within the plasmoids. We demonstrate that ions can be efficiently accelerated to energies exceeding the plasma magnetization parameter, $\gg \sigma$, and form a hard power-law energy distribution, $f_i \propto \gamma^{-1}$. This conclusion implies a highly efficient proton acceleration in the magnetospheres of young pulsars. Conversely, the energies of pairs are limited to either σ in the strong cooling regime or the radiation burnoff limit, γ_{syn} , when cooling is weak. We find that the high-energy radiation from pairs above the synchrotron burnoff limit, $\varepsilon_c \approx 16$ MeV, is only efficiently produced in the strong cooling regime, $\gamma_{\text{syn}} < \sigma$. In this regime, we find that the spectral cutoff scales as $\varepsilon_{\text{cut}} \approx \varepsilon_c(\sigma/\gamma_{\text{syn}})$, and the highest energy photons are beamed along the direction of the upstream magnetic field, consistent with the phenomenological models of gamma-ray emission from young pulsars. Furthermore, our results

place constraints on the reconnection-driven models of gamma-ray flares in the Crab Nebula.

2.2 Introduction

Over the past two decades, the physics of magnetic reconnection in the magnetospheres of compact astrophysical objects, such as neutron stars and black holes, has been extensively studied both theoretically and numerically. In particular, the key role of reconnection in the process of magnetic energy dissipation and particle acceleration has been widely recognized and accepted (Lyubarskii, 1996; Lyubarsky, 2005; Uzdensky et al., 2010; Sironi & Spitkovsky, 2014; Uzdensky & Spitkovsky, 2014; Uzdensky & Loureiro, 2016; Uzdensky, 2016; Sironi et al., 2016; Werner & Uzdensky, 2017). In some systems, the presence of large-scale reconnection sites has been demonstrated by global simulations, such as pulsar (e.g., Philippov & Spitkovsky, 2014; Chen & Beloborodov, 2014) and black hole magnetospheres (e.g., Parfrey et al., 2019; Bransgrove et al., 2021; Ripperda et al., 2022). In other systems, such as the coronae of X-ray binaries or jets of active galactic nuclei (AGNs), the formation of intermittent current layers associated with either turbulent energy cascade or MHD instabilities (e.g., kink instability, Begelman, 1998; Tchekhovskoy & Bromberg, 2016; Alves et al., 2018; Davelaar et al., 2020; Ortuño-Macías et al., 2022) has been theorized (Beloborodov, 2017b) or directly demonstrated from intermediate scale simulations (Servidio et al., 2011; Zhdankin et al., 2013, 2017, 2018; Comisso & Sironi, 2018; Davelaar et al., 2020; Chernoglazov et al., 2021; Sironi et al., 2021). Relativistic magnetic reconnection has thus established itself as one of the most important plasma-physical processes powering energy dissipation and non-thermal emission in various contexts of high-energy astrophysics.

Strong radiative cooling, both due to synchrotron emission and inverse-Compton scatterings, has implications for particle acceleration and the reconnection process itself (see, e.g., [Uzdensky, 2016](#)). The dynamics of the radiative reconnection has so far been mostly studied for 2D current sheets ([Jaroschek & Hoshino, 2009](#); [Uzdensky & McKinney, 2011](#); [Cerutti et al., 2012b,a](#); [Beloborodov, 2017b](#); [Werner et al., 2019](#); [Sironi & Beloborodov, 2020](#); [Mehlhaff et al., 2020](#); [Sridhar et al., 2021](#)). For the case of strong cooling (either synchrotron or inverse-Compton), radiative losses lead to the decrease of the plasma temperature, resulting in a strong compression of plasmoids (e.g., [Schoeffler et al., 2019](#); [Hakobyan et al., 2019](#)). Radiative cooling also substantially limits non-thermal particle acceleration. However, in the case of strong synchrotron cooling, the initial acceleration (“injection”) of particles takes place in or close to the X-points, where the radiative cooling is negligible ([Cerutti et al., 2014](#); [Kagan et al., 2016a](#); [Hakobyan et al., 2019](#)), allowing the particles to accelerate above the synchrotron “burn-off” limit and emit γ -ray synchrotron photons ([Uzdensky et al., 2011](#); [Cerutti et al., 2012a](#); [Kagan et al., 2016b](#)).

The case of 3D relativistic current sheets is substantially less explored ([Zenitani & Hoshino, 2008](#); [Guo et al., 2014](#); [Sironi & Spitkovsky, 2014](#); [Werner & Uzdensky, 2017](#); [Zhang et al., 2021a](#); [Werner & Uzdensky, 2021](#); [Guo et al., 2021](#); [Schoeffler et al., 2023](#)). It has been shown that there are competing instabilities and acceleration mechanisms specific to 3D. For example, [Zenitani & Hoshino \(2008\)](#), and [Cerutti et al. \(2014\)](#) demonstrated that in the linear stage, the relativistic drift-kink instability grows faster than the plasmoid instability. Later, [Sironi & Spitkovsky \(2014\)](#), and [Guo et al. \(2014\)](#) showed that the plasmoid instability still dominates in the non-linear regime, leading to the plasmoid-dominated fast reconnection stage in 3D current sheets and efficient particle acceleration. The 3D plasmoids (or flux tubes) resulting from the plasmoid instability are subject to MHD-kink instability, which ultimately leads to turbulence,

the nature and properties of which are largely unexplored (Huang & Bhattacharjee, 2016; Guo et al., 2021). The mechanisms for particle acceleration are also substantially different between 2D and 3D: while particle “injection,” i.e. acceleration to Lorentz factors comparable to the magnetization parameter, σ , proceeds similarly (e.g., Sironi & Spitkovsky, 2014; Werner & Uzdensky, 2017; Sironi, 2022), the acceleration to high energies in 3D is significantly more efficient. In particular, Dahlin et al. (2015), and Zhang et al. (2021a, 2023) have shown that in 3D, particles are capable of escaping the flux tubes and can be directly accelerated by the reconnection electric field. These studies were mainly focused on the non-radiative regime, and it remains to be seen how radiative cooling affects 3D current sheets (Cerutti et al., 2014; Schoeffer et al., 2023).

Another important aspect of the problem is the composition of current sheets. While most studies of relativistic reconnection focus on pair plasmas, in many scenarios, such as the magnetospheres of pulsars (e.g., Guépin et al., 2020, see Sec. 2.7.1) and supermassive black holes, we expect a mixture of ions (primarily protons) to be present. Importantly, ions are not susceptible to radiative cooling and have the potential to be accelerated more efficiently by reconnection. Thus far this question, in the context of radiative relativistic reconnection, has not been addressed systematically.

In this chapter, we present a set of large 3D PIC simulations of reconnecting current sheets populated by the electron-positron plasma with a small mixture of ions. In our simulations, pairs are subject to dynamically important synchrotron cooling, while ions are unaffected. The goal of our study is to understand the emerging distribution functions of the accelerated ions and pairs, their acceleration mechanisms, and the properties of high-energy radiation. We begin in Sec. 2.3 with a description of the numerical setup we employ. We then describe the structure of radiatively-cooled 3D current sheets and show that those with the strong cooling form more

compressed flux tubes (Sec. 2.4). In Sec. 2.5, we present our main findings on particle acceleration. In particular, we show that ions are very efficiently accelerated in current sheets with strong radiative cooling, while the Lorentz factors of pairs are limited by the upstream magnetization parameter by the synchrotron losses. In Sec. 2.6 we discuss the radiation spectra and the beaming of the emission, for varying strength of the synchrotron cooling. We find that 3D current sheets with efficient cooling are capable of emitting high-energy photons above the synchrotron burn-off limit, which are preferentially beamed along the upstream magnetic field. In contrast, current sheets with weak cooling emit photons with energies up to the burn-off limit in the direction perpendicular to the upstream field direction. We summarize our main findings and discuss their implications for magnetospheres of pulsars and black holes in Sec. 2.7.

2.3 Numerical Methods

For our simulations, we use the `Tristan v2` multi-species radiative PIC code (Hakobyan et al., 2023).

2.3.1 Initial configuration and boundary conditions

The simulation domain is initialized with a single current sheet in Harris equilibrium: the magnetic field profile is given by $\mathbf{B} = B_{\text{up}} \hat{x} \tanh(z/\delta_{\text{cs}})$, i.e., no guide field is present, where x and z are the coordinates along and transverse to the current sheet, correspondingly, and δ_{cs} is the initial half-thickness of the current sheet, and the thermal pressure in the current sheet is set to compensate the gradient of magnetic pressure $B^2/8\pi$. We initialize uniform thermal upstream pair plasma with a mixture of ions of overall number density, $n_{\text{up}} = n_i + n_e + n_{e+} = 2n_e$, and

a low temperature, $T_{\text{up}}/(m_e c^2) = 10^{-4}$. In all simulations presented in the main text, the ions constitute a fraction $f_i = 0.01$ of the total number of particles¹, and we set the ion-to-electron mass ratio to be $m_i/m_e = 1$. This is justified because in the process of reconnection at high pair magnetization, $\sigma = B^2/(4\pi n_e m_e c^2) \gtrsim 10^3$ (for the realistic mass ratio), ions are quickly accelerated to relativistic energies, $\gtrsim \sigma(m_e/m_i) \geq 1$, similarly to pairs (e.g., [Werner et al., 2018](#)). We do not expect a realistic ion-to-electron mass ratio to lead to different results as long as the plasma inertia is dominated by pairs, $n_i m_i/(2n_e m_e) = f_i \cdot (m_i/m_e) \lesssim 1$.² These assumptions are validated in the Appendix [A.1](#), where we vary both the mass ratio (up to $m_i/m_e = 25$) and the fraction of ions, f_i .

The magnetization parameter of the upstream plasma,

$$\sigma \equiv \frac{B_{\text{up}}^2/4\pi}{m_e n_{\text{up}} c^2 \left[(1 - f_i) + f_i \frac{m_i}{m_e} \right]} \approx \frac{B_{\text{up}}^2/4\pi}{m_e n_{\text{up}} c^2}, \quad (2.1)$$

is fixed at $\sigma = 50$. The current sheet is initialized as a hot electron-positron-ion plasma of density $n_{\text{cs}} = N_{\text{cs}} n_{\text{up}} \cosh^{-2}(z/\delta_{\text{cs}})$, and temperature $T_{\text{cs}}/(m_e c^2) = 2\sigma/N_{\text{cs}}$, where $N_{\text{cs}} = 3$ is the maximum overdensity of the layer. Thus, the thermal plasma pressure inside the current sheet balances the pressure of the upstream magnetic field, $B_{\text{up}}^2/8\pi$. The initial thickness of the current layer, $\delta_{\text{cs}} = 7d_e$, is large enough so the tearing instability is not seeded by the numerical noise.

Our simulation domain has outflowing boundary conditions in the $\pm\hat{x}$ -directions, with a thick absorber that gradually damps $\mathbf{E}$ and $\mathbf{B}$ fields to zero towards the edges of the box. In the $\pm\hat{z}$ -direction we implement a moving plasma injector, similar to [Sironi & Spitkovsky \(2014\)](#);

¹A simulation with no ions, but otherwise identical to 1d×1kCoo102, demonstrates similar overall dynamics of the reconnection process.

²This situation applies to the current sheets in pulsar magnetospheres and can be relevant to the pair-dominated coronae of accreting black holes.

Hakobyan et al. (2019, 2021), which provides a fresh supply of magnetized upstream plasma to replenish the outflowing population. In the third direction, $\pm\hat{y}$, we impose periodic boundary conditions. The aspect ratio of the domain in the plane of the current sheet is $L_x/L_y = 1$, and the size perpendicular to the current sheet is $L_z \sim 0.3...0.5L_x$, large enough to separate the current sheet and the plasma injector.

To trigger the reconnection, we remove thermal pressure in the central part of the current layer as done in Sironi et al. (2016). This procedure launches transients leading to the onset of reconnection and moving outward from the center at approximately the speed of light. To minimize the effects of boundary conditions on the statistical steady state, the radiative cooling and the outflow boundaries are switched on when the transient approaches the edge of the simulation box, at $0.5 L_x/c$, where L_x is the extent of the domain in the x -direction. In all of our analyses, we exclude the particles from the initial current layer, essentially ignoring the evolution of the sheet in its early transient stage. We consider the current sheet to enter the steady-state regime after about $t \gtrsim 1.5 L_x/c$ from the beginning of the simulation. We run all simulations for 3 additional light-crossing times, with a total duration of $4.5 L_x/c$. All statistical data is averaged over the final 3 light-crossing times, while the current sheet is in statistical steady-state.

To study the numerical convergence of our results, we explore resolutions of $d_e/\Delta x = 1...3$, where $d_e = (m_e c^2 / 4\pi n_{\text{up}} e^2)^{1/2}$ is the upstream plasma skin depth, and Δx is the cell size ($\Delta x = \Delta y = \Delta z$); the comparison of these two resolutions is presented in Appendix A.2, where we find only marginal differences in the reconnection dynamics and particle acceleration. The choice of σ and d_e determines the characteristic resolution of the Larmor orbits of particles, $r_L^{\text{up}}(\gamma) = \gamma m_e c^2 / (|e| B_{\text{up}})$, accelerated up to $\gamma \sim \sigma$: $r_L^{\text{up}}(\sigma)/\Delta x = (d_e/\Delta x)\sqrt{\sigma} \approx 7...21$, assuming particles gyrate in the perpendicular magnetic field equal to the upstream value. To

better characterize the acceleration mechanisms and minimize the boundary effects, we set our fiducial box sizes to be in the range of $L_x/r_L^{\text{up}}(\sigma) \approx 150\ldots 300$. Our timestep duration (the CFL number) is fixed at $c\Delta t/\Delta x = 0.225$ for most of our simulations. The summary of all the simulation parameters is given in Table 2.1; a discussion on how the synchrotron radiation is implemented (last column) follows further (see Sec. 2.3.2). In our strongest cooling simulation, 1dx1kCoo101, we employ a halved CFL number, $c\Delta t/\Delta x = 0.1125$, compared to our default choice of 0.225, in order to properly resolve the plasma oscillation period in regions of dense flux tubes, highly compressed due to strong radiative cooling (see Sec. 2.4). The upstream plasma is presented with a total concentration of both electrons and positrons averaging to 1 particle per cell³. To test the convergence, we run a simulation identical to 1dx1kCoo102 but with 8 particles per cell in total and obtain identical results. To mitigate the numerical artifacts from the finite number of particles per cell we employ 8 digital filter passes on the deposited currents with a $(1/4, 1/2, 1/4)$ stencil.

To analyze the properties of the current sheet in steady-state, we need to employ a robust method to distinguish the current layer carrying the plasma energized in reconnection from the fresh plasma upstream. The most convenient way is to separately trace particles from the two upstreams ($z \gtrsim 0$ and $z \lesssim 0$ regions), and define the mixing factor, $\mathcal{M}$, as suggested by Zhang et al. (2021a),

$$\mathcal{M} = 1 - 2 \left| \frac{n_{\uparrow}}{n} - \frac{1}{2} \right|. \quad (2.2)$$

Here $n_{\uparrow}$ is the local number density of particles originating from the $z > 0$ region, and n is the total number density. This definition implies $\mathcal{M} = 0$ in either of the two upstream regions and

³I.e., each cell either contains an electron and a positron injected at the same location to satisfy zero numerical charge or does not contain any simulation particles.

$\mathcal{M} \approx 1$ in the mixed region of the current layer. We use the threshold $\mathcal{M} = 0.7$ to separate the upstream from the reconnecting current sheet, however, our results are insensitive to this particular choice, as the upstream-to-downstream transition is quite sharp in terms of the values of $\mathcal{M}$.

Table 2.1: Summary of simulation parameters.

Simulation	$L_x \times L_y \times L_z$	$d_e/\Delta x$	$c\Delta t/\Delta x$	$\gamma_{\text{syn}}/\sigma$
3dx2kCool02	$2000^2 \times 1000$	3	0.225	0.2
3dx2kCool5	$2000^2 \times 1000$	3	0.225	5.0
3dx3kCool02	$3000^2 \times 1000$	3	0.225	0.2
3dx3kCool05	$3000^2 \times 1000$	3	0.225	0.5
3dx3kCool11	$3000^2 \times 1000$	3	0.225	1.0
3dx3kCool15	$3000^2 \times 1600$	3	0.225	5.0
3dx3kCoolInf	$3000^2 \times 1600$	3	0.225	∞
1dx1kCool02	$1000^2 \times 1000$	1	0.225	0.2
1dx2kCool02	$2000^2 \times 1000$	1	0.225	0.2
1dx1kCool01	$1000^2 \times 1000$	1	0.1125	0.1

2.3.2 Radiation

To model the synchrotron cooling, the code imposes a synchro-curvature radiation reaction force in the [Landau & Lifshitz \(1975\)](#) form, acting exclusively on the electrons and positrons (ions are unaffected by cooling):

$$\begin{aligned}
\mathbf{F}_{\text{rad}} &= \frac{\sigma_T}{4\pi} (\boldsymbol{\kappa}_R - \gamma^2 \tilde{B}_\perp^2 \boldsymbol{\beta}), \\
\boldsymbol{\kappa}_R &= (\mathbf{E} + \boldsymbol{\beta} \times \mathbf{B}) \times \mathbf{B} + (\boldsymbol{\beta} \cdot \mathbf{E}) \mathbf{E}, \\
\tilde{B}_\perp^2 &= (\mathbf{E} + \boldsymbol{\beta} \times \mathbf{B})^2 - (\boldsymbol{\beta} \cdot \mathbf{E})^2.
\end{aligned} \tag{2.3}$$

Here, $\beta = v/c$ is normalized 3-velocity, and $\gamma = (1 - \beta^2)^{-1/2}$ is the Lorentz factor⁴. In magnetospheres of compact objects, the radiative drag becomes dynamically important at very high Lorentz factors, $\gamma \gg 1$, depending on the strength of the magnetic field at the location of the reconnecting current sheet. In order to model a similar physical regime in simulations, we rescale the Thomson cross-section σ_T , using a dimensionless parameter, γ_{syn} (Uzdensky, 2018; Werner et al., 2019; Hakobyan et al., 2019; Sironi & Beloborodov, 2020), defined as follows:

$$|e|E_{\text{rec}} \equiv \frac{1}{4\pi} \sigma_T B_{\text{up}}^2 \gamma_{\text{syn}}^2. \quad (2.4)$$

Thus, the value of γ_{syn} corresponds to the Lorentz factor of particles for which the synchrotron drag force in the perpendicular background field of B_{up} is equal to the acceleration force $|e|E_{\text{rec}}$ (where for magnetic reconnection, we can typically estimate the reconnection-driven electric field E_{rec} as $\approx 0.1 B_{\text{up}}$).

The relative importance of radiative cooling for particle acceleration is given by the ratio of γ_{syn} and the upstream magnetization parameter, σ . Since the characteristic energy gain in X-points of the reconnecting sheet corresponds to $\gamma \gtrsim \sigma$ (Sironi & Spitkovsky, 2014; Werner et al., 2016), we call the cooling regime “strong” if the particles are strongly affected by cooling before their Lorentz factors reach σ around the X-point, i.e., if $\gamma_{\text{syn}} \lesssim \sigma$. The limit of $\gamma_{\text{syn}} > \sigma$ is referred to as the “weak” cooling regime.

The second quantity that needs re-scaling is the energy of radiated synchrotron photons, or, equivalently, Planck’s constant, $\hbar$. It is easy to see, that the particles with $\gamma \approx \gamma_{\text{syn}}$, with γ_{syn} being defined as in (2.4), moving in the perpendicular magnetic field of strength B_{up} radiate

⁴The total radiated power, $P_{\text{sync}} \propto \gamma^2 \tilde{B}_{\perp}^2$, corresponds to a square of the electric field in the rest frame of the moving particle, $\mathbf{E}' = \gamma(\mathbf{E} + \boldsymbol{\beta} \times \mathbf{B}) - (\gamma - 1)(\boldsymbol{\beta} \cdot \mathbf{E})\boldsymbol{\beta}/\beta^2$.

photons with the characteristic energy

$$\varepsilon_c = \frac{3}{2} \hbar \frac{|e| B_{\text{up}}}{m_e c} \gamma_{\text{syn}}^2 = \frac{9}{4 \alpha_F} \frac{E_{\text{rec}}}{B_{\text{up}}} m_e c^2 \approx 16 \text{ MeV}, \quad (2.5)$$

where $\alpha_F \approx 1/137$ is the fine-structure constant (see also [Uzdensky et al., 2011](#); [Cerutti et al., 2012a](#); [Uzdensky & Loureiro, 2016](#)). Importantly, this value is insensitive to the strength of the magnetic field and is thus a perfect benchmark to calibrate the photon energies. Thus, in our simulations, a particle with Lorentz factor of γ experiencing a magnetic field of $\tilde{B}_\perp$, defined in equation (2.3), by definition radiates a spectrum of photons which peaks near

$$\varepsilon_p = 16 \text{ MeV} \frac{\tilde{B}_\perp}{B_{\text{up}}} \left(\frac{\gamma}{\gamma_{\text{syn}}} \right)^2. \quad (2.6)$$

2.4 Structure of the Current Sheet

The dynamics of a current sheet, undergoing a 3D magnetic reconnection, is generally similar to that in 2D, but there are also significant differences specific to 3D that must be considered. The relativistic drift-kink instability is the fastest growing in 3D, and dominates the initial evolution. As the current sheet evolves, the plasmoid instability in the nonlinear phase becomes dominant (e.g., [Sironi & Spitkovsky, 2014](#)). The dynamical state of the current sheet at late stages is thus dominated by a number of large flux tubes (plasmoids), separated by kinetic-scale X-points. Similar to 2D, the plasmoids are continuously forming and undergoing merging processes. However, contrary to 2D, in 3D they are also being continuously “dissipated” by the MHD-kink instability ([Werner & Uzdensky, 2021](#)).

To visualize the complex 3D structure, Fig. 2.1 shows volume renderings of the plasma density in the fully-developed non-linear reconnection stage of 3D current sheets from two different simulations. Fig. 2.1a shows a density snapshot of the simulation with weak cooling ($\gamma_{\text{syn}}/\sigma = 5$; $3\text{d}\times 3\text{kC}\times 15$), while Fig. 2.1b corresponds to the strong cooling case ($\gamma_{\text{syn}}/\sigma = 0.2$; $3\text{d}\times 3\text{kC}\times 102$). The small ripples, visible in the perpendicular-to-field direction (y - z plane), are caused by the relativistic drift-kink instability. The imprint of the ideal kink instability is visible as the overall twist and deformation of the large flux tubes in both panels of Fig. 2.1. The lines in the figure represent the magnetic field lines and highlight the main features of the current sheet, including the X-points where field lines reconnect crossing the $z = 0$ plane, and the plasmoids where magnetic field lines are helically wrapped around high-density structures. The color of the field lines shows the strength of the magnetic field $|B|$, particularly demonstrating that it falls to zero in the X-points.

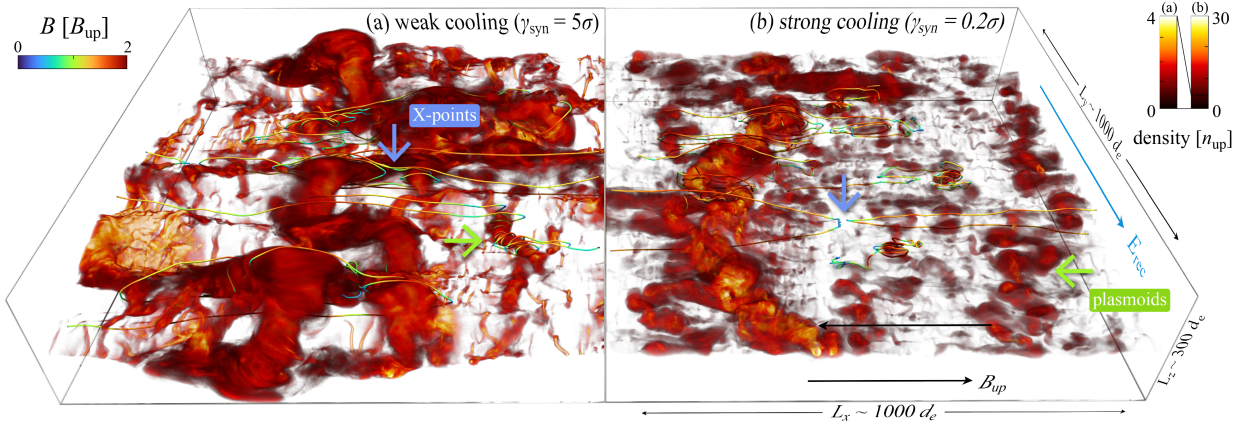


Figure 2.1: Volume rendering of the plasma density in the current sheet with weak (a), and strong (b) cooling at $t \sim 2.2L_x/c$. Magnetic field lines are also displayed, with the color representing the strength, $|B|$. From the field lines, one can easily identify the non-ideal regions – the X-points (selected ones are shown by blue arrows)– where the magnetic field vanishes (this corresponds to regions where the field lines are blue). The sizes of plasmoids (selected ones shown by green arrows) are significantly different in these two simulations, with the strong cooling run displaying smaller structures and larger compression (note that the colorbars for the two plots have different scales).

To systematically quantify the variations of the current sheet structure due to synchrotron cooling, we evaluate the current sheet width, w (along the z axis), for individual points in the horizontal (x - y) plane using the definition of the mixed region in Sec. 2.3.1 (this includes both the thin current sheet and plasmoids of various sizes). Fig. 2.2 shows the distribution of widths, $\mathcal{P}(w)$, for a large collection of points for four simulations with varying cooling strength (in the order of increasing cooling strength: 3dx3kCool1Inf, 3dx3kCool15, 3dx3kCool11, 3dx3kCool102). There are two important conclusions one can draw from this plot. First of all, the peak of the $\mathcal{P}(w)$ is slightly smaller than $w_{\text{peak}} \sim r_L^{\text{up}}(\sigma) = d_e \sqrt{\sigma}$, which corresponds to the Larmor radii of particles with the Lorentz factor close to σ in the upstream magnetic field, B_{up} . The presence of

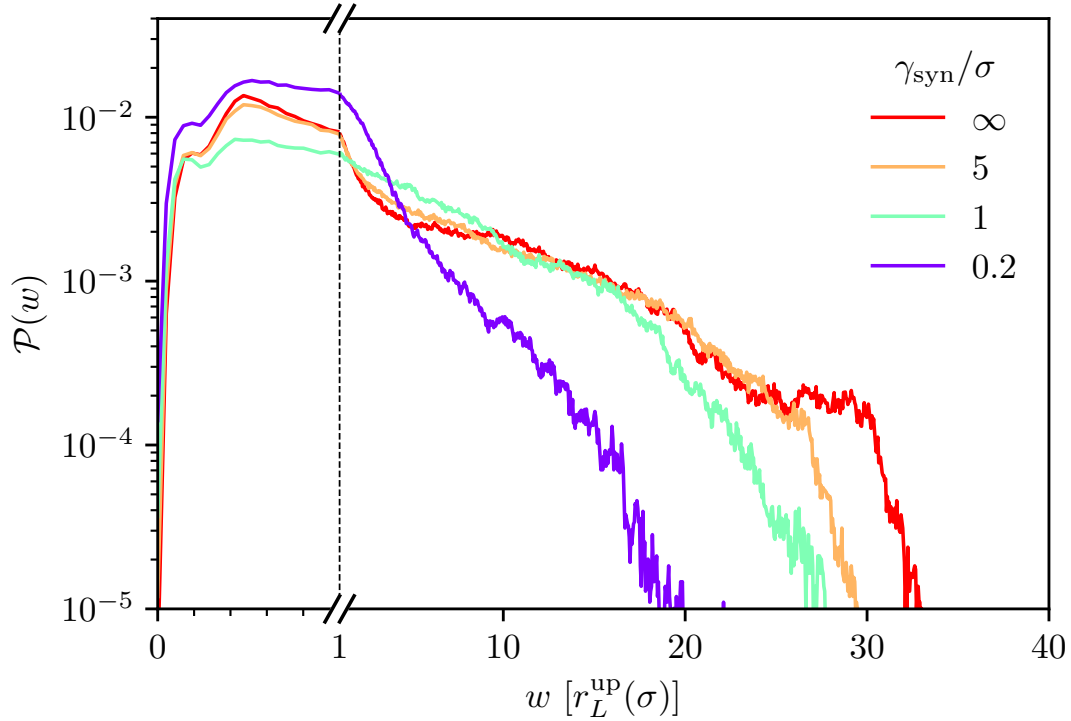


Figure 2.2: Distribution of transverse sizes (along z) of the current layer for different cooling regimes. The scale is zoomed in for $w < r_L^{\text{up}}(\sigma)$. Stronger cooled current sheets form on average smaller structures due to significant compression. The peak of the distribution, i.e., most of the current sheet, has a thickness below the Larmor radius of particles accelerated up to σ : $w \sim r_L^{\text{up}}(\sigma)$.

this peak is associated with the X-points, where particles are accelerated to characteristic Lorentz factors of $\gamma \lesssim \sigma$ (see Sec. 2.5.1). More importantly, current sheets with stronger cooling display a steeper slope and shorter distribution span for $\mathcal{P}(w)$, indicating that the plasmoids in the current layer become smaller and more concentrated as cooling strength increases (more detailed analysis of this issue was recently performed by Schoeffler et al. 2023).

Variation of the current sheet structure due to synchrotron cooling is reflected in one of the most important macroscopic parameters of the reconnection – its rate, β_{rec} . The evolution of the rate, measured as the inflow velocity of the plasma, $\beta_{\text{rec}} = (\mathbf{E} \times \mathbf{B})_z / |\mathbf{B}|^2$, averaged in the thin layer far upstream parallel to the current sheet, for different cooling regimes is shown in Fig. 2.3. As mentioned before, we focus our analysis in the time range $t \gtrsim 1.5 L_x/c$, when the simulation reaches a dynamic steady state. We find that the reconnection rate in the weak cooling case ($\gamma_{\text{syn}} \gtrsim \sigma$) is $\beta_{\text{rec}} \approx 0.06 \dots 0.1$ (cf. Zhang et al. 2021a). Notably, in the simulations with stronger cooling ($\gamma_{\text{syn}} \lesssim \sigma$), the reconnection rate is systematically larger, with values ranging around $\beta_{\text{rec}} \approx 0.12 \dots 0.14$. We associate this correlation with the change in the filling fraction of X-points in the current layer, which for the stronger cooled runs is approximately twice as large, compared to the run without cooling as seen from the peak values in Fig. 2.2. This conclusion is also consistent with the characteristically smaller plasmoid sizes in the strong cooling regime, as discussed above.

Plasmoids accumulate the hot plasma energized in the X-points. In 2D the energetic particles are well magnetized and confined, resulting in plasmoids being relatively coherent and clearly distinguishable from other relativistic outflows present in the sheet. In 3D the structure of plasmoids is richer and less coherent. Nonetheless, in both cases, their equilibrium is governed by an adiabatic balance between the gradient of thermal pressure of the hot plasma,

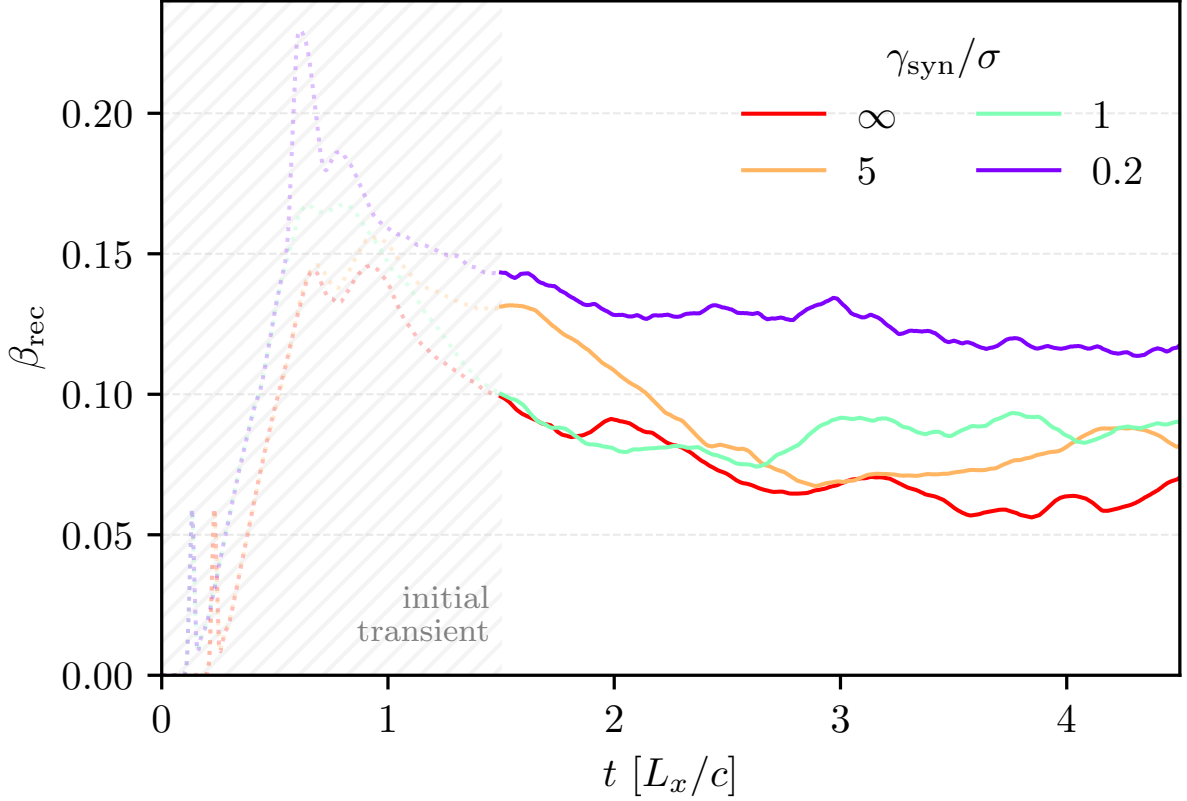


Figure 2.3: Time evolution of the reconnection rate in simulations with different cooling regimes. Measured rates are of order $\beta_{\text{rec}} \approx 0.06 \dots 0.14$, and are systematically larger for the case of strongest cooling, $\gamma_{\text{syn}}/\sigma = 0.2$.

P , and the Lorentz force: $\mathbf{j} \times \mathbf{B} = c \nabla P$, with $\mathbf{j}$ being the total current density. In general, plasma pressure is anisotropic and needs to be computed self-consistently as a moment of the particle distribution function. To compute it, we separate the bulk component of the particle momentum from the thermal one by defining the fluid (Eckart) frame for species s , moving with a velocity $U_s^\mu = N_s^\mu / \sqrt{N_{\nu,s} N_s^\nu}$, where $N_s^\mu \equiv \int f_s(u^\mu/u^0) d^3u$, and f_s and u^μ are the distribution and the four-velocity of particle species s . The plasma pressure tensor for species s , $P_s^{\mu\nu}$, is then computed as the projection of the stress-energy tensor: $P_s^{\mu\nu} = \Delta_\alpha^\mu \Delta_\beta^\nu T_s^{\alpha\beta}$, where $T_s^{\mu\nu} \equiv \int (p^\mu p^\nu / p^0) f_s d^3p$, $\Delta_s^{\mu\nu} \equiv \eta^{\mu\nu} + U_s^\mu U_s^\nu$ is the projection operator, and $\eta^{\mu\nu}$ is the Minkowski tensor ($p^\mu = m_s u^\mu$ is the four-momentum of the particle of species s). We find that the non-

diagonal elements of the pressure tensor compared to diagonal elements are typically at the level of 5% inside plasmoids in the weak-to-no cooling regime, with an increase to non-negligible $\approx 20\%$ for the strongest cooling regime, $\gamma_{\text{syn}}/\sigma = 0.2$. In addition, inside plasmoids, the diagonal component perpendicular to the magnetic field is suppressed by a factor of ~ 5 in the case of strong cooling.⁵ Guided by these findings, we consider the gradient of the full pressure tensor to analyze the force balance of flux tubes.

In typical 2D simulations (with or without cooling), the force balance inside plasmoids is achieved by strong compression evident in stronger magnetic fields and higher densities (Sironi et al., 2016). The presence of strong cooling exacerbates this effect⁶, by removing the pressure support, which results in extremely compressed structures (see, e.g., Hakobyan et al., 2019). In 3D, the picture is quantitatively different. The force-balance condition is demonstrated in Fig. 2.4, where in panels (a) and (c) we show the cross-section of the plasma density in x - z plane for two simulations: one without cooling (3dx3kCoolInf), and the other with strong cooling (3dx3kCool02). We pick intermediate-size plasmoids in both of these simulations (indicated by the dotted vertical lines in left panels) and in panels (b) and (d) plot the magnetic pressure, $B^2/8\pi$, the scalar plasma pressure, defined as $P = 1/3 \sum_s P_{s,i}^i$, as well as the z -components of the two forces: the pressure gradient, $c\nabla_i P^{iz}$, and the Lorentz force, $(\mathbf{j} \times \mathbf{B})_z$, as functions of z . In the no-cooling simulation (Fig. 2.4a, b), large plasmoids are typically weakly magnetized, while in the strong cooling regime (Fig. 2.4c, d), the strength of the magnetic field inside the plasmoids is several times larger than the upstream value, B_{up} , (compare blue solid lines in Fig. 2.4b, d). Despite this, even in the strongest cooling regime, the magnetic field

⁵The role of the pressure anisotropy due to synchrotron cooling is discussed by Zhdankin et al. (2022).

⁶This effect is investigated in detail by Schoeffer et al. (2019), who observed that this compression provides a positive feedback loop: magnetic field amplification by compression inside plasmoids leads to stronger synchrotron cooling, which in turn promotes further compression.

compression inside the plasmoids is much smaller than what is observed in 2D (Hakobyan et al., 2021; Schoeffler et al., 2023). Notably, the strongest B -field component inside the plasmoid (in the strong cooling regime) is the out-of-plane one, B_y . The generation of this out-of-plane magnetic field component, as well as the dissipation of the in-plane B_{xz} components (evident in the no-cooling regime), are both consequences of the MHD kink instability of 3D flux tubes. This instability dissipates the x - z component of the magnetic field, converting it to the plasma thermal energy (e.g., Werner & Uzdensky, 2021), while simultaneously acting as a dynamo by generating an out-of-plane field component, B_y , until the flux tube becomes stable to kinking (Werner & Uzdensky, 2017; Alves et al., 2018; Davelaar et al., 2020; Zhang et al., 2021b). Importantly, this interplay is inherent specifically to 3D reconnection and is thus completely overlooked in 2D.

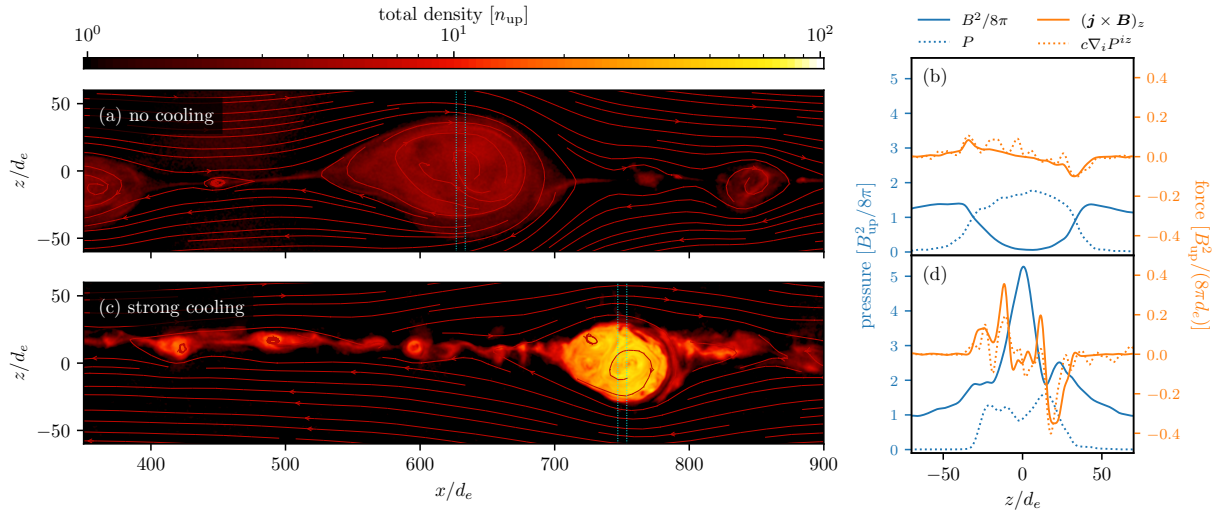


Figure 2.4: The cross-section of typical plasmoids for uncooled (a, b) and strongly cooled (c, d) current sheets. The left column (a, c) shows density slices and the magnetic field lines. The region highlighted with a cyan-dotted boundary is shown in 1D cuts in (b, d), where we plot the magnetic field, $B^2/8\pi$, and gas pressure, P , with the corresponding magnetic tension force, $(\mathbf{j} \times \mathbf{B})_z$, and the gradient of the pressure tensor, $c\nabla_i P^{iz}$. In the uncooled case, the plasmoids are in equilibrium in terms of magnetic and thermal pressure, with the thermal pressure typically dominating inside the plasmoids. On the other hand, the equilibrium inside the cooled plasmoids is more dynamic, with the magnetic field pressure being larger than the upstream value, and dominating the plasmoid cores.

The plasma inside flux tubes both in uncooled and strongly cooled simulations is compressed and heated to a rough equipartition with the upstream magnetic field pressure: $P \approx B_{\text{up}}^2/8\pi$. In the latter case, the pressure is mainly delivered by higher plasma density (compare densities in panels (a) and (c) in Fig. 2.4), since strong synchrotron cooling prevents plasmoids from heating up by effectively radiating away the components of particle momenta perpendicular to the magnetic field. Orange lines in Fig. 2.4b and d show the two force terms: the gradient of the plasma pressure tensor, and the magnetic force, both plotted against the vertical axis z . The balance between these two forces is well satisfied for the flux tube in the uncooled simulation (Fig. 2.4b), and is reasonable in the strongly cooled simulation (Fig. 2.4d).

In this section, we largely neglected the presence of ions (insusceptible to synchrotron cooling) in plasmoids. While their contribution to the total pressure was self-consistently included in our calculations, their relative contribution was typically small for the timescales considered, e.g., in Fig. 2.4, due to their smaller number density. However, at later times ions can accelerate to much higher Lorentz factors than the pairs (see Sec. 2.5.1), at which point their lack in the number density is compensated by the higher energy. When the ion pressure dominates in flux tubes, we expect the structures of plasmoids to be similar to the case of no-cooling (Fig. 2.4a, b and Appendix A.1). The applicability of this regime for the current sheets in pulsar magnetospheres is studied in detail in Sec. 2.7.1.

2.5 Particle acceleration

During the non-linear stage of reconnection, cold particles from upstream – both pairs and ions – are advected into the current sheet, where they are energized in either the X-points or the

relativistic outflows from the X-points (e.g., [French et al., 2022](#)). After this initial “injection” stage most of these particles enter into the plasmoids. In 2D, strong compression of the magnetic field inside plasmoids leads to adiabatically slow secondary acceleration of these particles to high energies limited by the system size ([Petropoulou & Sironi, 2018](#); [Hakobyan et al., 2019](#)). However, in 3D, the finite size of plasmoids in the out-of-plane direction allows the most energetic particles to freely escape from their bounds. Moreover, as described above, compression of the magnetic field inside plasmoids is significantly smaller in 3D. Thus, the same secondary acceleration mechanism as in 2D cannot operate efficiently in 3D. On the other hand, in 3D, another acceleration channel has been observed for uncooled pair plasmas by [Zhang et al. \(2021a, 2023\)](#), with the large-scale reconnection electric field in the upstream, E_{rec} , being its main driver.

The dynamics of plasma in X-points is insensitive to the presence of synchrotron cooling, as the magnetic field vanishes in these regions. However, as soon as the energized pairs leave the X-points, they lose their energy shortly after being exposed to a strong perpendicular magnetic field component. As a result, the secondary acceleration channel in the reconnection upstream is strongly suppressed for pairs. Ions, on the other hand, are unaffected by cooling, and can thus tap the upstream electric field and get accelerated to potentially large Lorentz factors, $\gg \sigma$. This can clearly be seen in [Fig. 2.5](#), where we plot the energy spectra of both pairs and ions. For the simulations with strong-to-marginal cooling ($\gamma_{\text{syn}}/\sigma = 0.2\dots 1$), the distribution of pairs (solid lines) typically spans up to $\sim \sigma$, with the slope being close to γ^{-2} at highest energies. Ions (dashed lines), on the other hand, regardless of the cooling strength, extend to Lorentz factors significantly above σ , with the cutoff energy itself increasing in time. In the uncooled simulation (red lines), the cutoff in the distribution of pairs is exactly the same as for the ions, suggesting that the underlying secondary acceleration mechanisms are identical. Surprisingly, in simulations

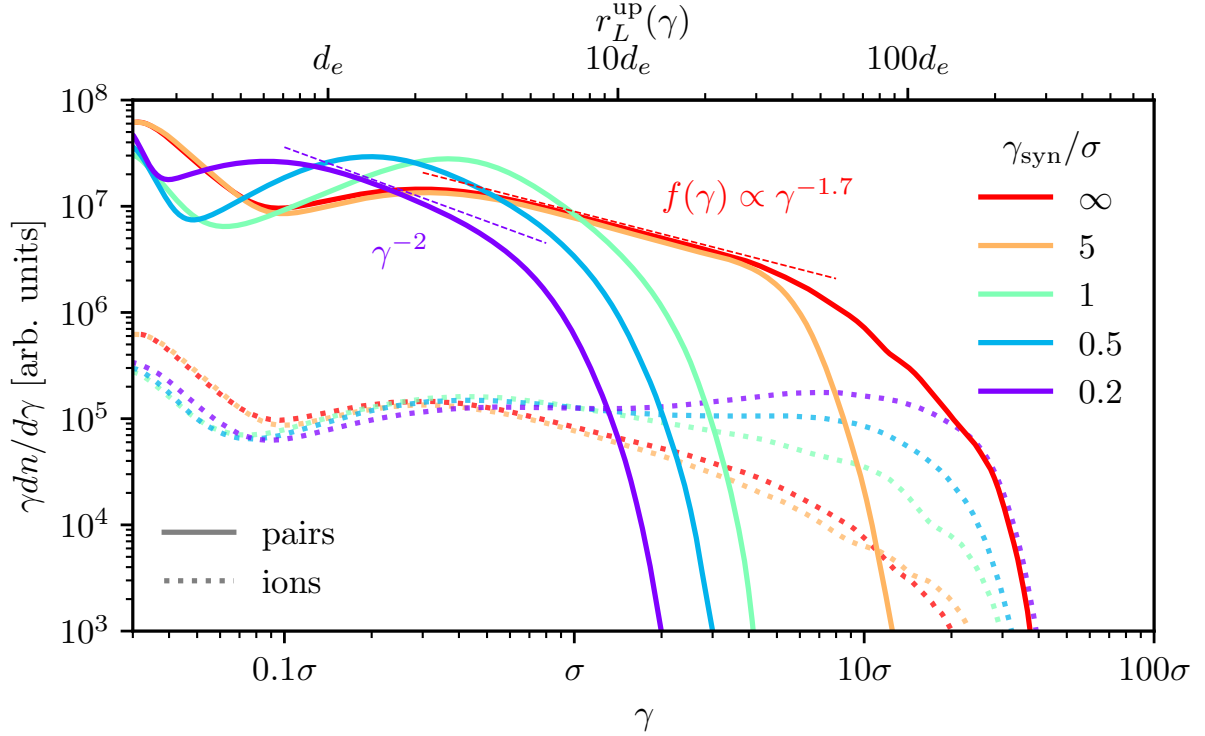


Figure 2.5: Energy distributions of the pairs (solid lines) and the ions (dotted lines) at $t = 2.25 L_x/c$ for different cooling regimes: purple lines correspond to strong cooling $\gamma_{\text{syn}}/\sigma = 0.2$, and the red ones for the uncooled runs, $\gamma_{\text{syn}}/\sigma = \infty$. Dynamically strong cooling effectively limits the maximum energy of pairs to $\sim \sigma$, but at the same time facilitates the formation of harder distribution for the ions. For reference, we also show the corresponding Larmor radii, r_L^{up} .

with stronger cooling, ions form a much harder power law distribution, with the typical slope close to γ^{-1} for the strongest cooling case, and steepening to $\gamma^{-1.7}$ for the uncooled case (this effect is explored in Sec 2.5.1).

We further quantify the contributions of different acceleration channels in Fig. 2.6, where we plot the energy spectra of both ions and pairs in the uncooled and strongly cooled simulations. To test the dominant acceleration mechanism up to $\gamma \sim \sigma$ (the injection stage), we choose the particles that experienced the non-ideal electric field, $E > B$, during their lifetime (dashed lines in Fig. 2.6). In the inset panels of Fig. 2.6 we also plot the distribution of energies, $\Delta\gamma$, gained during the passage of $E > B$ regions. Notably, the particles that have passed through an $E > B$

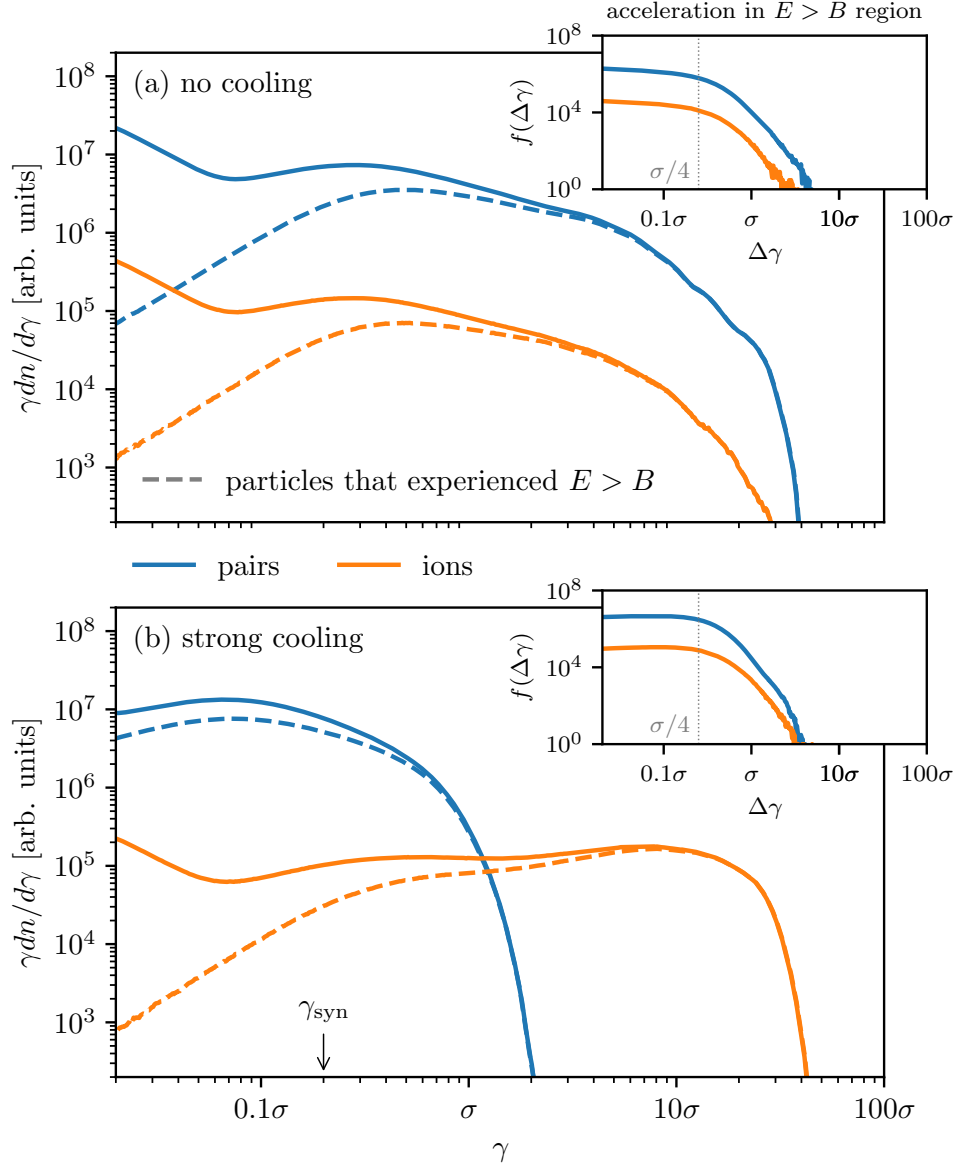


Figure 2.6: Distribution functions of the pairs (blue) and the ions (orange) for both the uncooled (a) and the strongly cooled (b) simulations. Dashed lines show parts of the distributions contributed by the particles having experienced $E > B$ during their lifetime in the sheet. Inset panels show the distribution of energy gains, $\Delta\gamma$, of particles during their contiguous presence in these regions. The non-ideal electric field in the X-points only accelerates particles to Lorentz factors of a fraction of σ . At the same time, this process is essential to promote the particles (pairs and ions) for further acceleration up to the highest energies: in relativistic outflows along the layer and the large-scale reconnection electric field upstream. As a result, the high energy tail of the spectrum is fully formed by the particles passed through $E > B$.

region dominate at energies $\gtrsim \sigma$ (cf. [Sironi, 2022](#)). However, the net energy gain of particles in X-points in both the uncooled and strongly cooled cases does not exceed a fraction of σ , and thus cannot account for the high-energy power-law tail of either the pairs or the ions (e.g., [Werner et al., 2016](#); [Uzdensky, 2022](#)). The main energy gain, thus, occurs due to the acceleration by ideal electric fields. For pairs, as we argue further in Sec. 2.5.2, “pick-up” acceleration by outflows from the X-points allows them to gain energies up to σ even in the strong cooling case.

2.5.1 Ion acceleration

In our simulations, ions are modeled as separate particle species with a mass equal (or comparable) to that of pairs. Additionally, ions do not undergo synchrotron cooling. In Fig. 2.7, we present 2D projections of the ion trajectories obtained from the 3D simulation with strong synchrotron cooling, $\gamma_{\text{rad}} = 0.2\sigma$, 3dx3kCool02. Each point on the trajectory is color-coded according to the Lorentz factor of the ion, this information is also displayed on the inset axes. The thickness of the line shows whether the ion is upstream (thick) or inside the current sheet (thin). We initiate the tracking by picking two representative ions with energies $\gamma \sim \sigma$ at $t = 2.25 L_x/c$ (corresponding to the same moment shown in Fig. 2.5, when the spectrum shape is in the steady-state), and follow their evolution until the end of the simulation $t = 4.5 L_x/c$. Interestingly, the behavior of the two ions differs significantly. One ion reaches a high Lorentz factor, $\sim 30\sigma$, while the energy of the other ion oscillates around a few σ . Notably, the ion that gains large energy experiences most of its acceleration while moving outside of the current layer ([Zhang et al., 2021a, 2023](#)), as evident from the thickness of the line representing the trajectory in Fig. 2.7. These trajectories correspond to bouncing of accelerated ions between the two converging upstream

flows⁷. In contrast, the low-energy ion displays a tangled trajectory due to its entrapment inside a plasmoid or frequent scattering on magnetic field inhomogeneities inside the current sheet. This behavior is likely associated with self-generated turbulence (e.g., [Guo et al., 2021](#)).

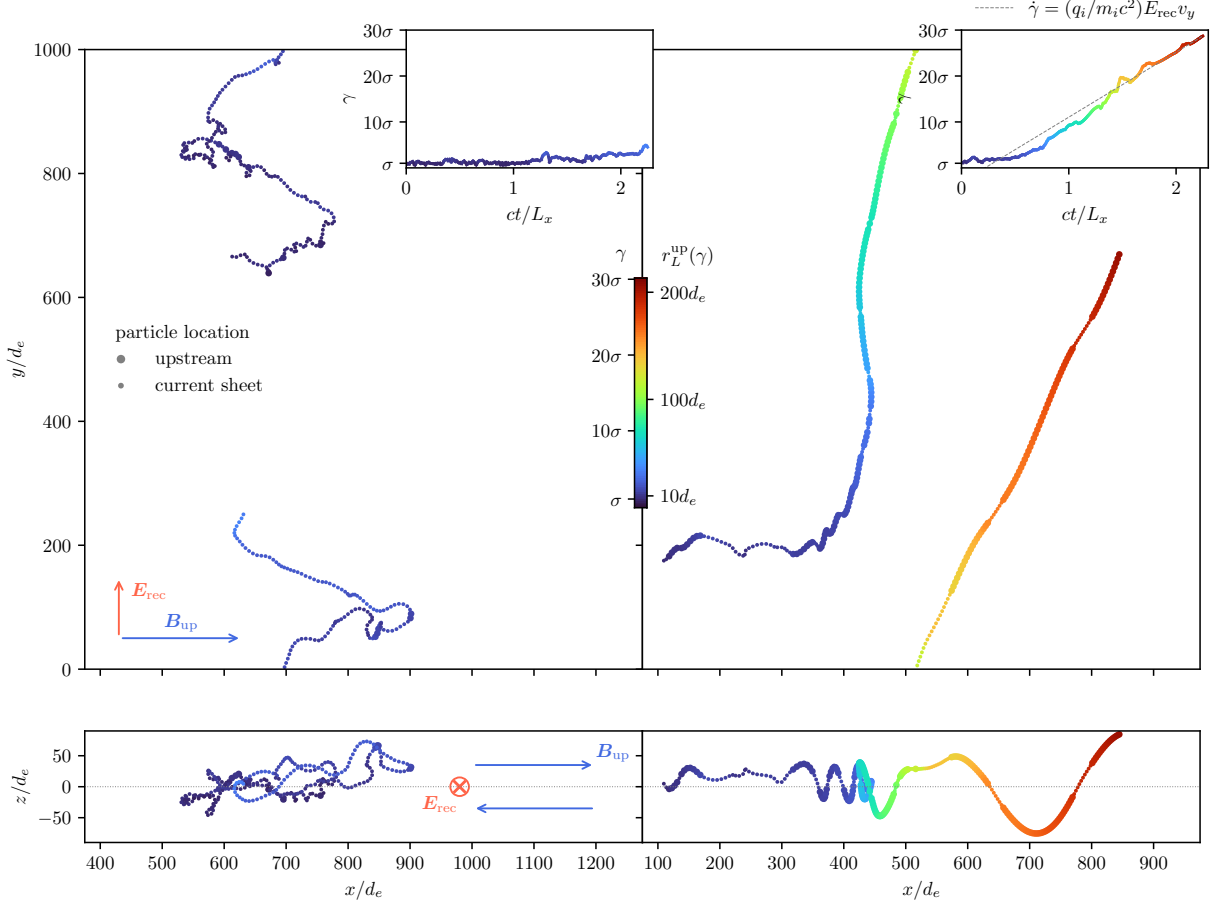


Figure 2.7: Archetypal trajectories (projected to the x - y and the x - z plane) and the energy evolution of two ions. The color of each dot corresponds to the Lorentz factor of the ion at that particular time, while the size of the dot indicates whether the ion is inside the current sheet (small dots), or upstream (large dots). The ion on the right accelerates to high energies ($\gg \sigma$) by the large-scale electric field, E_{rec} , in the region upstream of the layer, while the one on the left gains only a marginal amount of energy ($\sim \sigma$). The more energetic ion on the right primarily moves upstream of the current layer along the reconnection electric field (y direction) and accelerates linearly in time. The slower particle, on the other hand, is trapped inside the current sheet and experiences constant scatterings on local inhomogeneities.

The acceleration of the more energetic ion takes place linearly in time, with its trajectory

⁷Note that the classical Speiser orbits correspond to the acceleration inside the current sheet, as opposed to the acceleration upstream (for more discussion, see [Zhang et al. 2021a](#)).

showing almost no scatterings. Notably, the preferred direction of motion for this ion is in y , while the acceleration rate is almost constant. Both of these facts point towards the large-scale reconnection-driven electric field, $\mathbf{E}_{\text{rec}} \approx \mathbf{E}_y$, being the dominant acceleration mechanism. The rate of acceleration for these trajectories can be written as

$$\dot{\gamma} = \frac{q_i}{m_i c} \left(\mathbf{E}_{\text{rec}} \cdot \frac{\mathbf{v}}{c} \right) \implies \Delta\gamma \approx \beta_{\text{rec}} \beta_y \frac{q_i B_{\text{up}}}{m_i c} \Delta t, \quad (2.7)$$

where we used $E_{\text{rec}} = \beta_{\text{rec}} B_{\text{up}}$ (Zhang et al., 2021a). This slope is overplotted with a dashed line in the inset of the left panel of Fig. 2.7, showing excellent agreement with the acceleration history. Since ions are by construction not susceptible to synchrotron cooling, we find that this acceleration mechanism operates similarly for both the strong cooling and no-cooling regimes.

Despite the ion acceleration mechanism being identical for different cooling regimes, the distribution of ions is much harder in the stronger cooling cases, as highlighted with dashed lines in Fig. 2.5. Two parameters from Eq. (2.7) that could potentially contribute to different acceleration rates are the reconnection rate, β_{rec} ,⁸ and the beaming of ions towards the direction of the reconnection electric field, β_y . To inspect the relative importance of the two effects, in Fig. 2.8 we plot time-dependent statistical distributions of particle Lorentz factors, γ (panels a1...a4), their velocity component along the reconnection electric field in y , β_y (panels b1...b4), and the electric field strength in y (panels c1...c4). The first two columns represent ions from the uncooled simulation, while the last two – from the simulation with the strongest cooling, $\gamma_{\text{syn}}/\sigma = 0.2$. Since we are primarily interested in the post-injection phase, in the odd-numbered columns we select ions that have reached energies $\gamma_{\text{in}} \approx \sigma$ (the core of the distribution), while in

⁸As it was already mentioned in Sec. 2.4, the global reconnection rate is larger for the strongest cooling regime, which in turn means that the large-scale electric field, E_y , is typically stronger.

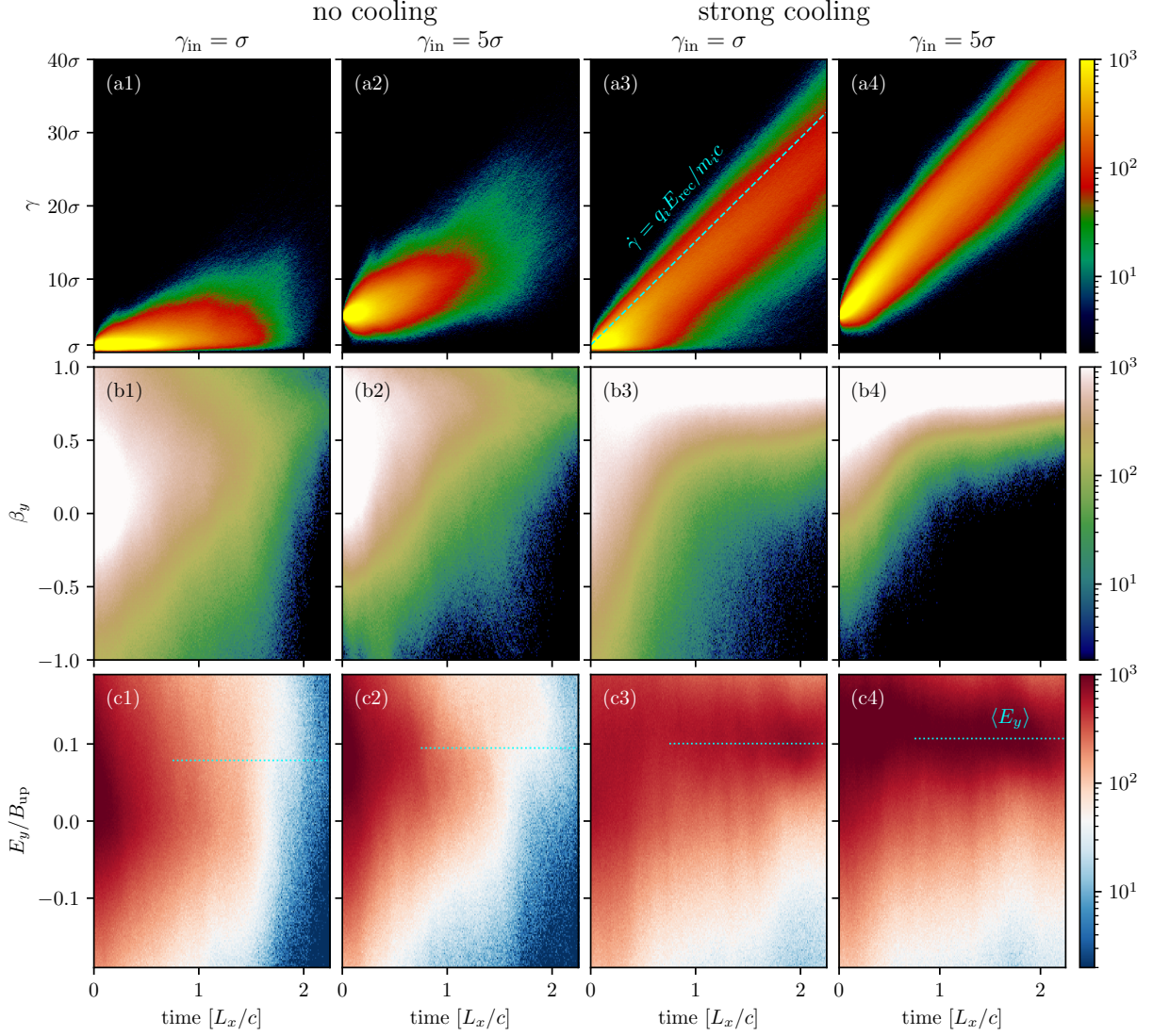


Figure 2.8: Time evolution of a sample of ions that reach energies, correspondingly, $\gamma_{\text{in}} = \sigma$ (odd columns), and $\gamma_{\text{in}} = 5\sigma$ (even columns). Results from two simulations are shown: with no cooling (first two columns), and strong cooling current (last two columns). The upper row (a1...a4) shows the evolution of the Lorentz factor of the ions. The middle row (b1...b4) shows the evolution of the y component of their 3-velocities (along the reconnecting electric field E_{rec}). The lower row (c1...c4) shows the value of the reconnection electric field E_y experienced by these ions. In the strong cooling regime, most of the ions successfully accelerate in the reconnection electric field and experience little-to-no scatterings (evident from the narrow spread in their energy evolution, a3-a4). In the uncooled regime, the energy gain is more chaotic as ions constantly scatter inside the current sheet.

the even-numbered ones, we pick ions with $\gamma_{\text{in}} \approx 5\sigma$ (the middle of the tail of the distribution).

If we consider the population of particles with energies $\gamma \approx \sigma$ (consider a slice at $t = 0$ in Fig. 2.8b1, b3), we see that the motion of ions is either not beamed at all in the y -direction (the case with no cooling), or has only marginal beaming (the case with strong cooling). By this time, the energy gain of these ions, $\gamma_{\text{in}} \approx \sigma$, is dominated by the Fermi acceleration in the current sheet (either via a Fermi kick or the “pick-up” mechanism). This prelude is consistent with the fact that the ions at these energies are primarily moving in the plane of the reconnecting magnetic field, x - z , i.e., perpendicular to the reconnection electric field, $\beta_y \approx 0$, (cf. French et al., 2022). At a later time, a fraction of these particles begin accelerating upstream of the layer along the y direction, i.e., along $\mathbf{E}_{\text{rec}}$; that fraction of “lucky” particles ultimately determines the power-law slope in the distribution. As it is evident from the first column of Fig. 2.8, in the uncooled run only a small fraction of ions follows the $\dot{\gamma} \propto \mathbf{E}_{\text{rec}} \cdot \mathbf{v}$ linear acceleration path, with the bulk beaming in y -direction remaining marginal. This is due to the fact, that the cross-section of typical plasmoids in the simulation with weak-to-no cooling is significantly larger (see Fig. 2.2), making them more capable of trapping the freely accelerating ions. On the other hand, in the strong cooling case almost all of the particles that reach $\gamma_{\text{in}} \approx \sigma$ are “lucky” to land on the efficiently accelerating trajectories, beamed in the direction of the electric field in the upstream (see the third column in Fig. 2.8; these trajectories are similar to the one shown in the right column of Fig. 2.7). As the “lucky” particles become more energized, they are essentially bound to remain on these accelerating trajectories. This can be clearly seen in the second and the fourth panels of Fig. 2.8, where we show the same tracking statistics, but only for the particles that have already reached $\gamma_{\text{in}} \approx 5\sigma$. For this selection, the fraction of “lucky” particles is substantially larger. This suggests that even in the uncooled case the ion energization mechanism to $\gamma \gg \sigma$ is

the same large-scale acceleration in the upstream, albeit a smaller amount of accelerated particles, free-streaming in the upstream, results in a steeper power-law slope in their distribution.

We observe a general tendency that the acceleration rate for individual “lucky” ions is larger for the strongly cooled simulations. This can be clearly seen from the typical strength of the electric field experienced by the ions in Fig. 2.8c1...c4: for stronger cooling the electric field is systematically larger than in the uncooled case, even if we consider particles beyond $\gamma_{\text{in}} \approx 5\sigma$. We associate this difference with the fact that the global reconnection rate, β_{rec} , is larger for the strong cooling case, as shown in Fig. 2.3, which in turn means larger $E_y = E_{\text{rec}} \approx \beta_{\text{rec}} B_{\text{up}}$. Another notable feature is the evolution of the distribution of electric field values, E_y , felt by ions (especially for $\gamma_{\text{in}} \approx \sigma$, panels c1 and c3). At early times, while particles still have relatively low energies, they move primarily in the current sheet, where the electric field is mostly chaotic. When accelerated to large enough energies, most of these ions are moving primarily in the upstream region (cf. Fig. 2.7), where the electric field is coherent and almost exclusively points perpendicular to the magnetic field direction y .

2.5.2 Pair acceleration

Acceleration mechanisms for pairs in the weak-to-no cooling regimes are identical to that for ions: pairs are first accelerated in the X-points and their outflows by non-ideal and ideal electric fields up to Lorentz factors comparable to σ , after which a fraction of them escape to the upstream where they are linearly accelerated further by the large-scale E_{rec} . However, when the cooling is enabled, even if it is dynamically weak, $\gamma_{\text{syn}} \gtrsim \sigma$, this linear acceleration cannot energize the pairs to arbitrarily high energies. In this case, the typical value of the effective

perpendicular magnetic field, $\tilde{B}$, defined in (2.3), for particles accelerating upstream is close to B_{up} . This implies that the upstream acceleration in the weak cooling regime is limited by the burnoff limit, $\gamma \lesssim \gamma_{\text{syn}}$. In Fig. 2.5, this effect can clearly be seen by comparing the uncooled simulation with the one with $\gamma_{\text{syn}} = 5\sigma$. While in the uncooled case, the distributions of both pairs and ions extend to several tens of σ in energy, in the weakly cooled simulation the distribution of pairs essentially cuts off at $\approx \gamma_{\text{syn}} = 5\sigma$, with ions being relatively unaffected.

In the strong cooling regime, the upstream acceleration is essentially prohibited, as the strong synchrotron cooling disallows particles to leave the X-points by crossing magnetic field lines. As shown in Fig. 2.5, the spectrum of accelerated leptons does not extend beyond $\approx \sigma$. In this regime, two acceleration mechanisms are particularly efficient: the X-point acceleration by the non-ideal electric field, and the “pick-up” mechanism when particle is injected non-adiabatically into a high velocity reconnection exhaust. In Fig. 2.6b, we show that particle energy gain in the $E > B$ zones is not sufficient to extend the power-law to $\gamma \approx \sigma$; this implies that pairs with the highest energies are additionally accelerated by the “pick-up” mechanism.

An example of a typical high-energy electron trajectory in the strong cooling regime is shown in Fig. 2.9 (3dx2kCoo102). In panels (a)...(e) we plot the x - z slices of the plasma density, and the magnetic field lines at different timesteps, with the current positions of the particle indicated with colored circles. The instantaneous four velocities of the particle are indicated with colored arrows. We also over-plot in blue the magnetic field line at the position of the particle, with the width of the line indicating its strength. On the right of Fig. 2.9 we show the time evolution of the Lorentz factor γ , with the dashed lines indicating the times of the snapshots on the left. In panel (a) the particle is upstream from the layer, moving towards the current sheet with a small Lorentz factor, $\gamma \sim 1$. In panel (b), the particle enters the X-point and is acceler-

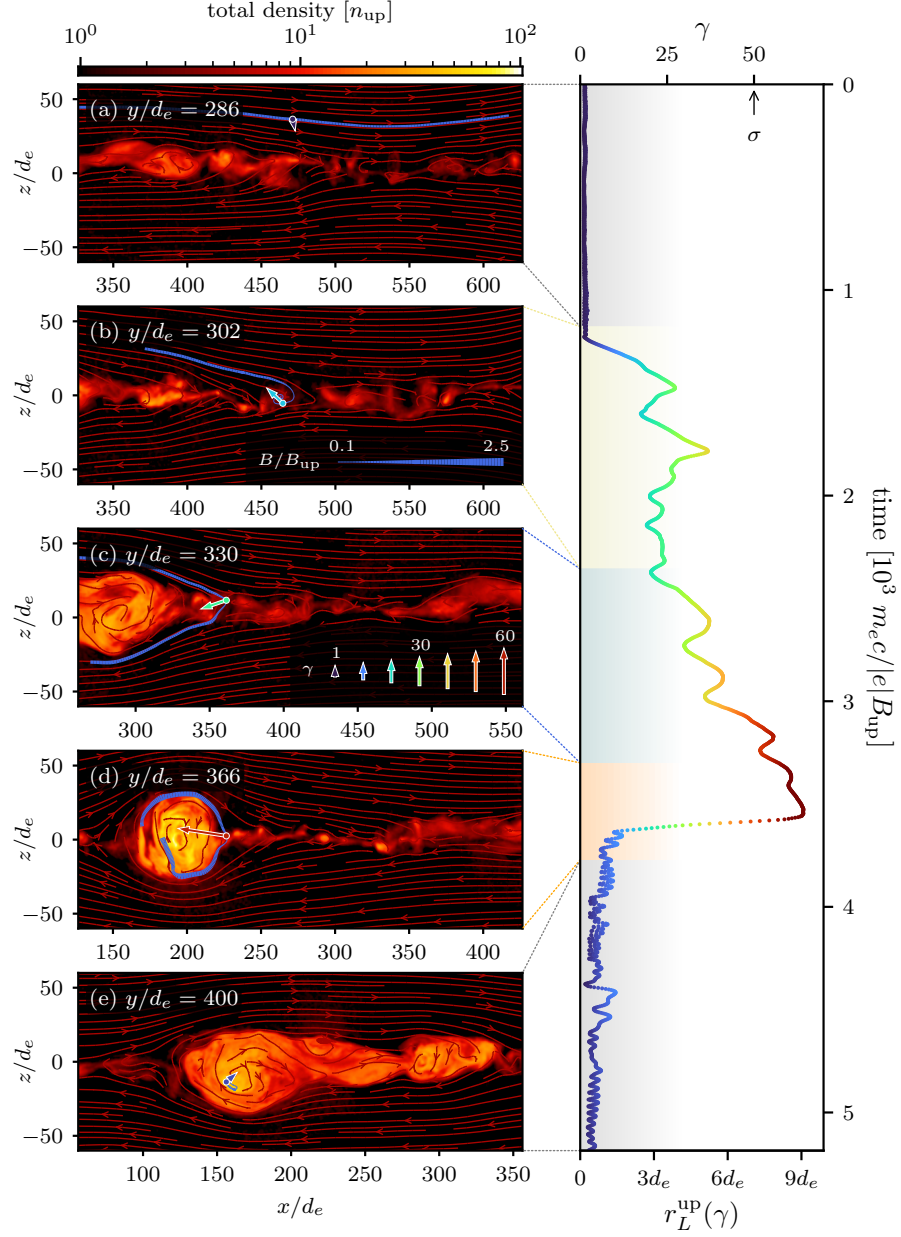


Figure 2.9: Time evolution of a typical accelerated electron in the strongly cooled simulation. The right panel shows the energy evolution of the particle with time. Panels (a)...(e) show the density slice and the magnetic field lines in the x - z plane at a particular time of the simulation (corresponding to the right panel). The position and the velocity of the electron are shown with a colored circle and an arrow (color corresponds to its energy at that particular time). The magnetic field line passing through the position of the particle is shown as a blue line, with its thickness corresponding to the strength of the magnetic field. The electron starts cold upstream of the current layer (a) and is accelerated to the Lorentz factor up to $\approx \sigma$ in the current sheet (b), after which it is picked up by the outflow from the X-point and gains energy further (c). Ultimately, the electron experiences an intensive slow-down when it collides with the magnetic field inside the plasmoid (d), where it remains throughout the rest of its lifetime (e).

ated by the non-ideal electric field, $E > B$, to $\gamma \approx 0.5\sigma$. After this rapid acceleration event, the particle moves along with the relativistic outflow from the X-point, constantly crossing the upstream/downstream boundary, which is perturbed vertically due to RDKI (see below). As the particle hits the upstream, it experiences subtle energy losses indicated as oscillations in its energy, γ . After that, the particle is picked up by a contracting magnetic field line performing a slingshot-like motion. After that, as shown in panel (c), for an extended period of time particle moves along with a rapidly receding field line while being slowly accelerated to $\gamma \gtrsim \sigma$. Once the slingshot-like motion of the field line comes to a halt, the particle is injected into the forming plasmoid, and experiences a catastrophic cooling event, being exposed to a large $\tilde{B}_\perp$ (panel (d); this event is studied in more details later). During its interaction with the plasmoid, the particle radiates most of its momentum and is then trapped inside the plasmoid until the end of the simulation, with a low Lorentz factor, $\gamma \sim \text{few}$, as shown in panel (e).

We further describe the 3D motion of the same accelerated electron in Fig. 2.10a...d, where we show 3D volume renderings of the plasma density and the magnetic field lines in the vicinity of the electron, also shown in Fig. 2.9. Panels on the right display the energization history of the particle, as well as the interesting quantities zoomed at two distinct moments of the particle's evolution (roughly corresponding to panels (a)/(b), and (c)/(d) of the same figure). In the upper half of Fig. 2.10, we specifically focus on the time when the particle experiences successive scattering events while moving in the RDKI-perturbed current layer, after being pre-energized in the X-point (corresponding to Fig. 2.9b). As indicated in the upper right panels, just before the scattering event, the particle crosses the null line, $|\mathbf{B}| \approx 0$, and approaches the peak of the RDKI perturbation, where both $\tilde{B}_\perp$ and $|\mathbf{B}|$ grow rapidly (time $t \approx 1.8$ in units indicated on the plot; field strengths at particle position are shown with blue and orange lines). $\tilde{B}_\perp$ reaches the critical

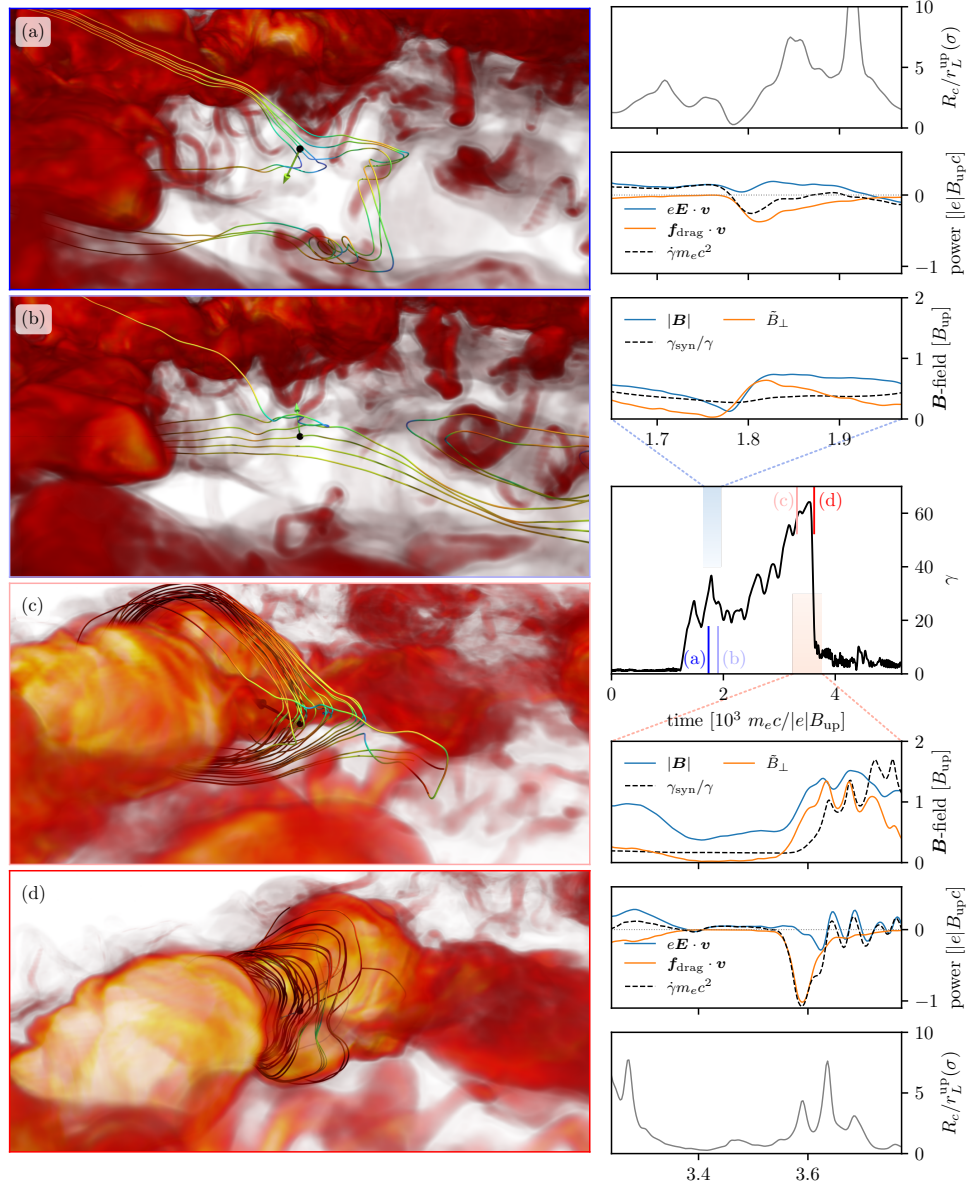


Figure 2.10: Close-up view of the acceleration dynamics of the electron shown in Fig. 2.9. Panels (a...d) show 3D snapshots of the density zoomed in at the position of the particle (shown with a black sphere). We also show the vector of its velocity and the magnetic field lines close to the location of the particle. Line colors represent the amplitude of the magnetic field, with the colorbar being identical to the one in Fig. 2.1. The middle panel on the right shows the energy evolution of the electron, with the moments of the 3D snapshots indicated by vertical lines. Three extra panels at the top and the bottom of the right column show the time evolution of several important quantities, limited in time to the intervals of the scattering (top) and the catastrophic cooling events (bottom). In particular, we plot the strength of the magnetic field and the $\tilde{B}_\perp$, experienced by the electron, as well as the critical magnetic field, $B_{up}\gamma_{syn}/\gamma$. We also show the contributions of two forces, the acceleration by the electric field, $e\mathbf{E} \cdot \mathbf{v}$, and the cooling, $\mathbf{f}_{drag} \cdot \mathbf{v}$, to the energy evolution of the particle, $\dot{\gamma}$. Finally, we show the curvature radius of the magnetic field at the location of the particle.

value, $\approx B_{\text{up}}\gamma_{\text{syn}}/\gamma$ (see below), after which deceleration occurs due to the radiative drag force, f_{drag} (shown with an orange line), with the particle being slowed down until the field drops below critical (time $t \approx 1.9$).

The catastrophic cooling event, shown in the lower half of Fig. 2.10 (panels c,d; see also Fig. 2.9d), is different from the subtle scattering event discussed above. In this case, the particle is accelerated by the “pick-up” mechanism to energy $\gamma \gtrsim \sigma$, while moving along a slingshot-like trajectory. Before $t \approx 3.5$, although the velocity of the particle is perpendicular to the local magnetic field (as indicated in Fig. 2.10c, and Fig. 2.9c, and d), the effective perpendicular magnetic field, $\tilde{B}_{\perp}$ is almost zero, allowing the particle to avoid being radiatively cooled. At the time $t \approx 3.55$, when the particle enters into the dense flux tube, $\tilde{B}_{\perp}$ increases dramatically (above the critical value, $B_{\text{up}}(\gamma_{\text{syn}}/\gamma)$), and the particle experience a strong radiative drag force (as shown in the second panel on the lower-right of Fig. 2.10 with an orange line, f_{drag}). In this catastrophic cooling episode, the particle loses most of its energy and is further trapped inside the flux tube (also shown in Fig. 2.9e).

In both the scattering and the catastrophic cooling examples, the full radiation reaction force, given by Eq. 2.3, consists of both synchrotron and curvature contributions. As shown in the top and the bottom panels on the right in Fig. 2.10, in both cases the local radius of curvature of the field line⁹ during the deceleration is comparable to the characteristic Larmor radius of the particle with the Lorentz factor $\gamma \approx \sigma$, i.e., it is close to the gyro-radius of the particle. Under these conditions, the synchrotron and curvature radiation mechanisms are indistinguishable, with the radius of curvature of field lines being microscopic and set by the local field perturbations inside the current layer, i.e., the curvature radius is comparable to the characteristic sizes of the

⁹Calculated as $((\hat{\mathbf{b}} \cdot \nabla)\hat{\mathbf{b}})^{-1}$, where $\hat{\mathbf{b}} \equiv \mathbf{B}/|\mathbf{B}|$ is a unit vector along the local magnetic field.

flux tubes.

After the acceleration phase, pairs experience a highly inhomogeneous magnetic field while propagating inside the current sheet. Synchrotron cooling strength for a given particle with a Lorentz factor γ is controlled by the effective perpendicular magnetic field strength at the position of the particle, $\tilde{B}_\perp$. Assuming a balance between the energy gain from the electric field of characteristic strength, $|e|E_{\text{rec}}$, and the radiative drag force, $|e|E_{\text{rec}}(\tilde{B}_\perp/B_{\text{up}})^2(\gamma/\gamma_{\text{syn}})^2$, one can define the critical value for the effective magnetic field: $\tilde{B}_\perp \approx B_{\text{up}}(\gamma_{\text{syn}}/\gamma)$. As discussed above, pairs can experience short episodes of stronger $\tilde{B}_\perp$, when they rapidly cool down on the timescales of gyration. This condition thus implies, that in simulations with a strong cooling, the highest-energy pairs (with $\gamma \approx \sigma$) that sustain their motion for long periods of time experience a significantly weaker effective perpendicular magnetic field, either biasing towards local weak-field regions or outflows from the X-points.

To directly measure the effect of energy-dependent field inhomogeneities, in Fig. 2.11 we plot the distribution of energies, γ , and effective magnetic field strengths experienced by particles, $\tilde{B}_\perp$, for pairs in simulations with varying cooling strength. The dashed line indicates the position of the critical magnetic field, $\tilde{B}_\perp \approx B_{\text{up}}(\gamma_{\text{syn}}/\gamma)$. Evidently, for all the cases particle distributions do not cross the critical threshold, with the highest-energy particles, $\gamma \approx \sigma$, following trajectories with $\tilde{B}_\perp \lesssim B_{\text{up}}(\gamma_{\text{syn}}/\gamma)$ (cf. [Hakobyan et al., 2023b](#)). Importantly, the combination of particle energy, γ , and the effective magnetic field it experiences, $\tilde{B}_\perp$, determines the peak energy of synchrotron emission, $\varepsilon_p \propto \gamma^2 \tilde{B}_\perp$, and in the context of the highest energy pairs, this determines the cutoff of the overall radiation spectrum. The results shown in Fig. 2.11 suggest that the naive prediction of the radiation spectrum from the pair distribution (or vice-versa) typically used in one-zone models can be significantly modified by this non-trivial dependence of $\tilde{B}_\perp$ on

the particle energy, especially in systems with dynamically strong synchrotron cooling.

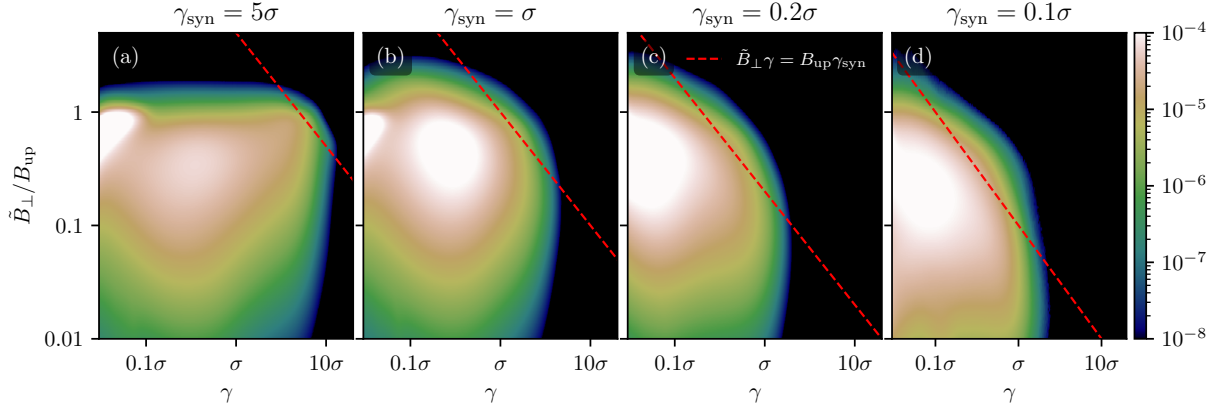


Figure 2.11: The effective perpendicular magnetic field, $\tilde{B}_\perp$, experienced by the pairs of different energies, γ . The 2D histograms are shown for all the cooling regimes, from weak, $\gamma_{\text{syn}}/\sigma = 5$, to strong, $\gamma_{\text{syn}}/\sigma = 0.1$. The red dashed lines show the estimated critical magnetic field, $\tilde{B}_\perp\gamma = B_{\text{up}}\gamma_{\text{syn}}$. Notably, the acceleration of the highest-energy pairs is inhibited by the synchrotron cooling, and, as a result, most of the pairs do not cross the critical line.

As we have seen above, in strongly cooled simulations inhomogeneities of the magnetic field in the current layer do not allow pairs to accelerate indefinitely. In particular, we can estimate the characteristic timescale of catastrophic energy losses for the highest-energy pairs as $t_c \lesssim \gamma/(\beta_{\text{rec}}\omega_L) \approx \beta_{\text{rec}}^{-1}(r_L^{\text{up}}(\sigma)/c)(\gamma/\sigma)$, where $\omega_L = |e|B_{\text{up}}/(m_e c)$.¹⁰ We directly measure the continuous residence time, Δt , at high energies, $\gamma \gtrsim \gamma_{\text{thr}}$, for all the pairs in our strongly cooled runs ($\gamma_{\text{syn}} = 0.2\sigma$). The distribution of residence times, normalized to $r_L^{\text{up}}(\sigma)/c$, is shown in Fig. 2.12 for two values of γ_{thr} : $2\gamma_{\text{syn}}$, and σ . To ensure our results are independent of the size of the box, we also over-plot the results with varying domain sizes and separation of scales. Evidently, most of the particles spend a very short amount of time, $\Delta t \lesssim t_c$, at high energies, after which they rapidly cool down, as we have discussed in the example above. Notice, that this timescale is typically much shorter than the light-crossing time of our simulation domain, L_x/c ,

¹⁰Here we used the definition of γ_{syn} from (2.4) and approximated $\tilde{B}_\perp \approx B_{\text{up}}(\gamma_{\text{syn}}/\gamma)$. Then one can write $\gamma/t_c = f_{\text{drag}}/(m_e c) = (4\pi m_e c)^{-1}\sigma_T B_\perp^2 \gamma^2$, and substitute γ_{syn} , and $\tilde{B}_\perp$.

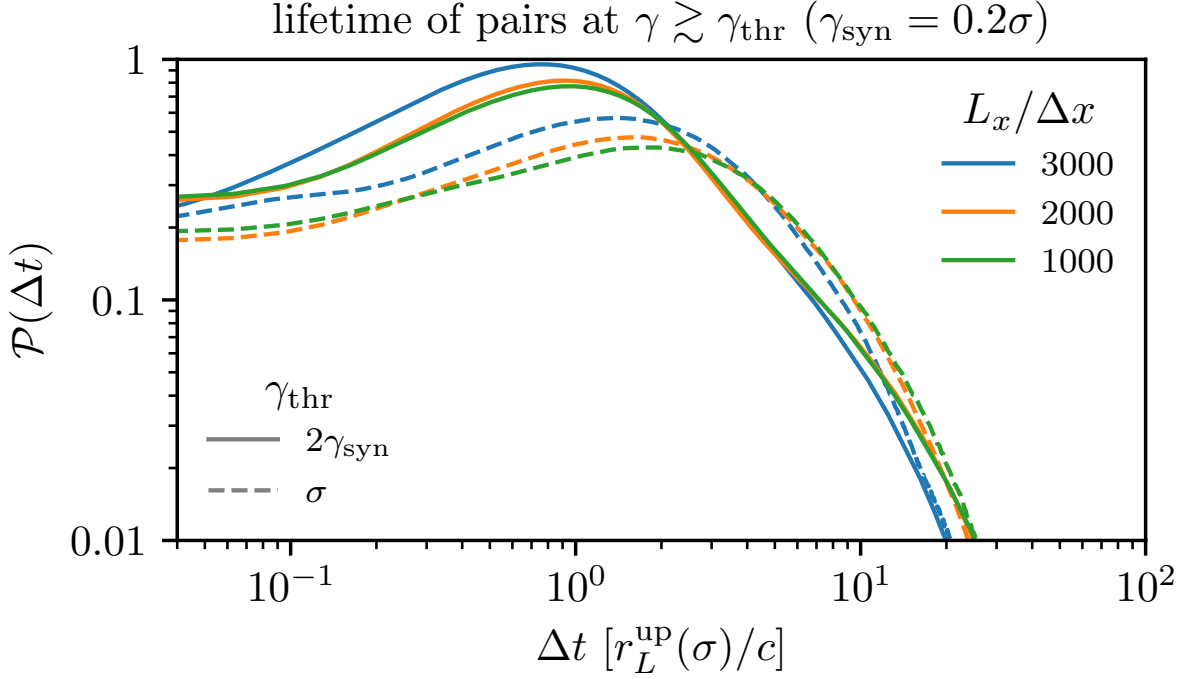


Figure 2.12: Distribution of the lifetimes of strongly cooled pairs at high energies ($\gamma \gtrsim 2\gamma_{\text{syn}}$ and $\gamma \gtrsim \sigma$) for different sizes of the simulation domain. For both thresholds, the lifetime is close to one gyro-period of the most energetic particles, $r_L^{\text{up}}(\sigma)/c$ and, importantly, is independent of the domain size.

which implies two important points. First of all, this indicates that our simulation domain is large enough so that the pairs are free to experience many acceleration-deceleration events during their lifetime inside the current layer. Despite this, pairs are unable to gain energy indefinitely in successive acceleration events, as each of these episodes is followed by a rapid energy loss. Second, the statistical balance between these two episodes essentially enforces the critical field criterion, $\tilde{B}_\perp \gamma \approx B_{\text{up}} \gamma_{\text{syn}}$, allowing us to extrapolate the upper bound on pair acceleration to the current layers with realistic scale separations.

2.6 Radiation from the current sheet

The primary probe of the underlying acceleration physics from the reconnecting sheets in astrophysical systems is the properties of the outgoing high-energy radiation. So far we have only focused on the acceleration dynamics of ions and pairs under the influence of the synchrotron drag force, in this section we study the properties of the produced emission.

In Sec. 2.5.2, we demonstrated that in the weak-to-moderate cooling regime, $\gamma_{\text{syn}} \gtrsim \sigma$, pair acceleration is limited to $\gamma_{\text{max}} \lesssim \gamma_{\text{syn}}$. We expect the spectrum of the high-energy radiation to cut off at energies comparable to the burnoff limit, ε_c , from (2.5). In contrast, in the strong cooling regime, $\gamma_{\text{syn}} \lesssim \sigma$, pairs can accelerate to $\gamma_{\text{max}} \sim \sigma$ (see, e.g., Fig. 2.5), and we expect the radiation spectrum to extend above the burnoff limit, as follows from Eq. (2.6). Following the discussion in Sec. 2.3.2, the estimate of the cutoff energy for the strong cooling regime can be obtained from the critical value of the effective perpendicular magnetic field, $\tilde{B}_\perp \approx B_{\text{up}}(\gamma_{\text{syn}}/\gamma)$:

$$\varepsilon_{\text{cut}} = \varepsilon_c \left(\frac{\tilde{B}_\perp}{B_{\text{up}}} \right) \left(\frac{\gamma_{\text{max}}}{\gamma_{\text{syn}}} \right)^2 = \varepsilon_c \left(\frac{\sigma}{\gamma_{\text{syn}}} \right). \quad (2.8)$$

To verify our expectations, in Fig. 2.13, we show photon spectra, $\varepsilon^2 dn_\varepsilon/d\varepsilon \equiv \varepsilon^2 f_\varepsilon$, from simulations with different cooling regimes. Each of these spectra is averaged in the interval of $1.5 < tc/L < 4.5$. To construct the spectrum, we assume that every particle emits a continuous synchrotron spectrum (e.g., Rybicki & Lightman, 1979):

$$\varepsilon \frac{dn_\varepsilon}{d\varepsilon} \propto \frac{\varepsilon}{\varepsilon_p} \int_{\varepsilon/\varepsilon_p}^{\infty} K_{5/3}(x) dx, \quad (2.9)$$

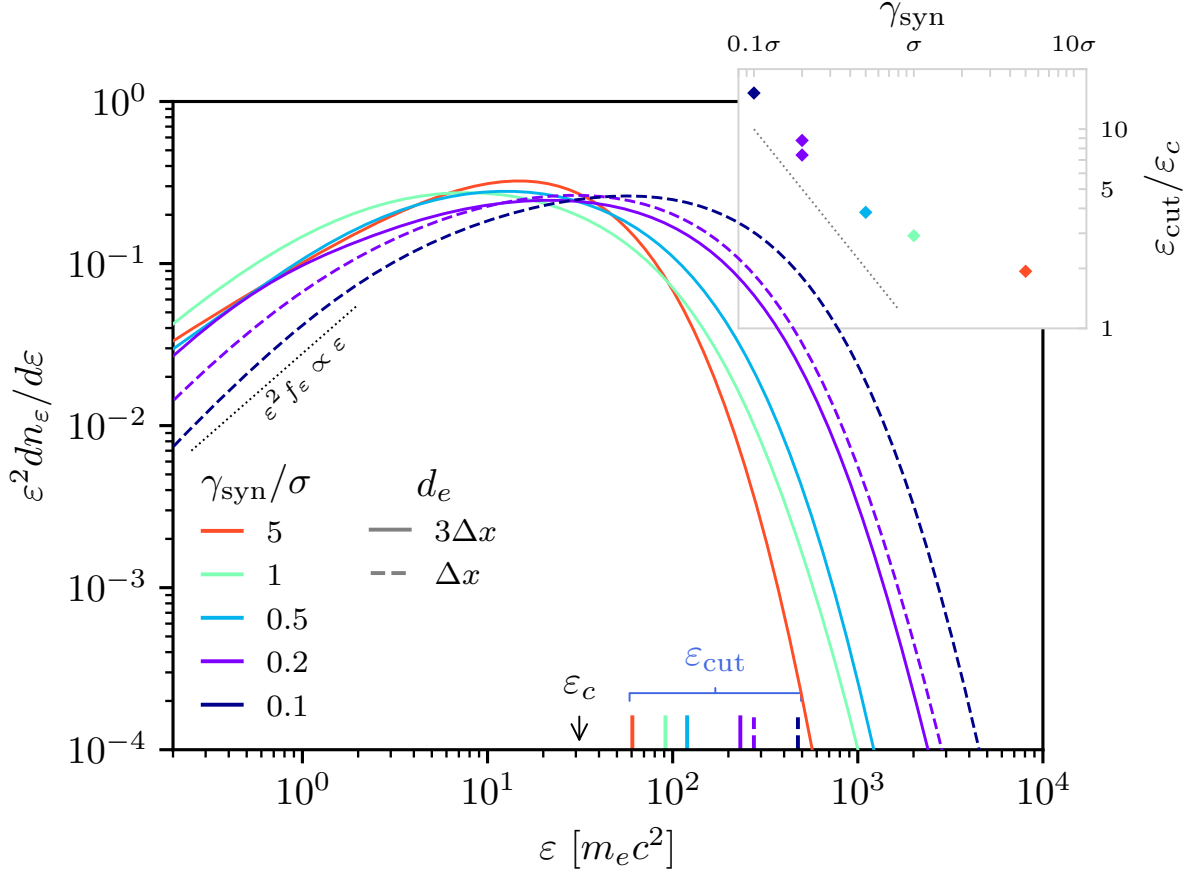


Figure 2.13: Normalized flux density of the emitted radiation for different cooling regimes. The additional panel shows the position of the exponential cutoff of the spectra as a function of the cooling strength. Current sheets with stronger cooling produce a more extended emission spectrum, with significant power above the burnoff limit, ϵ_c , and a cutoff being in good agreement with the prediction of Eq. 2.8: $\epsilon_{\text{cut}} \sim \epsilon_c (\sigma/\gamma_{\text{syn}})$. Current sheets with weak cooling, on the other hand, produce almost no emission above the burnoff limit.

where $K_{5/3}$ is the modified Bessel function, and the energy ϵ_p is defined according to Eq. (2.6). All of the spectra rise linearly at low energies, $\epsilon^2 f_\epsilon \propto \epsilon$, which is expected from the hard power-law of reconnection-accelerated particles, $f \sim \gamma^{-1}$. The rise is then followed by a broad peak, and a near-exponential decay at high energies. To find the position of the exponential cutoff, we fit the high-energy parts of the spectra as $\epsilon^2 f_\epsilon \propto \exp(-\epsilon/\epsilon_{\text{cut}})$. The inferred values, ϵ_{cut} , are shown in Fig. 2.13 with colored vertical lines, as well as in the inset panel at the top right.

Evidently, our results are consistent both with the prediction of Eq. (2.8) for the strong cooling, $\varepsilon_{\text{cut}} \approx \varepsilon_c(\sigma/\gamma_{\text{syn}})$, as well as for the weak cooling, where the cutoff approaches $\varepsilon_{\text{cut}} \approx \varepsilon_c$.

The angular distribution of the produced emission is another milestone toward applying our results to realistic astrophysical systems. Previous analysis of the emission anisotropy was done theoretically (e.g., [Uzdensky et al., 2011](#); [Cerutti et al., 2012a](#)), in 2D simulations ([Cerutti et al., 2012b,a, 2013](#); [Kagan et al., 2016a](#); [Mehlhaff et al., 2020](#)) or in 3D simulations with a limited separation of scales ([Cerutti et al., 2014](#)). In Fig. 2.14 we show the distribution of photon energies, $f(\theta, \varepsilon)$, with respect to θ , the angle between the photon (emitted along the instantaneous motion of the electron/positron) momentum and the upstream magnetic field, $\hat{x}$, for different values of the cooling strength. In the weakest cooling case, $\gamma_{\text{syn}} = 5\sigma$ (Fig. 2.14a), the highest energy pairs are accelerated in the upstream, along the reconnection electric field, $E_{\text{rec}} \approx E_y$. This results in high-energy photons being emitting preferentially perpendicular to the upstream magnetic field ([Zhang et al. 2021a](#), and Sec. 2.5.2). For the strong cooling (Fig. 2.14c, d), the emission of photons with energies above the burnoff limit, $\varepsilon \gtrsim \text{few} \cdot \varepsilon_c$, is strongly anisotropic and is beamed towards the direction of the upstream field. The bulk of these photons is emitted by the mechanism described in Sec. 2.5.2 and particularly demonstrated in the lower half of Fig. 2.10: the particle is accelerated by the “pick-up” mechanism in the X-point outflow, in the direction of the upstream magnetic field (Fig. 2.10c), and later produces the short radiation burst, after colliding face-on with the plasmoid.

To understand the spatial distribution of emitting regions, in Fig. 2.15 we plot a volumetric map of the radiated synchrotron power, $P_{\text{sync}} = \sum_{e^\pm} \mathbf{F}_{\text{rad}} \cdot \mathbf{v}$, with the $\mathbf{F}_{\text{rad}}$ defined in Eq. (2.3) for simulations with weak (Fig. 2.15a), and strong cooling (Fig. 2.15b), (the summation, $\sum_{e^\pm}$, is done over all the pairs in the given cell). We separately plot the total radiated power (green)

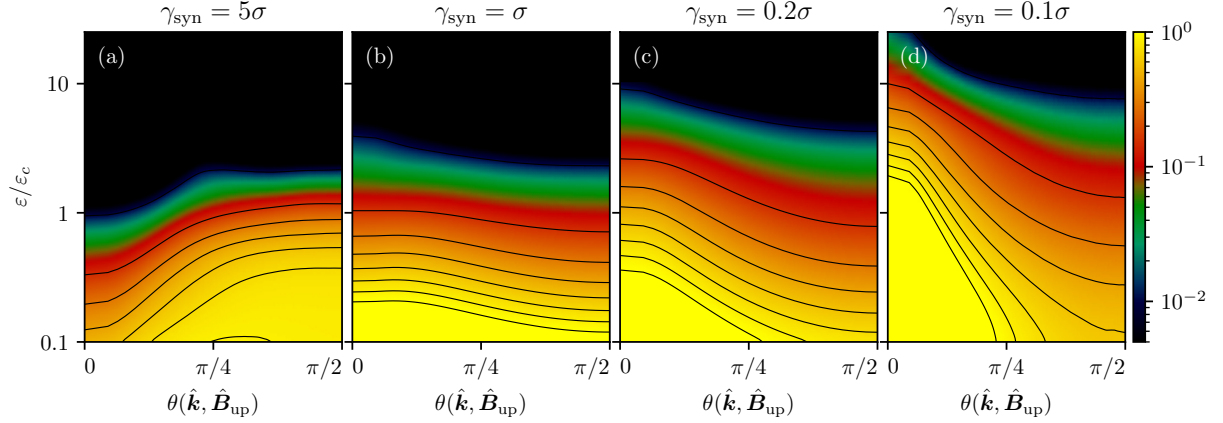


Figure 2.14: Angular anisotropy of the radiation at different photon energies for different cooling regimes. The angle is measured with respect to the direction of the upstream magnetic field $\mathbf{B}_{\text{up}} \equiv B_x \hat{\mathbf{x}}$. In the weak cooling case, $\gamma_{\text{syn}} = 5\sigma$, emission above the burnoff limit is significantly suppressed, and most of the high-energy photons are emitted perpendicular to the upstream magnetic field. On the other hand, in the strong cooling regime, most of the high-energy photons above the burnoff limit are emitted along the upstream magnetic field.

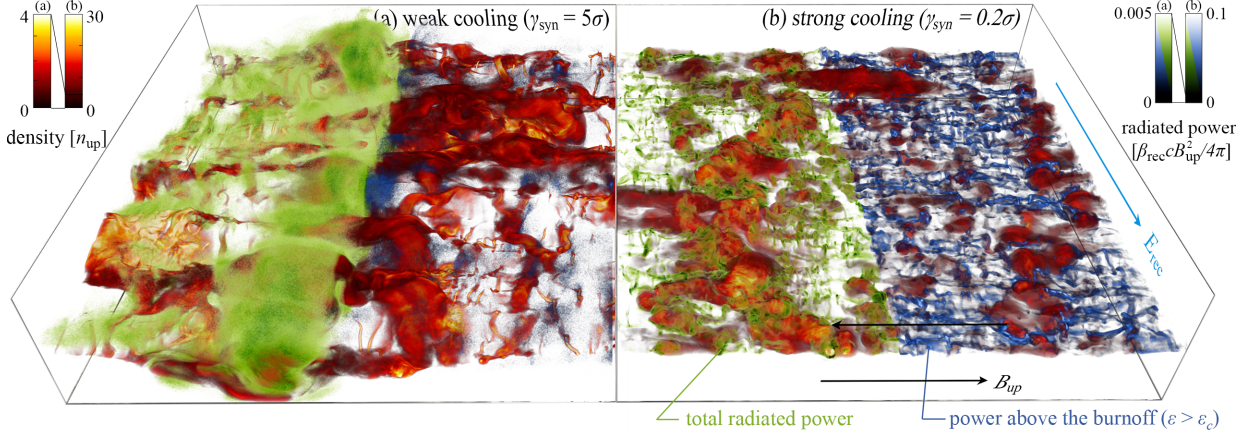


Figure 2.15: Spatial distribution of regions producing the synchrotron radiation shown together with the total plasma density, for the same snapshots as in Fig. 2.1. Green volume rendering corresponds to the total synchrotron intensity, while the blue one shows the emission intensity above the burnoff limit, ε_c . Panel (a) corresponds to the weak cooling, while panel (b) is for the strong cooling. The emission in the strongly cooled current sheet is mainly associated with the peaks of the RDKI and the outer shells of the plasmoids (both for the total emission and the emission above the burnoff). The much fainter emission from the weakly cooled current sheet (note the difference in color scales between the two regimes) is significantly more uniform, with the emission above the burnoff being essentially absent.

and the power carried by the photons above the burnoff limit, $\varepsilon > \varepsilon_c$ (blue). For visual guidance, we use the same snapshot and viewing angle as in Fig. 2.1, and also plot the total plasma density. To properly quantify the radiated power, we normalize it by the incoming Poynting flux, $\beta_{\text{rec}} c B_{\text{up}}^2 / (4\pi)$.

In the case of weak cooling (Fig. 2.15a), the density of radiation losses is very weak, i.e., significantly below the Poynting flux inflow (indicated by the colorbar scales for panels (a) and (b)), substantially uniform, and marginally biased towards plasmoids. This is due to the fact that a substantial fraction of the highest energy particles are accelerated by E_{rec} in the upstream of the current sheet, and trapped by plasmoids. On the other hand, for the strong cooling regime (Fig. 2.15b) a large fraction ($\sim 10\%$) of the incoming Poynting flux is being radiated away, with most of the power going into photons comparable to and above the burnoff limit, ε_c . There are two main zones producing the radiation in this case: the outflows from the X-points, and the boundaries of plasmoids. The former is associated with the emission of the X-point-accelerated particles, which interact with the small-scale field inhomogeneities in the outflows of the current sheet (see also Fig. 2.10a, b). Radiation near the plasmoid boundaries is associated with catastrophic cooling events when the energetic particles collide face-on with the strong magnetic field surrounding the plasmoids (Fig. 2.10c, d). These particles lose energy very intensively at a very short timescale and are ultimately trapped by plasmoids while having very low Lorentz factors, $\gamma \ll \sigma$ (also visible in Fig. 2.9e). Because of this, in the strong cooling regime, plasmoid interiors have almost no contribution to the total radiated power, in contrast to the weak cooling regime.

Emission above the burnoff limit (blue regions) is almost negligible in the case of weak cooling (Fig. 2.15a), which is consistent with our discussion above, and is also highlighted in the emission spectra in Fig. 2.13. For the strong cooling (Fig. 2.15b), the power above the burnoff

limit (blue) is similar to the total power (green), with the main difference being that the highest-energy emission is suppressed in plasmoid interiors. Thus, in the strong cooling regime, the inner volume of plasmoids only emit soft photons, with the photons above the burnoff being emitted primarily near their exteriors, in the outflows from the X-points, and at the local field inhomogeneities (such as the RDKI ripples, see also Fig. 2.10).

2.7 Discussion

In this paper, we study the 3D dynamics of magnetic reconnection in radiatively cooled pair-dominated plasma with a small mixture of ions. Our study focuses on different regimes of synchrotron cooling, parametrized by the burnoff Lorentz factor, γ_{syn} , which quantifies the dynamical importance of cooling compared to acceleration by the reconnection electric field. In the weak cooling regime, $\gamma_{\text{syn}} \gtrsim \sigma$, particle acceleration is effective both for pairs and ions, both of which form a power-law distribution function with an index $f(\gamma) \propto \gamma^{-1.7}$ that extends to energies substantially above σ . Here, the highest energy particles (ions and pairs) are mainly accelerated by the large-scale reconnection electric field upstream of the current sheet. In contrast to that, in the strong cooling regime, $\gamma_{\text{syn}} \lesssim \sigma$, the acceleration of pairs is severely limited by the radiative losses, and their maximum Lorentz factor reaches only $\gamma \approx \sigma$. Surprisingly, the acceleration of uncooled ions in this regime becomes more efficient as they form a very hard power-law distribution, with the distribution approaching γ^{-1} . As we uncovered in Sec. 2.5.1, this hardening of the power law in ion distribution is due to plasmoids being compressed by radiative cooling, which limits their ability to trap the accelerated particles and prevent acceleration in the reconnection electric field. Simultaneously, the strong cooling increases the filling fraction of

the X-points in the current layer, effectively enhancing the reconnection rate and, as a result, the strength of the electric field.

The difference in acceleration mechanisms results in a number of important implications for the properties of the high-energy radiation. In the case of weak cooling, the spectral cutoff is close to the synchrotron burnoff limit, $\varepsilon_c \approx 16$ MeV, as the particles accelerated in the upstream have large pitch angles and, thus, their energy cannot substantially exceed γ_{syn} . Conversely, in the case of strong cooling particles are accelerated up to $\gamma \approx \sigma \gtrsim \gamma_{\text{syn}}$ in X-points and relativistic outflows, which results in significant emission beyond the burnoff limit, with the cutoff being close to $\varepsilon_{\text{cut}} \approx \varepsilon_c(\sigma/\gamma_{\text{syn}})$. In addition, the difference in acceleration mechanisms results in a significant difference in the anisotropy of the high-energy radiation. For strong cooling, the most energetic photons are emitted by particles accelerated by the “pick-up” mechanism, which primarily move in the direction of the upstream magnetic field. Thus, the highest energy radiation is emitted in the direction along the upstream magnetic field. For the weak cooling, the most energetic pairs propagate along the reconnection electric field, and thus in this regime the most energetic photons are emitted along E_{rec} and perpendicular to the upstream magnetic field B_{up} .

2.7.1 Acceleration of ions in pulsar magnetospheres

Current sheets in the magnetospheres of pulsars are predominantly filled with the electron-positron plasma. However, a small mixture of ions is also expected (Arons, 2003; Amato & Arons, 2006; Fang et al., 2012; Guépin et al., 2020). Ions can be extracted from the atmosphere of the neutron star near the polar cap region and can carry volumetric and separatrix return current (Timokhin, 2006; Philippov & Spitkovsky, 2018). Depending on the inclination angle between

the magnetic and the rotational axes of the pulsar, a significant fraction of these ions can be carried toward the reconnecting current sheet. Due to their large mass, ions are relatively unaffected by the synchrotron cooling and may thus be accelerated more efficiently.

The assumption of negligible cooling for ions can be justified for the observed population of pulsars. Using the definition of the synchrotron burnoff Lorentz factor, Eq. (2.4), taking the Thomson cross-section of ions, $\sigma_T^i = \sigma_T Z^4 / A^2 (m_e/m_p)^2$, $m_i = Am_p$ (m_p is the mass of a proton), and $q_i = Z|e|$, we find that for the ions $\gamma_{\text{syn}}^i = AZ^{-3/2}(m_p/m_e)\gamma_{\text{syn}}$, with γ_{syn} being the burnoff Lorentz factor for pairs (see Sec. 2.3.2). Since most of the high-energy emission takes place in regions of the current layer close to the light cylinder, we may further employ $B_{\text{LC}} \approx B_\star (R_\star/R_{\text{LC}})^3$ as a proxy for B_{up} , with $R_\star$, and $B_\star$ being the radius of the neutron star, and the magnetic field strength near its surface, calculated as $B_\star = 10^{12} \text{ G} \sqrt{(P_\star/1 \text{ s})(\dot{P}_\star/10^{-15})}$, and $R_{\text{LC}} = cP_\star/2\pi$ being the light cylinder radius ($P_\star$ is the rotation period of the star). We then find

$$\gamma_{\text{syn}}^i \approx 10^{10} \frac{A}{Z^{3/2}} \left(\frac{P_\star}{1 \text{ s}} \right)^{3/2} \left(\frac{B_\star}{10^{12} \text{ G}} \right)^{-1/2}, \quad (2.10)$$

which can be compared with the total potential drop in the current layer, $q_i \beta_{\text{rec}} B_{\text{LC}} R_{\text{LC}}$, normalized to $m_i c^2$:

$$\gamma_{\text{max}}^i \approx 4 \cdot 10^3 \frac{Z}{A} \left(\frac{P_\star}{1 \text{ s}} \right)^{-2} \left(\frac{B_\star}{10^{12} \text{ G}} \right). \quad (2.11)$$

Then the assumption of negligible cooling for ions holds when $\gamma_{\text{syn}}^i \gtrsim \gamma_{\text{max}}^i$. We compare these two quantities for protons ($A = Z = 1$) for the population of about 3000 radio pulsars from the ATNF catalog (Manchester et al., 2005; The Australia Telescope National Facility (ATNF), 2023)

by overplotting the isocontours on the $P_\star - B_\star$ diagram in Fig. 2.16. The gray-shaded region in the parameter space indicates the strong cooling regime for ions, $\gamma_{\text{syn}}^i < \gamma_{\text{max}}^i$, and evidently, none of the observed radio pulsars fall into that category. This justifies our assumption made throughout the paper on neglecting the radiation drag for the ions. As seen from Fig. 2.16, a small population of the most energetic pulsars with small-enough rotation period and large surface magnetic fields (e.g., Crab) are able to accelerate protons to Lorentz factors $\gtrsim 10^6$, i.e., above PeV energies. For heavier nuclei with $Z \approx A > 1$, the cooling is slightly more efficient, as $\gamma_{\text{syn}}^i/\gamma_{\text{max}}^i \propto Z^{-1/2}$, and may thus inhibit their acceleration even for the most energetic pulsars. Acceleration near the light cylinder is also severely radiation-limited for the parameters of fast new-born magnetars (see, e.g., Arons, 2003; Fang et al., 2012). In this case, however, if the reconnection continues to be efficient further from the light cylinder, additional acceleration is possible in the wind, where the magnetic field is weaker and the radiative losses are less severe.

The mass density fraction of ions in pulsar current sheets is low, $f_i \approx (m_i/m_e)\alpha/\lambda \approx 0.01$ (which is the value we employed in our simulations), where $\lambda = n_\pm/n_{\text{GJ}} \sim 10^4$ is the pair multiplicity, and $\alpha = n_i/n_{\text{GJ}} \sim 0.1$ is the ion density relative to the Goldreich-Julian density ¹¹, $n_{\text{GJ}} = |\mathbf{\Omega} \cdot \mathbf{B}|/2\pi ce$, where $\mathbf{\Omega}$ is the angular velocity of the pulsar. The energy budget of even this small fraction of ions, $n_i m_i c^2 \langle \gamma_i \rangle$, can become substantial as they accelerate to high energies, which will inevitably lead to steepening in their distribution function. In the parameter regime where synchrotron cooling is strong for pairs ($\gamma_{\text{syn}} < \sigma_{\text{LC}}$), and weak for ions ($\gamma_{\text{syn}}^i \gtrsim \gamma_{\text{max}}^i$), the energy of pairs is limited to $\gamma_{\text{max}} \approx \sigma_{\text{LC}}$, while the distribution of ions can extend to the full potential drop, γ_{max}^i , Eq. (2.11). Assuming a power-law index of -1 for ions and $\langle \gamma_\pm \rangle \approx \gamma_{\text{syn}}$

¹¹Low fractions of the extracted ions are due to the fact that the bulk of the return magnetospheric current is carried by produced electron-positron pairs (e.g., Timokhin & Arons, 2013; Chen & Beloborodov, 2014; Philippov & Spitkovsky, 2018).

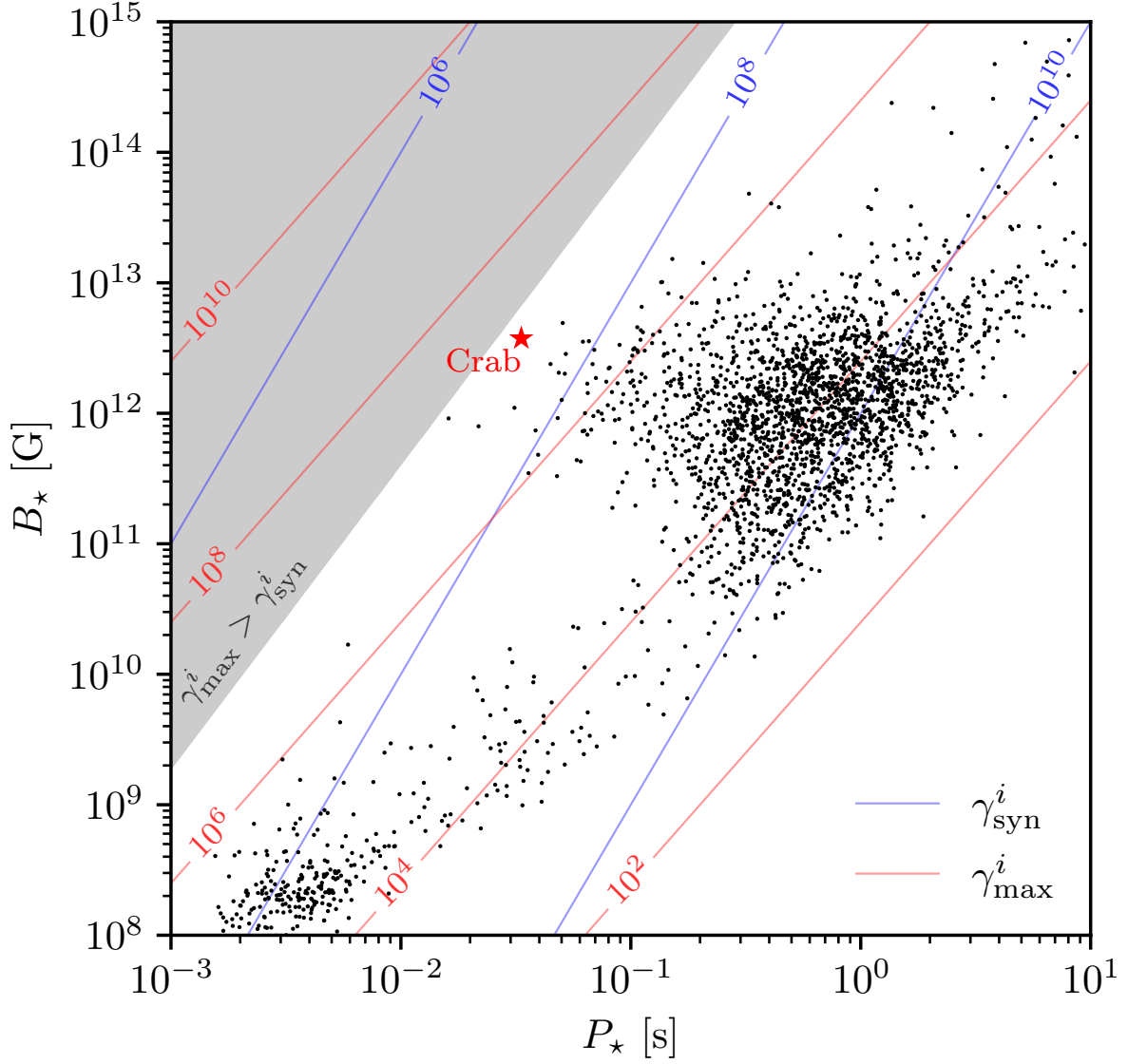


Figure 2.16: Estimates for burnoff limit of protons, γ_{syn}^i , (blue lines) and the Lorentz factor of protons accelerated in the full potential drop near the light cylinder, γ_{max}^i , (red lines) for the pulsars in the ATNF catalog. Axes correspond to the rotation period of the star, P_* , and the magnetic field strength near the surface of the star. The shaded region corresponds to the regime when the synchrotron cooling of protons becomes dynamically important: $\gamma_{\text{max}}^i > \gamma_{\text{syn}}^i$. Evidently, for all of the observed pulsars, the radiation drag of ions is negligible in the context of ion acceleration.

for pairs (cf. our simulation 3dx3kCool02), we may rewrite the total energy budget for both species as:

$$\begin{aligned}\mathcal{E}_{\pm} &\approx \lambda n_{\text{GJ}} m_e c^2 \gamma_{\text{syn}}, \\ \mathcal{E}_i &\approx \alpha n_{\text{GJ}} m_i c^2 \gamma_{\text{max}}^i.\end{aligned}\tag{2.12}$$

On the other hand, the magnetization near the light cylinder, σ_{LC} , can be expressed as $\sigma_{\text{LC}} = (P_*/4\pi\lambda)(|e|B_{\text{LC}}/m_e c) = 5\gamma_{\text{max}}^{\pm}/\lambda$, where $\gamma_{\text{max}}^{\pm} = \beta_{\text{rec}}|e|B_{\text{LC}}R_{\text{LC}}/m_e c^2$ is defined as the Lorentz factor of an electron or a positron, which taps β_{rec} fraction of the total voltage drop. We can thus find $\mathcal{E}_{\pm}/\mathcal{E}_i = (5/\alpha)(\gamma_{\text{syn}}/\sigma_{\text{LC}})$, where we used the relation $m_e \gamma_{\text{max}}^{\pm} = m_i \gamma_{\text{max}}^i$. For a Crab-like pulsar with $\gamma_{\text{syn}}/\sigma_{\text{LC}} \approx 0.1$, we find that pairs dominate in terms of the total energy budget, i.e., $\mathcal{E}_{\pm} \gg \mathcal{E}_i$, even when the ions are accelerated to the full voltage. This implies that the feedback of ions on the reconnection dynamics is negligible, and the ion spectrum extends to γ_{max}^i without significant deviations from the $f(\gamma) \propto \gamma^{-1}$ predicted in our simulations. In fact, we do not observe the spectrum to steepen even when $\mathcal{E}_{\pm}/\mathcal{E}_i \approx 1/3$, as happens in the simulation 3dx3kCool02 at late times.

Another channel of ion acceleration in pulsar magnetospheres is the energy gain by the parallel electric field in the discharge region close to the polar cap. The amplitude of the unscreened voltage is limited by pair cascade, which in young pulsars is triggered when the energetic pairs reach Lorentz factors, $\gamma_{\text{th}} \sim 10^6 \dots 10^7$. The same potential drop will correspond to the maximum Lorentz factor of protons $\gamma_{\text{th}}^i \sim 10^3 \dots 10^4$ (smaller by a factor of m_i/m_e). These values are much smaller than the ones corresponding to the acceleration in the current sheets of Crab-like pulsars, $\gamma_{\text{max}}^i \gtrsim 10^6$, reinforcing the notion that the acceleration in the current sheet beyond the light cylinder is the most efficient channel of ion energization in magnetospheres of young pulsars.

2.7.2 High-energy emission from young pulsars and supermassive black holes

In Sec. 2.6, we discussed the predictions for the observed high energy spectrum from the accelerated pairs. To summarize, in the strong cooling regime applicable in the context of the magnetospheric current sheets of young pulsars, the flux density of the synchrotron emission rises linearly $\nu F_\nu \propto \nu$ (with ν being the photon frequency) up to the energy comparable to the synchrotron burnoff limit, 16 MeV. The broad peak at energies higher than the burnoff is then followed by a steep cutoff at energies close to $16 \cdot (\sigma/\gamma_{\text{syn}})$ MeV (cf. [Hakobyan et al., 2023a](#)). For the young pulsars, such as the Crab, where the values for the gamma-ray cutoff are between one-to-few GeV ([Abdo et al., 2013](#)), our prediction yields $\sigma_{\text{LC}}/\gamma_{\text{syn}} \sim 100$ (i.e., the cooling regime is strong). For Crab pulsar, where $\gamma_{\text{syn}} \approx 4 \cdot 10^4$ close to the light cylinder, this implies that the magnetization parameter near the current layer is around $\sigma_{\text{LC}} \approx 10^6$. Notably, the presence of pairs with energies $\gtrsim \text{few} \cdot 10^6 m_e c^2$ is a strict requirement to explain the pulsed emission above TeV energies in the Crab pulsar ([Ansoldi et al., 2016b](#)).

The beaming of the high energy emission is another central outcome of our work. Modeling of gamma-ray lightcurves from global simulations of pulsar magnetospheres ([Bai & Spitkovsky, 2010](#); [Cerutti et al., 2016](#); [Kalapotharakos et al., 2018](#)) discovered that in order to explain the observed lightcurves the emission has to be beamed along the magnetic field lines, in the corotating frame of the pulsar. It has thus been not entirely clear how magnetic reconnection in the outer-magnetospheric current sheet could be responsible for such energization, as the naive expectation implied that particles were beamed in the direction of the reconnecting electric field, perpendicular to the direction of the magnetic field. Our findings outlined in Sec. 2.6 and in Fig. 2.14, however, indicate that in the regime of strong synchrotron cooling, $\gamma_{\text{syn}} < \sigma_{\text{LC}}$, most

of the high-energy emission is indeed beamed along the upstream magnetic field, as a result of particles being reaccelerated by the “pick-up” mechanism, while moving along the relativistic outflows from the X-points in the direction of the upstream field (see also [Cerutti et al. 2012b](#) for the 2D study).

Magnetospheres of low-luminosity accreting supermassive black holes have also been predicted to host intermittently formed reconnecting current layers with sizes comparable to few-to-ten black-hole gravitational radii (e.g., [Ripperda et al., 2020, 2022](#)). In particular, for the black hole at the center of the M87 galaxy, due to abundant pair-loading, the regime of reconnection is predicted to be somewhat similar to the reconnection in Crab pulsar ([Ripperda et al., 2022](#)), with $\sigma \sim 10^7$, and $\gamma_{\text{syn}} \sim 10^6 \dots 10^7$. In that regard, our simulations are largely consistent with the results of 2D simulations by [Hakobyan et al. \(2023b\)](#), where the authors predict that most of the power is radiated at 10...100 MeV range, with the most energetic pairs, $\gamma \sim \max(\sigma, \gamma_{\text{syn}}) \sim 10^7$, also producing the observed TeV signal (via inverse Compton scattering of soft disk photons). Remarkably, the luminosity ratio between the TeV signal, and the jet power, 0.1%, is close to the luminosity ratio between the TeV and GeV signal from the Crab pulsar, suggesting that the Compton amplification parameters for both of these vastly different systems are somewhat similar.

Finally, in most of the discussion above, we focused on the strong cooling regime, where significant emission is expected to be observed above the burnoff limit. One important instance where it is unclear whether this regime applies is the Crab Nebula, where reconnection has been proposed (see, e.g., [Cerutti et al., 2014](#); [Lyutikov et al., 2016](#)) as a possible mechanism for the gamma-ray flares at energies $\gtrsim 160$ MeV observed by the Fermi satellite ([Tavani et al., 2011](#); [Buehler et al., 2012](#)). Here, the upstream pair plasma is relativistically hot, $\langle \gamma \rangle \gg 1$,

and the characteristic energy gain per particle is given by the hot magnetization parameter, $\sigma_h = B^2/(4\pi n m_e c^2 h)$, where $h \sim \langle \gamma \rangle$ corresponds to the plasma enthalpy per particle. The strong cooling regime then applies only if $\sigma_h \langle \gamma \rangle \gtrsim \gamma_{\text{syn}} \sim 10^9$, where the estimate for γ_{syn} is made assuming the strength of the magnetic field is close to milligauss. Our findings suggest that a significant flux at energies above $\gtrsim 160$ MeV can only be produced when $\sigma_h \gtrsim 10 \cdot (\gamma_{\text{syn}}/\langle \gamma \rangle) \approx 1000(\gamma_{\text{syn}}/10^9)/(\langle \gamma \rangle/10^7)^{-1}$. It is currently unclear whether such conditions can be realized in the Nebula.

Chapter 3: Coherence of Multi-Dimensional Pair Production Discharges in Polar Caps of Pulsars

3.1 Abstract

We report on the first self-consistent multidimensional particle-in-cell numerical simulations of non-homogeneous pair discharges in polar caps of rotation-powered pulsars. By introducing strong inhomogeneities in the initial plasma distribution in our simulations, we analyze the degree of self-consistently emerging coherence of discharges across magnetic field lines. In 2D, we study discharge evolution for a wide range of physical parameters and boundary conditions corresponding to both the absent and free escape of charged particles from the surface of a neutron star. We also present the results of the first 3D simulations of discharges in a polar cap with a distribution of the magnetospheric current appropriate for a pulsar with 60° inclination angle. For all parameters, we find the coherence scale of pair discharges across magnetic field lines to be of the order of gap height. We also demonstrate that the popular “spark” model of pair discharges is incompatible with the universally adopted force-free magnetosphere model: intermittent discharges fill the entire zone of the polar cap that allows pair cascades, leaving no space for discharge-free regions. Our findings disprove the key assumption of the spark model about the existence of isolated distinct discharge columns.

3.2 Introduction

Pulsars are highly magnetized (up to $\sim 10^{13}\text{G}$) rapidly-spinning neutron stars (NS) (e.g., [Philippov & Kramer, 2022](#)). They emit pulsed non-thermal radiation across a wide range of frequencies generated by dense pair plasma in their magnetospheres. Most of the plasma is produced in quantum-electrodynamic (QED) discharges in polar caps – regions around magnetic poles close to the NS surface. There, a rotation-induced electric field pulls electrons from the atmosphere of the NS and accelerates them to high energies. As these electrons propagate along the curved magnetic field lines, they emit high-energy curvature photons, which subsequently decay into electron-positron pairs as first suggested by [Sturrock \(1971\)](#). For a long time, the de facto standard model for pair production assumed a stationary unidirectional space-charge-limited flow (SCLF) of plasma extracted from the NS atmosphere. Here, electrons are accelerated by the electric field due to charge starvation caused by a small mismatch between the local Goldreich-Julian (GJ, [Goldreich & Julian 1969](#)) and the actual plasma charge densities, $\rho - \rho_{\text{GJ}}$ ([Arons & Scharlemann, 1979](#); [Muslimov & Tsygan, 1992](#)). However, it was later demonstrated that the current density needed to support the twist of open magnetic field lines in the pulsar magnetosphere, j_{mag} , could not be sustained by a unidirectional plasma flow ([Timokhin, 2006](#)), in general requiring a flow with counter-streaming particles. In such flows, charge starvation, and thus the accelerating electric field, appears when there is a difference between the GJ current density, $j_{\text{GJ}} = \rho_{\text{GJ}}c$, and the current density required by the magnetosphere, $j_{\text{mag}} - j_{\text{GJ}}$; thus, the accelerating electric field is “current-density-driven”, rather than “charge-density-driven” ([Beloborodov, 2008](#); [Timokhin, 2010](#); [Timokhin & Arons, 2013](#)). Self-consistent particle-in-cell (PIC) plasma simulations ([Timokhin, 2010](#); [Timokhin & Arons, 2013](#)), which took into account particle ac-

celeration, photon emission, and pair production, demonstrated that, regardless of the ability of particles to escape from the NS surface, pair discharges always operate in a cyclic fashion. These findings are reminiscent of the model of [Ruderman & Sutherland \(1975\)](#) (RS) with no particle outflow from the NS surface. The bursts of pair formation are followed by a quiet outflow of plasma where the electric field is screened; once a substantial amount of the produced plasma leaves the polar cap, the electric field is restored, and the cycle restarts.

For most pulsars, the radio emission is generated in the inner magnetosphere not far from their polar caps (e.g., [Lyne et al., 2004](#); [Philippov & Kramer, 2022](#)). The observed brightness temperature of the radio emission is incredibly high, $10^{25} - 10^{30}$ K ([Lorimer & Kramer, 2005](#)), indicating a coherent nature of the radiation mechanism. Most of the existing models of pulsar radio emission assume the development of plasma instabilities in the pair plasma above the polar cap when the accelerating electric field is already screened ([Melrose et al., 2021b](#)). Recently, [Philippov et al. \(2020\)](#) proposed a novel radio emission mechanism driven by the screening of the accelerating electric field during pair discharge. Their 2D PIC simulations demonstrated the excitation of electromagnetic superluminal modes when the discharge occurs obliquely with respect to the background magnetic field, i.e., in this model, the radio emission is caused by the spatial inhomogeneity of the pair discharges. Several studies have confirmed the excitation of electromagnetic waves through the [Philippov et al. \(2020\)](#) mechanism using 2D local ([Cruz et al., 2021](#); [Benáček et al., 2024](#)) and global ([Bransgrove et al., 2023](#)) PIC simulations. However, the scale at which the discharge maintains spatial coherence across magnetic field lines has not yet been explored, partially because of the limited separation of plasma and global scales and the simplified physics of pair production in multi-dimensional simulations. Spatial decorrelation of the polar discharge across the field lines might seem plausible as a physics basis to explain some

aspects of the short-term nulling of the radio emission (e.g., [Ng et al., 2020](#)).

Observed radio emission shows variability over a wide range of timescales. This variability includes the microstructure of individual radio pulses at a level of ~ 100 microseconds and smaller, as well as the quasi-stable drift of individual sub-pulses seen in some pulsars (e.g., [Drake & Craft, 1968](#); [Janagal et al., 2023](#)). To explain these phenomena, RS proposed that pair discharges occur in a set of localized columns, which they called “sparks”, separated by plasma-depleted regions. The potential drop in the sparks gives rise to the perpendicular to the magnetic field component of the electric field, which causes sparks’ quasi-regular drift with respect to the NS. The spark model is still being used for modeling drifting subpulses (e.g. [Gil & Sendyk, 2000](#); [Gil et al., 2003](#); [Basu et al., 2020](#)). The intermittency of the cyclic discharge, as revealed by simulations, is a natural candidate for explaining the microstructure. On the other hand, the existence of long-lived isolated sparks in intermittent discharges is questionable.

In this Letter, we aim to explore the existence and sustainability of the spatial de-correlations in the pair discharges for the polar cap models with no ([Ruderman & Sutherland, 1975](#)) as well as free ([Arons & Scharlemann, 1979](#)) escape of particles from the NS surface. Using 2D and 3D PIC simulations, we study how pair discharges fill the polar cap and the characteristic discharge scales perpendicular to the magnetic field. This chapter is organized as follows. In Section [3.3](#), we describe the physical model employed in our simulations of polar cap discharges. In Section [3.4](#), we present the main results of our simulations. We demonstrate that the discharge desynchronizes on transverse scales comparable to the length of the gap along the field line and conclude that localized sparks can not exist in polar caps of pulsars with force-free magnetospheres. Section [3.5](#) discusses the applications of our results to observed pulsars.

3.3 Overview of the physical model

3.3.1 Dynamics of the electromagnetic fields

Global solutions for the pulsar magnetospheres in the force-free (FFE) approximation (e.g., [Contopoulos et al., 1999](#); [Spitkovsky, 2006](#); [Timokhin, 2006](#); [Kalapotharakos et al., 2012](#)) demonstrate that the open field lines develop a twist that is self-regulated by conditions at the light cylinder. The global toroidal magnetic field, $\mathbf{B}_\varphi$, sets the distribution of the field-aligned magnetospheric current, $\mathbf{j}_{\text{mag}} \equiv (c/4\pi)\nabla \times \mathbf{B}_\varphi$, in the polar cap zone. We decompose the electromagnetic fields into a sum of the FFE fields, $\mathbf{B}_{\text{FFE}} = \mathbf{B}_0 + \mathbf{B}_\varphi$, $\mathbf{E}_{\text{FFE}} = -\boldsymbol{\Omega} \times \mathbf{r} \times \mathbf{B}_0/c$, where $\mathbf{B}_0$ is the background (e.g., dipolar) poloidal magnetic field, and corrections, $\delta\mathbf{E}$ and $\delta\mathbf{B}$, created in the process of adjustment of the electric current in the gap, $\mathbf{j}$, to the global magnetospheric current, $\mathbf{j}_{\text{mag}}$. By substituting the field decomposition into Maxwell equations, we obtain the evolution equations for $\delta\mathbf{E}$ and $\delta\mathbf{B}$:

$$\frac{\partial}{\partial t}\delta\mathbf{E} = c\nabla \times \delta\mathbf{B} - 4\pi(\mathbf{j} - \mathbf{j}_{\text{mag}}), \quad (3.1)$$

$$\frac{\partial}{\partial t}\delta\mathbf{B} = -c\nabla \times \delta\mathbf{E}. \quad (3.2)$$

The electric field has also to obey Gauss's law:

$$\nabla \cdot \delta\mathbf{E} = 4\pi(\rho - \rho_{\text{GJ}}), \quad (3.3)$$

where ρ is the charge density, and $\rho_{\text{GJ}} = \nabla \cdot \mathbf{E}_{\text{FFE}}/4\pi = -\boldsymbol{\Omega} \cdot \mathbf{B}/2\pi c$ is the GJ charge density.

Unless mentioned otherwise, we consider the GJ density to be negative.

We assume that the distribution of magnetospheric current, j_{mag} , is constant in our simulations. This is well justified in actively pair-producing pulsars, in which case the global toroidal field, B_φ , evolves on timescales comparable to the light crossing time of the light cylinder, R_{LC}/c , which is much longer than the characteristic time in the polar cap zone, R_{PC}/c . Equations (3.1) and (3.2) have two special stationary, $\partial/\partial t = 0$, solutions:

1. There is abundant plasma everywhere to support the magnetospheric current, $j = j_{\text{mag}}$, which leads to a fully FFE solution, $\delta\mathbf{E}, \delta\mathbf{B} = 0$.
2. A field line has no plasma loading and $j = 0$ in the gap zone. Then, the full twist of the field lines becomes zero, $\delta B = -B_\varphi$.

In the previously studied 1D approximation, the set of Eqns. (3.1-3.2) reduces to the $\partial(\delta E_{\parallel})/\partial t = -4\pi(j_{\parallel} - j_{\text{mag}})$, since only components of the field and current parallel to the background field can be evolved in 1D (Beloborodov, 2008; Timokhin, 2010; Timokhin & Arons, 2013). By construction, these solutions describe deviations from the FFE state, i.e., solution 1. Our multidimensional simulations solve the full system of Eqns. (3.1-3.2) together with the constraint Eqn. (3.3). In contrast to the 1D simulations, if the pair creation process is unable to sustain the global magnetospheric current, the total twist can become zero, and the system reaches state 2.

The limitations of numerical methods do not allow the study of processes in pulsar polar caps for realistic values of physical parameters except in 1D. One must rely on a scaled model with sometimes exaggerated parameters to make such a study possible in 2D and 3D. Despite the simplifications, our numerical setup preserves all the important properties of the real polar cap at the base of the FFE magnetosphere, as described below.

3.3.2 Magnetospheric current distribution

The work by [Gralla et al. \(2016, 2017\)](#) provided fits of the spatial distribution of the magnetospheric current over the polar cap of oblique rotators. They found that for misaligned rotators, a significant fraction of the polar cap is capable of supporting active pair production, i.e., possesses super-GJ current, $j_{\text{mag}}/j_{\text{GJ}} > 1$, where again $j_{\text{GJ}} = \rho_{\text{GJ}}c$, or the volume return current, $j_{\text{mag}}/j_{\text{GJ}} < 0$ ([Timokhin & Arons, 2013](#)).

For a 2D simulation performed in Cartesian coordinates, we choose the profile of the magnetospheric current to be qualitatively similar to the solution for the aligned rotator:

$$j_{\text{mag}}^{2\text{D}}(x) = 2 - Ax^4 + Bx^6 + Cx^{40}, \quad (3.4)$$

here $x = r_{\perp}/R_{\text{PC}}$ is the coordinate transverse to the background magnetic field, B_0 , which is chosen to be uniform. Coefficients A , B , C are chosen to satisfy the conditions of total zero current and smooth transition to the zone with no current:

$$\int_0^{R_{\text{PC}}} j_{\text{mag}}^{2\text{D}} dx = 0, \quad j_{\text{mag}}^{2\text{D}}(1) = 0, \quad \left. \frac{d}{dx} j_{\text{mag}}^{2\text{D}} \right|_{x=1} = 0. \quad (3.5)$$

This choice results in approximately half of the polar cap field lines having the volume return magnetospheric current, $j_{\text{mag}} < 0$, and half of the field lines carry a super-GJ current, $j_{\text{mag}} > j_{\text{GJ}}$ ¹.

The current profile is visualized in Fig. 3.1a. Our choice excludes the extremely narrow zone of the intense return current at the boundary of the open field line region, the “separatrix” current,

¹Our current density distribution has a higher maximum value and results in larger portions of the polar cap occupied by the super-GJ current density compared to a real axisymmetric pulsar. However, qualitatively, it preserves the properties of an actual current density distribution.

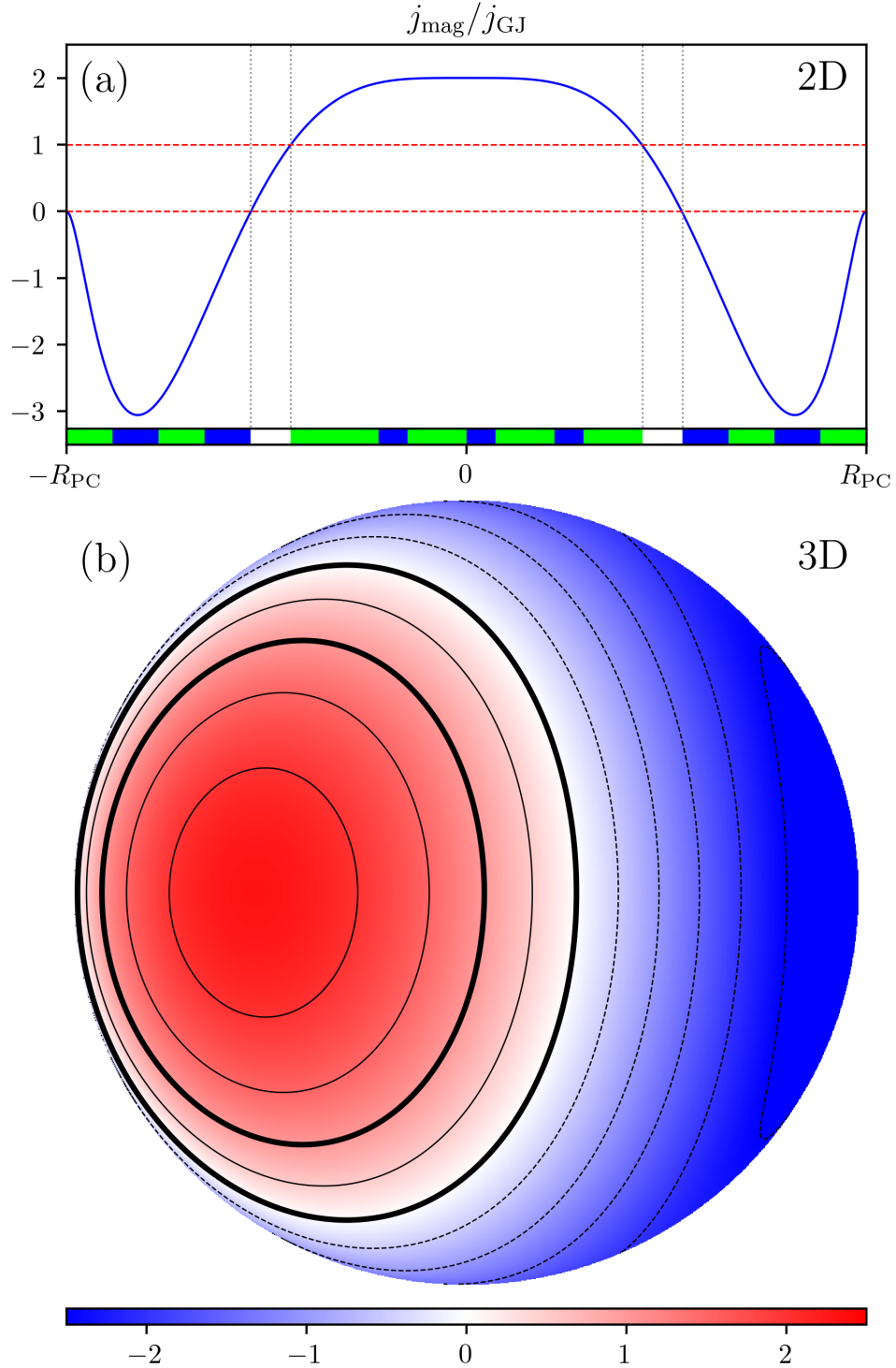


Figure 3.1: Distribution of the magnetospheric current across the polar cap used in our (a) 2D and (b) 3D simulations, which mimic the currents in the magnetospheres of an aligned and 60° -inclined rotators. The net current across the polar cap is zero, with outgoing and return current regions each covering roughly half of the polar cap. In panel (b), the thick black lines represent the boundaries of super GJ and return current regions. The colored inset in panel (a) marks the boundaries of initial patches (see Sec. 3.3.5).

which is present in magnetospheric solutions for any obliquity. Studying the discharge in the separatrix requires fully global simulations (e.g., Philippov et al., 2015; Hu & Beloborodov, 2022).

In 3D simulations, the approximation for the distribution of the magnetospheric current over the polar cap is taken from Gralla et al. (2016):

$$\begin{aligned} \frac{j_{\text{mag}}^{\text{3D}}}{\rho_{\text{GJ}} C}(\theta, \phi) \approx & \frac{1}{(1 - \Omega_Z/\Omega)} [J_0(\arcsin(\theta/\sqrt{\alpha_0})) - \\ & - J_1(\arcsin(\theta/\sqrt{\alpha_0})) \tan i \cos \phi]. \end{aligned} \quad (3.6)$$

Here, $J_{0,1}$ are Bessel functions of the zeroth and first order, $\theta = r_{\perp}/R_{\star}$ gives the distance from the center of the polar cap, and ϕ is an azimuthal angle. The constant parameters i , Ω_Z , α_0 describe the global magnetosphere, namely $i = 60^\circ$ is an inclination angle, $\Omega_Z \approx (2C/5)\Omega$ is the angular velocity of the frame dragging, and $C = 2M/R_{\star} \approx 0.5$ is stellar compactness, and α_0 defines the size of the polar cap for the dipolar magnetic field. The analytical formula (3.6) has a non-zero total current through the polar cap, requiring a separatrix current. In our numerical setup, the mean current $\langle j_{\text{mag}} \rangle_{\text{PC}}$ is subtracted from (3.6) to enforce the zero total volumetric current. The resulting distribution of the magnetospheric current is shown in Fig.3.1b.

3.3.3 QED pair production

1D studies by Timokhin (2010); Timokhin & Arons (2013) showed that pair production in polar zones operates intermittently, i.e., episodes with strong electric field component along the magnetic field and intense particle acceleration are followed by plasma-filled states with fully screened electric fields. During the active phase, electrons and positrons are accelerated into two

oppositely directed beams with high Lorentz factors $\gamma_b \sim 10^7$ for real pulsars (Timokhin, 2010).

During propagation along magnetic field lines with curvature radius ρ , high-energy particles emit γ -rays of energy ε at a rate of

$$\frac{dN_{\text{ph}}}{dt d\varepsilon} = \frac{1}{\sqrt{3}\pi} \frac{e^2}{\hbar^2 c} \frac{1}{\gamma_b^2} \int_{\frac{\varepsilon}{\varepsilon_{\text{ph}}^*}}^{\infty} K_{5/3}(x) dx, \quad (3.7)$$

where $K_{5/3}$ is Macdonald function, and $\varepsilon_{\text{ph}}^*$ is characteristic energy of a curvature photon

$$\varepsilon_{\text{ph}}^* = \frac{3}{2} \hbar \frac{c}{\rho} \gamma_b^3. \quad (3.8)$$

Photons are emitted tangentially to the magnetic field lines, and the angle between the photon propagation and magnetic field direction, ψ , increases. In our simulations, the angle increases linearly with distance from the emission point, z : $\sin \psi = \int dz / \rho(z)$. We do not include the displacement of photons from one field line to another during their propagation. This is justified because the transverse propagation distance, $\sim l_{\text{gap}}(l_{\text{gap}}/\rho)$, where l_{gap} is the longitudinal size of the gap, is significantly smaller compared to the polar cap radius, R_{pc} , in pulsars with $l_{\text{gap}} \lesssim R_{\text{pc}}$ considered in this work.

High-energy photons propagating at a non-zero angle with respect to the strong magnetic field eventually decay into electron-positron pairs. The cross-section of the photon absorption is given by (Erber, 1966):

$$d\sigma = 0.23 \frac{B}{B_q} \sin \psi \frac{\alpha_F}{\lambda_c} \exp\left(-\frac{8}{3\chi}\right) dz, \quad (3.9)$$

where $\chi = (B/B_q) \tilde{\varepsilon}_{\text{ph}} \sin \psi$, $B_q = m_e^2 c^3 / e \hbar \approx 4.41 \times 10^{13} \text{G}$ is the critical magnetic field, and $\tilde{\varepsilon}_{\text{ph}} = \varepsilon_{\text{ph}} / m_e c^2$ corresponds to the photon energy measured in units of $m_e c^2$. After rapid

radiation of the momentum component perpendicular to the magnetic field, secondary pairs move with the following 4-velocity along the magnetic field,

$$u_{\parallel} = \frac{|\cos \psi_a|(\tilde{\varepsilon}_{\text{ph}}^2 - 4)^{1/2}}{(\tilde{\varepsilon}_{\text{ph}}^2 \sin^2 \psi_a + 4 \cos^2 \psi_a)^{1/2}} \sim \frac{1}{\sin \psi_a} \sim 10^2 - 10^3, \quad (3.10)$$

where ψ_a is the angle of the photon propagation direction at the point of absorption. This process generally results in large pair multiplicity, $\mathcal{M} = n_{\pm}/n_{\text{GJ}} \gg 1$ (Timokhin & Harding, 2019).

Realistic scales of particle and photon energies have so far been achieved only in one-dimensional simulations. In 2D and 3D, we adopt the following re-scaling procedure. The maximum Lorentz factor of particles accelerated in a vacuum electric field on a scale of the transverse size of the polar cap $\gamma_{\text{PC}} = 0.5(R_{\text{PC}}/d_e^{\text{GJ}})^2$, where $d_e^{\text{GJ}} = c/\sqrt{4\pi\rho_{\text{GJ}}e/m_e}$ is the cold skin depth defined for the density of the GJ plasma. Here, we used an electric field linearly increasing with distance, appropriate for super-GJ and volume return currents (Timokhin & Arons 2013, also see Appendix B.1). The cooling efficiency due to curvature radiation is parameterized by the radiation reaction limit, γ_{rad} . It is defined as the Lorentz factor corresponding to the balance between the accelerating electric force and the radiation reaction force, $eE_{\text{PC}} = (2/3)e^2\gamma_{\text{rad}}^4/\rho^2$. Finally, the energies of the curvature photons are expressed as $\tilde{\varepsilon}_{\text{ph}} = (\gamma/\gamma_{\text{emit}})^3$, where γ_{emit} parameterizes the Lorentz factor of a particle emitting a curvature photon with an energy of an electron rest mass, $(3/2)\hbar(c/\rho)\gamma_{\text{emit}}^3 = m_e c^2$. Together with the curvature radius of field lines, ρ , and the strength of the magnetic field B/B_q , these energy scales fully determine the problem. For Vela with the rotational period $P = 0.089\text{s}$ and surface magnetic field $B_{\star} \approx 2 \times 10^{12}\text{G}$, the characteristic Lorentz factors $\gamma_{\text{emit}} = 8 \times 10^5$, $\gamma_{\text{rad}} = 1.8 \times 10^8$, $\gamma_{\text{PC}} = 6.5 \times 10^9$, resulting in the hierarchy of energy scales as $\gamma_{\text{emit}} \ll \gamma_{\text{rad}} \ll \gamma_{\text{PC}}$. Our simulations downscale all these num-

bers (see Table 3.1) but preserve the above hierarchy. We provide additional details of numerical implementation and QED-related rescaling relations in Appendix B.1.

Table 3.1: Summary of Simulation Parameters for Discharge Simulations.

Simulation	$l_{\text{gap}}/R_{\text{PC}}$	$R_{\text{PC}}/d_{\text{e}}^{\text{GJ}}$	γ_{PC}	γ_{rad}	γ_{emit}	ρ/R_{PC}	B_0/B_q
SCLFsmall	$l_{\text{gap}} \ll R_{\text{PC}}$	12000	7.2×10^7	6.7×10^5	10^4	13	2.0
SCLFcurv	$l_{\text{gap}} \ll R_{\text{PC}}$	12000	7.2×10^7	6.7×10^5	10^4	13	2.0
SCLFtail	$l_{\text{gap}} \ll R_{\text{PC}}$	12000	7.2×10^7	1.35×10^5	4×10^3	13	2.0
SCLFlarge	$l_{\text{gap}} \lesssim R_{\text{PC}}$	8000	3.2×10^7	8×10^5	3×10^4	20	0.1
Rsmall	$l_{\text{gap}} \ll R_{\text{PC}}$	8000	3.2×10^7	8×10^5	3×10^4	20	1.0
Rlarge	$l_{\text{gap}} \lesssim R_{\text{PC}}$	8000	3.2×10^7	8×10^5	7×10^4	20	0.1
SCLF3D	$l_{\text{gap}} \lesssim R_{\text{PC}}$	1000	5×10^5	10^4	10^3	5	0.2

3.3.4 Atmosphere

The surface of the NS is covered by a thin electron-ion atmosphere (e.g., [Haensel et al., 2006](#)). This atmosphere provides a reservoir of charged particles that an unscreened electric field can pull into the magnetosphere. To model the space-charge limited flow, hereafter SCLF, we place a thermal plasma just above the inner boundary of the simulation domain, where we set conducting boundary conditions for electromagnetic fields and outflow for plasma particles. The thermal plasma has a Boltzmann spatial distribution $n = n_{\text{peak}} \exp(-z/h)$. Here, z is a coordinate along the background magnetic field, B_0 , the peak density $n_{\text{peak}} \approx 10n_{\text{GJ}}$, and $h = kT/(m_e g) \approx 10d_{\text{e}}^{\text{GJ}}$ is a scale-height. The Boltzmann distribution is supported by adding a constant gravitational force $-m_e g$ to the particle equation of motion². At each time step, the injector replenishes escaping particles within one scale height from the stellar surface to maintain

²We turn this force off inside the magnetosphere, since in real pulsars, the Lorentz force vastly exceeds the gravitational force.

the Boltzmann distribution. The temperature T is selected to be low to prevent the situation when the quasi-neutral thermal outflow carries the magnetospheric current from the tail of the Boltzmann distribution (Beloborodov & Thompson, 2007). We verify that our setup produces correct 1D SCLF solutions for both sub- and super-GJ currents in Appendix B.1. Additionally, we perform simulations adopting the RS model, where the atmospheric supply of plasma is absent. In this case, plasma is generated solely through pair production in discharge events.

3.3.5 Initial plasma state

Above the atmosphere, we initialize plasma with multiplicity $\mathcal{M} \sim \text{few}$ to provide the magnetospheric current $j = j_{\text{mag}}$ by sub-relativistic counter-streaming. This is done to avoid finding the configuration of vacuum electric fields in multi-dimensional settings and allows to start with initial conditions $\delta\mathbf{E} = 0$. In the 1D setup, we verified that the late time evolution of the discharge is insensitive to starting with vacuum or plasma-filled conditions. To satisfy the constraint given by Eqn. (3.3), we add an uncompensated charge, $\rho = \rho_{\text{GJ}}$. As pointed out in Timokhin & Arons (2013), gap formation in super-GJ current is sensitive to the sign of the spatial gradient of the GJ density. In particular, the gap appears close to the stellar surface if the gradient of $|\rho_{\text{GJ}}|$ is positive, as dictated by the general-relativistic correction caused by frame-dragging (Beskin, 1990; Muslimov & Tsygan, 1992; Philippov et al., 2015). Thus, we adopt the model of decreasing GJ charge density in the simulation box, $\rho_{\text{GJ}} = \rho_{\text{GJ}}^0(1 + 0.8z/L_z)$, where L_z is the size of the simulation box along the magnetic field³.

Simulations initialized with a uniform distribution of plasma across magnetic field lines

³Realistic gradient caused by the frame-dragging is significantly smaller: $|\rho_{\text{GJ}}| = (\Omega - \Omega_Z)B_0/(2\pi c) = |\rho_{\text{GJ}}^0|(1 + Az/L_z)$, where $|\rho_{\text{GJ}}^0| = \Omega B_0(1 - 2C/5)/2\pi c$ and $A \approx (6C/5)(L_z/R_*)/(1 - 2C/5)$. In 1D simulations performed for realistic parameters, we verified that the location of the gap is similar to more realistic values of the gradient, e.g., $A \approx 0.07$ corresponding to Vela pulsar.

result in a coherent discharge across the whole polar cap. For the duration of our simulations, this idealized coherence persists. In this paper, we aim to study the “natural” state of pair discharges when discharges on all field lines do not necessarily start simultaneously, so we introduce initial inhomogeneities in plasma and let the system evolve. To desynchronize the discharges on neighboring field lines, we split the polar cap into the domains or patches. Initially, each patch is populated with a high-density plasma, $n \gg n_{\text{GJ}}$, which can be thought of as the remaining plasma from the previous discharge episode. To maintain these high-density regions, we slowly inject $e^- - e^+$ plasma into the patch to maintain $\delta \mathbf{E} \approx 0$ and the absence of particle acceleration. At some point, we stop the plasma injection, and the discharge restarts when most of the plasma falls into the atmosphere or escapes the box as a result of counter-streaming demanded by the magnetospheric current. We stop supplying the initial plasma at different times on neighboring patches, which leads to the desynchronization of phases of discharges. Following self-consistent evolution after all patches are cleared of the initial plasma allows us to study the phase coherence of discharges across the polar cap. The initial partition of the polar cap into patches is shown by a colored inset in Fig. 3.1a.

3.3.6 Numerical details

We perform simulations of pair production discharges using the `Tristan-v2`, multi-species radiative PIC code (Hakobyan et al., 2023). We present results of 2D and 3D simulations with resolution $\Delta x = d_e^{\text{GJ}}$, where Δx is the cell size. In 2D, we verified that our results are unchanged for $\Delta x = d_e^{\text{GJ}}/2$. To ensure that the hot skin depth, $d_e^{\text{hot}} = c/\sqrt{4\pi n e^2 \langle 1/\gamma^3 \rangle / m_e}$, is always resolved in the simulation domain, we limit the multiplicity of the produced electron-

positron pairs at $\mathcal{M}_{\text{lim}} = 50$. Our system size is $R_{\text{PC}}/d_e^{\text{GJ}} = 8 \times 10^3 \div 1.2 \times 10^4$ for 2D, resulting in the Lorentz factors corresponding to the full vacuum potential drop across the polar cap, $\gamma_{\text{PC}} = 0.5(R_{\text{PC}}/d_e)^2 \approx 5 \times 10^7$. In 3D, the size of the polar cap is $R_{\text{PC}} = 10^3 d_e^{\text{GJ}}$, and correspondingly $\gamma_{\text{PC}} \sim 10^6$. The simulation domain size is $2R_{\text{PC}}$ along the stellar surface and $2R_{\text{PC}}$ along the field lines. We study both the SCLF and RS models with varying parameters, corresponding to small, $l_{\text{gap}} \ll R_{\text{PC}}$, and large, $l_{\text{gap}} \lesssim R_{\text{PC}}$, gap sizes in 2D. Additionally, we perform a simulation in the SCLF model of a marginally small gap, corresponding to the current distribution in the 60° -inclined rotator. The summary of the parameters of different simulations is given in Table 3.1. We employ 8 filter passes on the deposited currents to reduce the particle noise. The fields are evolved using modified stencils (Blinne et al., 2018), which significantly improve the isotropy of the numerical dispersion relation of the electromagnetic waves. This method also leads to a substantial reduction of the electromagnetic noise in our multi-dimensional simulations compared to the evolution employing the standard Yee solver.

3.4 Results

1D simulations in the SCLF regime found qualitatively different regimes of the discharge operation, dependent on the value of the magnetospheric current. (Timokhin & Arons, 2013).

1. Super-GJ current, $j_{\text{mag}}/j_{\text{GJ}} > 1$. The gap opens close to the stellar surface. A beam of high-energy electrons is accelerated outwards by the electric field in the gap and produces photons and secondary pairs that escape into the magnetosphere. Positrons in the gap accelerate towards the star and produce pairs bombarding the surface.
2. Sub-GJ current, $0 < j_{\text{mag}}/j_{\text{GJ}} < 1$. Here, a sub-relativistic beam of electrons of density

$\approx n_{\text{GJ}}$ is extracted from the atmosphere, supplying enough charge density and current to screen the electric field. No pair production happens in this zone.

3. Return current, $j_{\text{mag}}/j_{\text{GJ}} < 0$. The gap opens at high altitudes, close to the upper boundary of the simulation domain⁴. Since the gap is located at a large distance from the star, only the beam of particles propagating toward the star leads to efficient pair production.

In the RS model, [Timokhin \(2010\)](#) demonstrated that the magnetic field lines with the $j_{\text{mag}}/j_{\text{GJ}} > 1$ and $j_{\text{mag}}/j_{\text{GJ}} < 0$ operate overall similarly to the SCLF case. On the other hand, field lines with sub-GJ current $0 < j_{\text{mag}}/j_{\text{GJ}} < 1$ become capable of particle acceleration and subsequent pair production. Our multi-dimensional simulations support these basic qualitative conclusions but also reveal novel dynamics across magnetic field lines.

3.4.1 Space-Charge-Limited Flow

In this Section, we present the results of simulations of a discharge in the 2D polar cap with the supply of plasma from the stellar atmosphere (see simulations `SCLFsmall`, `SCLFcurv`, `SCLFtail`, `SCLFlarge` in Table 3.1). In all of these simulations, we start with three patches loaded with low-density plasma, allowing quick opening of the gap on these field lines. Two of the patches are located in the zone of return current, and one is in the zone of the super-GJ current.

⁴It is likely that this gap initially opens at the null-surface, where $\rho_{\text{GJ}} = 0$, see evidence in global simulations in, e.g., [Bransgrove et al. \(2023\)](#).

3.4.1.1 Small Gap, Constant Curvature of Field Lines

In polar caps of real pulsars, the magnetic field can deviate from a dipolar configuration due to the presence of local multipolar components, resulting in magnetic field configurations that do not necessarily possess axial symmetry around the center of the polar cap. Our configurations with constant curvature radius represent such polar caps. The results of the simulation `SCLFsmall` with a small gap, $l_{\text{gap}} \ll R_{\text{pc}}$ ⁵, and constant field line curvature, ρ , are shown in Fig. 3.2. In the super-GJ zone, the gap forms close to the stellar surface. Initially, electric potential along the gap is limited by the transverse size of the patch, resulting in slow acceleration and low energies of electrons extracted from the atmosphere. The mean free path of photons emitted by these electrons is large, and pair production happens beyond the gap, i.e., in a region with $\delta E \approx 0$. Since produced pairs have a non-zero momentum directed outward, none of them can return into the gap zone to screen the electric field. At the same time, the region beyond the gap is filled with high-density secondary plasma, preventing the gap from further expansion. Under these conditions, the gap becomes quasi-stationary: its size does not change in time. We find that the gap operates in a quasi-stationary regime when the size of the gap along the field lines is large, $l_{\text{gap}} \gg w$, where w is the transverse width of the patch. Quasi-stationary separated narrow gaps are seen in the active patches in the zone with super-GJ current at early times, as shown in the left column of Fig. 3.2.

When a new patch between the two already active patches with quasi-stationary gaps is cleared of plasma, the unscreened electric field in the plasma-starved zone communicates across

⁵Sizes of gaps are quoted for the zone with super-GJ current, in the state when all patches are cleared of the initial plasma. Parameters of the simulation in Table 3.1 are calibrated using 1D simulations for corresponding values of j_{mag} .

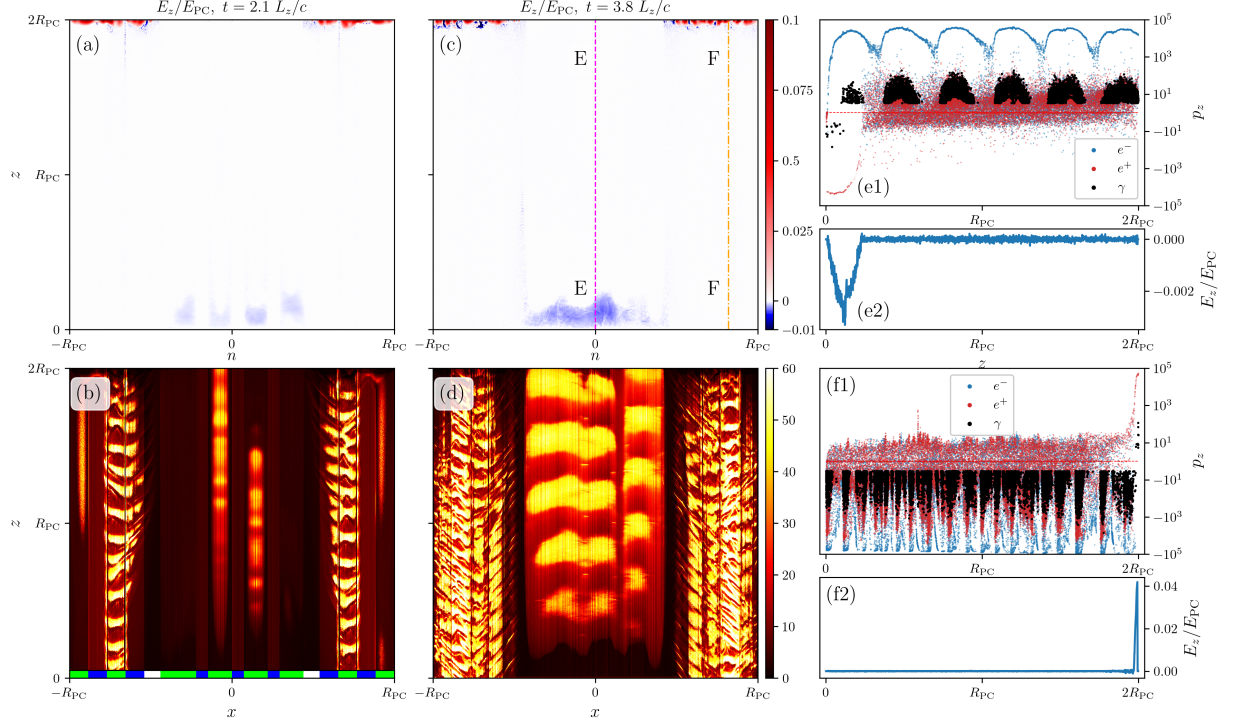


Figure 3.2: Evolution of the discharge in the 2D polar cap with a small gap, $l_{\text{gap}} \ll R_{\text{PC}}$, and free particle escape from the stellar surface (SCLF). The first column depicts (a) the electric field and (b) plasma density during the initial phase of the evolution, where patches occupy a large area of the polar cap. At this time, the discharge along the super-GJ field lines behaves in an almost steady-state manner. The second column illustrates a later evolutionary stage when all patches are cleared of the initial plasma, leading to a cyclic discharge. The asymmetry of clouds of produced plasma suggests the decorrelation of discharges on neighboring field lines. Lines “E” and “F” indicate the positions of 1D cuts of the $z - p_z$ phase space and electric field for the region of the super-GJ current (e1,e2) and (f1,f2) - for return current. In the phase space plots, blue, red, and black dots represent the position and momentum of electrons, positrons, and pair-producing photons, correspondingly.

the field lines. This leads to a significant increase in the accelerating potential, proportional to the increase of the transverse size of the gap. Electrons, accelerated to higher energies, emit more energetic photons that are absorbed within the gap, leading to the cyclic screening of the electric field. The example of the merging of a few quasi-stationary gaps can be seen in the central column of Fig. 3.2 (subplots c-d). These snapshots correspond to the later time when all patches are cleared of initial plasma, and the whole polar cap is available for discharge.

In a cyclic discharge, the unscreened electric field on a particular field line exists during $\tau_{\text{active}} \approx l_{\text{gap}}/c$. Thus, the transverse size of the region affected by the unscreened electric field on the field line cannot exceed $c\tau_{\text{active}} = l_{\text{gap}}$. Hence, if two field lines are separated by a distance exceeding $2l_{\text{gap}}$, their electric fields can not communicate. This condition sets the limitation to the transverse coherence scale of a discharge, $l_{\perp} \sim l_{\text{gap}}$. The transverse variation of shapes of the front of the unscreened electric field and clouds of secondary plasma seen in Fig. 3.2d present a hint of the desynchronization of discharges at different field lines. We note that the size of the cyclic gap on field lines carrying the super-GJ current is larger compared to the one expected for a vacuum electric field. Here, the electric field extracts electrons from the stellar atmosphere and positrons from the plasma cloud produced in the previous discharge episode into the gap zone. The presence of both electrons and positrons results in a smaller deviation of the plasma charge density from the GJ value, leading to smaller electric field in the gap, $\sim 0.1|E_0|$, where $E_0 = 4\pi\rho_{\text{GJ}}l_{\text{gap}}(|j_{\text{mag}}/j_{\text{GJ}}| - 1)$ corresponds to the case of only electrons carrying the current (see the 1D slice of the electric field in subplot (e2) of Fig. 3.2).

On field lines carrying the return current, a strong electric field appears at the upper boundary of the simulation box, where a lack of electrons develops. The current in the gap zone is carried by positrons moving into the magnetosphere. This leads to strong accelerating electric fields reaching $E_z \sim E_0 = 4\pi|\rho_{\text{GJ}}|l_{\text{gap}}(|j_{\text{mag}}/j_{\text{GJ}}| + 1)$. The electric field in the return current zone⁶ is shown in the subplot (f2) of Fig. 3.2. Photons emitted by the high-energy positron beam decay into pairs moving into the magnetosphere. The unscreened electric field reverses some of the electrons towards the star and accelerates them to high energies, igniting the discharge moving inwards (see a 1D cut of the phase space $z - p_z$ in subplot (f1) of Fig. 3.2). Because

⁶Note that the color bar in the electric field panels (subplots a and b) has different lower and upper bounds.

of the much stronger electric fields, the gap size in the return current zone is smaller compared to the gap in the super-GJ zone. Since here $l_{\text{gap}} \ll w$, the discharge within the first open patch is immediately in an intermittent regime (subplot (b) of Fig. 3.2), for chosen parameters of the gap and widths of initial patches⁷. When all patches are cleared of initial plasma, the discharge proceeds intermittently on all field lines. The complex shape of the clouds of secondary plasma additionally points to the transverse desynchronization of the gaps at the scale of $l_{\perp} \sim l_{\text{gap}}$.

3.4.1.2 Small Gap, Strong Desynchronization

Both regular (e.g., [Manzali et al., 2007](#); [Rigoselli et al., 2022](#)) and millisecond ([Miller et al., 2019, 2021](#)) pulsars show thermal X-ray emission from their surfaces. The source of heating behind this radiation is thought to be the bombardment of the stellar surface by the ultra-relativistic beam of particles accelerated in a gap. These backward propagating beams can be seen in 1D cuts of phase spaces $z - p_z$, shown in subplots (f1)-(e1) of Fig. 3.2. For a stationary gap, the X-ray luminosity can be estimated as $L_X = 0.5L_0(\gamma_{\text{gap}}/\gamma_{\text{pc}})$, where L_0 is a spin-down luminosity and γ_{gap} is the maximum Lorentz factor of particles in the gap. It is long known that this estimate gives significant overheating of the stellar surface compared to the observations (e.g., [Beskin, 2018](#)). For the cyclic discharge, the surface heating is attenuated by a repetition rate, $f = \tau_{\text{active}}/\Delta\tau_{\text{rep}} \ll 1$, where τ_{active} is again the lifetime of the unscreened electric field, and $\Delta\tau_{\text{rep}}$ is the time between successive discharges ([Timokhin & Harding, 2015](#)).

We expect some long-lived oscillations of an electric field to remain inside the secondary plasma cloud after the main screening event in a gap ([Tolman et al., 2022](#)). These oscillations

⁷At the same time, the discharge in the outermost isolated patches initially occurs in a quasi-stationary regime due to the low available potential on field lines with small return currents. It remains stationary until the patch next to it becomes available for the discharge. After a neighboring patch opens, the electric fields of the two patches synchronize, and discharge continues intermittently.

can reverse a fraction of injected secondary particles and deliver them to the gap zone. Suppose the process of particle reversal persists for a long time, $\gg \tau_{\text{active}}$. In that case, the electric field behind the cloud of secondary plasma remains screened, and the repetition rate f is small.

In our multi-dimensional simulations, we observe $\Delta\tau_{\text{rep}} \approx \text{few} \times \tau_{\text{active}}$, which is not sufficient to explain the X-ray observations. We believe this is because of the low allowed plasma multiplicity in the clouds and, hence, low-density backflow – a conjecture that we will address in the forthcoming work. In order to decrease the repetition rate, we inject additional extended tails behind the escaping clouds of secondary plasma. We trace high-density clouds of curvature photons and inject additional pairs in cells where the photon cloud passed less than $0.3L_z/c$ ago. The parameters of the simulation are chosen such that the discharges are in the cyclic regime within each initial patch, which corresponds to an even smaller gap size compared to `SCLFsmall` and stronger de-synchronization. As shown in Fig. 3.3b, simulation `SCLFtail` based on this prescription indeed results in long tails and low repetition rates. While the size and the lifetime of initial patches are identical to ones in simulation `SCLFsmall`, the structure of discharges is substantially different (the density snapshot for `SCLFsmall` is reproduced in Fig. 3.3a for convenience). Discharges happen completely independently within each transverse zone $l_{\perp} \sim l_{\text{gap}}$ with no synchronization between each other. We coin the resulting structure as a “lava lamp” discharge. This result follows from our previous statement that synchronization of neighboring discharges proceeds during overlapping plasma-starved episodes. The probability for two initially desynchronized gaps to appear next to each other in neighboring patches is $\approx f$. Hence, for low repetition rates, $f \ll 1$, the probability of spontaneous synchronization is very low. At the same time, the discharge remains coherent within one patch of size $\sim l_{\text{gap}}$. In our simulations, this coherence is preserved by the weak variation of the repetition rate across

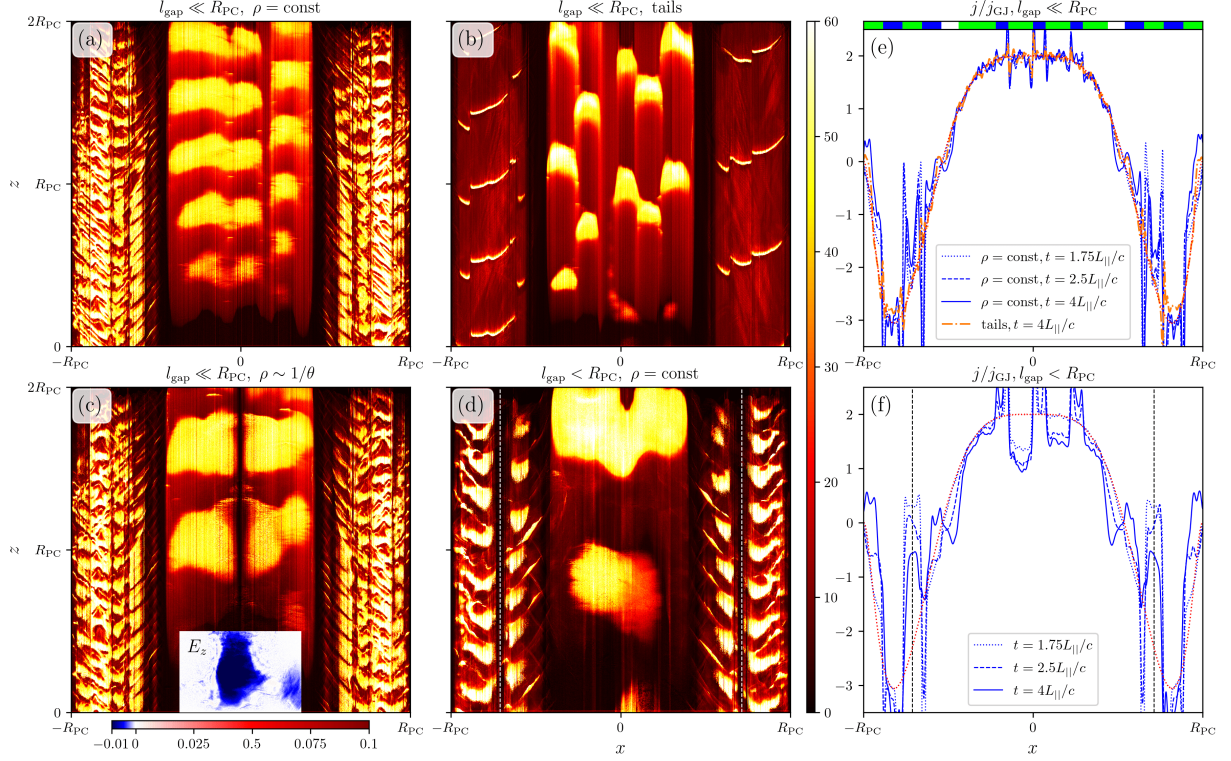


Figure 3.3: Comparison of plasma density distributions for various simulations in the SCLF regime. (a) Small gap $l_{\text{gap}} \ll R_{\text{PC}}$ with constant curvature (same as used in Fig. 3.2, simulation SCLF_{small} from Table 3.1). Panel (b) shows significant decorrelation of discharges on different field lines for a small gap and artificially reduced repetition rate (simulation SCLF_{tail}). Panel (c) presents the simulation SCLF_{curv} with the same parameters but with dipolar curvature, $\rho \sim 1/\theta$. The inset displays the field-aligned component of the electric field, highlighting the inclined fronts of field screening. Snapshot (d) depicts the case of a large gap, $l_{\text{gap}} \lesssim R_{\text{PC}}$ (SCLF_{large}). The right column shows plasma currents averaged along the field lines at different times for (e) small and (f) large gaps, indicating that smaller gaps more effectively sustain the twist of the magnetospheric field lines.

magnetic field lines.

3.4.1.3 Small Gap, Quasi-Dipolar Field

For the case of a commonly assumed star-centered dipole, the field line curvature diverges, $\rho \approx 4\sqrt{R_{\star}R_{\text{LC}}}/3\theta$, where θ is the polar angle, making pair production impossible near the magnetic axis. To study discharges in a dipolar geometry, in simulation SCLF_{curv} we impose the

curvature of the field lines in the form $\rho = \rho_0(R_{\text{PC}}/x)$, where values of ρ_0 and R_{PC} are given in Table 3.1, and x is the distance from the magnetic pole (center of the polar cap).

In Fig. 3.3c, we show a snapshot of the plasma density in the simulation. We find that close to the edge of the polar cap, in the return current zone, the discharge behavior is very similar to the simulation `SCLFsmall`. The field lines carrying the super-GJ current, however, have a substantially larger radius of curvature than the outer field lines. Comparison of snapshots with constant and dipolar curvatures, Fig. 3.3a,c, demonstrates that the gap in the super-GJ current zone is overall larger in the simulation with dipolar curvature than in the simulation with the constant curvature. This behavior can be easily understood considering the analytical scaling for the gap size (Timokhin & Harding, 2015):

$$l_{\text{gap}} \propto \xi^{-3/7} \rho^{2/7} P^{3/7} B^{-4/7}, \quad (3.11)$$

where $\xi \approx j_{\text{mag}}/j_{\text{GJ}}$ characterizes the strength of the electric field, P is the pulsar's rotational period, and B is the magnitude of the surface magnetic field. The variation of the gap size across the field lines in the super-GJ zone is determined by two factors: the strength of the accelerating electric field and the radius of curvature of the field lines. Since both $\xi \approx j_{\text{mag}}/j_{\text{GJ}}$ and the radius of curvature increases towards the pole, these two effects nearly cancel each other. Consequently, we find no substantial variation in the gap size everywhere except in the region very close to the magnetic axis where the field curvature diverges. Here, the gap size becomes infinite because the energy of curvature photons is very low (see Eq. 3.7), and a hole in the plasma density forms.

The de-synchronization of the discharges on neighboring field patches in the return current zone is practically identical to the simulation `SCLFsmall` (see Section 3.4.1.1). Synchronization

of the discharges in the super GJ-current zone is more active due to the overall larger size of the gap (also see Section 3.4.1.4). Here, synchronization of two patches on opposite sides of the magnetic axis occurs during the plasma-starved episodes because the unscreened electric field always exists in the central plasma density hole. Since the gap size varies substantially in a narrow region close to the magnetic axis, the front of the electric field screening is highly inclined with respect to the magnetic field, as shown in the inset of Fig. 3.3c. As pointed out by Philippov et al. (2020); Cruz et al. (2021), these inclined electric field screening fronts are plausible sources of superluminal ordinary modes, which can transform into the observed radio emission as they escape from the magnetosphere.

3.4.1.4 Large Gap

In this Section, we study a discharge with a gap size comparable to the size of the polar cap $l_{\text{gap}} \sim R_{\text{pc}}$, corresponding to less energetic pulsars. In Fig. 3.3d, we show a snapshot of the plasma density distribution in the simulation SCLF_{large}. At early times, a few isolated narrow patches are available for discharge in the region with the super-GJ current. In each of the patches, a quasi-stationary gap forms as described in Section 3.4.1.1. After all patches are cleared of the initial plasma, the discharges on neighboring field lines synchronize across most of the super-GJ zone of size $l_{\perp} \sim l_{\text{gap}}$. The resulting coherent discharge proceeds intermittently.

Similar conclusions apply to the discharge in the return current region, where we observe a pair of coherent discharges in each zone, see Fig. 3.3d. They are separated by a plasma density hole (white dashed lines in Fig. 3.3d mark its center), which has a peculiar origin. Here, if the patch width is too small to establish a high enough electric potential to trigger pair production,

the whole field line can become clear of pairs because of the counter-streaming demanded by the magnetospheric current. Simultaneously, an electric field extracts massive ions from the stellar atmosphere, which are not able to emit high-energy photons. In our simplified setup that does not take the transverse motion of photons into account, zones that once become clear of pairs can no longer ignite a discharge. These field lines possess an unscreened electric field that allows transverse communication of neighboring patches.

3.4.2 Ruderman-Sutherland Gap

The SCLF model ([Arons & Scharlemann, 1979](#)) corresponding to the free escape of particles from the stellar atmosphere seems more physically motivated; however, the RS model with no plasma supply from the star is still widely used to explain various observations (e.g., [Wen et al., 2020](#); [Primak et al., 2022](#); [Janagal et al., 2023](#)). In this Section, we consider discharges in the RS model of a polar cap. The profile of the magnetospheric current and fragmentation of the polar cap into patches is identical to the simulations in the SCLF model, as described in Section [3.4.1.1](#). The RS model is typically applied to a pulsar with counter-aligned magnetic moment and angular velocity vectors, the “anti-pulsar”. In this geometry, the GJ density is positive. In comparison with SCLF simulations, we change the sign of both the GJ density (see Section [3.3.5](#)) and the magnetospheric current.

In Fig. [3.4](#), we show snapshots of the accelerating electric field and plasma density for the simulations with small (simulation `RSsmall`) and large (`RSlarge`) gaps. The evolution of discharges in the RS model is very similar to that of the SCLF model: the discharges operate in an intermittent regime, and the directions of discharges are identical to those in the SCLF polar

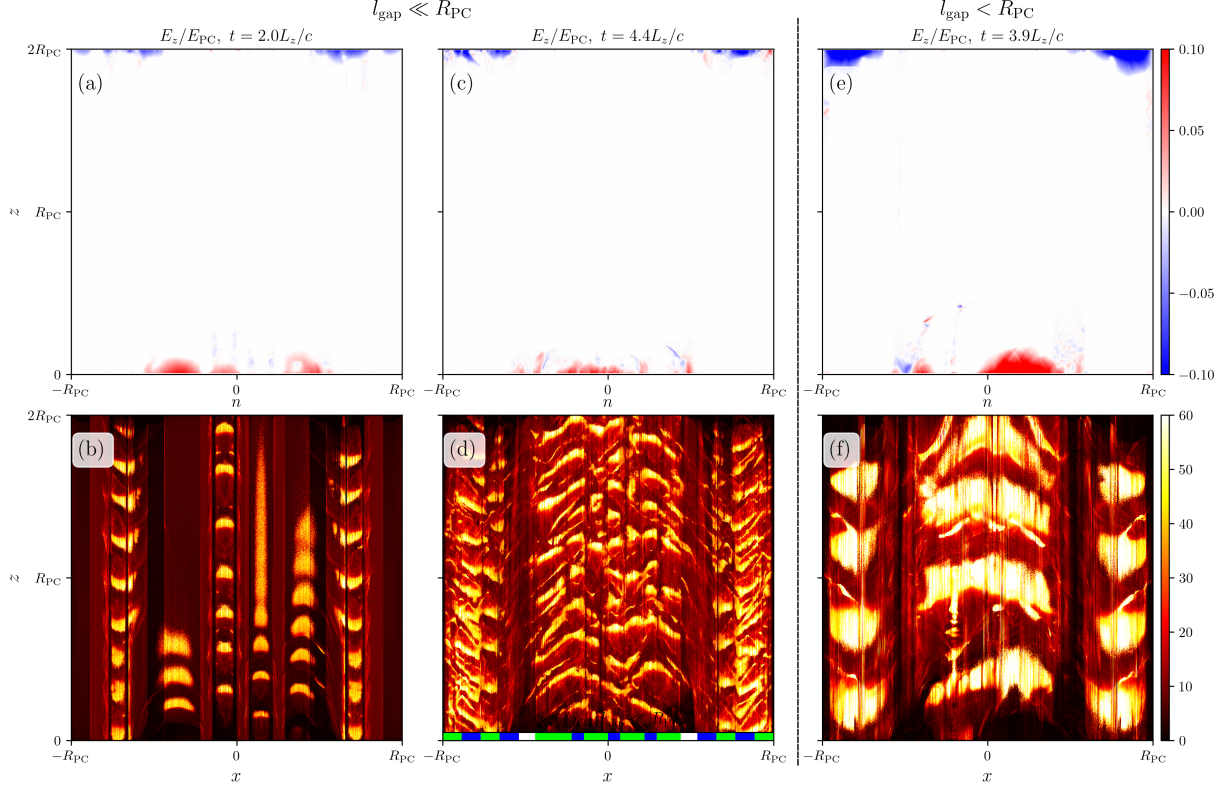


Figure 3.4: Evolution of the discharge in the 2D RS polar cap with a small, $l_{\text{gap}} \ll R_{\text{PC}}$, RS_{small} , and large, $l_{\text{gap}} \lesssim R_{\text{PC}}$, RS_{large} , gaps. The first column shows (a) electric field and (b) plasma density distribution during the initial phase of the evolution in RS_{small} , when initial patches still occupy a large area of a polar cap. The second column, panels (c)-(d), illustrates a later evolutionary stage when all patches are cleared of the initial plasma. The right column shows (e) the electric field and (f) density snapshots for a large gap, RS_{large} , at a later time when the initial patches are cleared of plasma.

cap. Namely, the gap opens at the stellar surface, and clouds of secondary plasma move toward the magnetosphere in zones of the positive current. In the regions of return current, a gap opens at high altitudes and launches plasma moving toward the star. It is important to note that the area with sub-GJ magnetospheric current also possesses an accelerating electric field and pair-producing discharge, which is in contrast with the SCLF model. The RS and SCLF models also differ in the zones of the super-GJ current, where the RS model has a gap with large parallel electric field $\approx E_0 \approx 4\pi\rho_{\text{GJ}}l_{\text{gap}}$. In the RS model, the current is carried by particles of one sign extracted from the cloud created during the previous discharge. In the SCLF model, the electric

field is much smaller because particles of both signs are present in the gap.

In the case of a small gap, $l_{\text{gap}} \ll R_{\text{PC}}$, Fig. 3.4a-d, the discharge operates in the intermittent regime already when the first patches are cleared of the initial plasma (Fig. 3.4a,b), in the zones of both super-GJ and return currents. At a later time, when all of the patches are cleared, the discharges communicate through the vacuum episodes similar to the discharges in the SCLF model. The complex shape of the clouds of secondary plasma demonstrates that the discharges are not fully synchronized across magnetic field lines. At the same time, a relatively large repetition rate, f , leads to smaller de-synchronization of the discharges (Fig. 3.4 c,d). We expect a similar “lava lamp” pattern of de-synchronized discharges in a situation of a realistically low repetition rate, $f \ll 1$, as described in Section 3.4.1.2. In the polar cap with a large gap, R_{Slarge} , $l_{\text{gap}} \lesssim R_{\text{PC}}$, see Fig. 3.4e-f, a quiet stage at the beginning of simulations with narrow, isolated patches is later replaced by a well-developed cascade with nearly synchronized discharges when all patches are cleared of the initial plasma. This dynamics is broadly similar to the evolution in the SCLF simulation with a large gap, $\text{SCLF}_{\text{large}}$.

3.4.3 3D Gap, Space-Charge-Limited Flow

In this Section, we study a 3D discharge corresponding to the current distribution in the 60° -inclined rotator. Here, almost half of the polar cap is occupied by the field lines with super-GJ current, and the other half is occupied by the field lines with negative return current (see Fig.3.1b). The choice of the simulation parameters (simulation $\text{SCLF}_{3\text{D}}$ in Table 3.1) provides a small gap, $l_{\text{gap}} \ll R_{\text{PC}}$, in the region of the return magnetospheric current, and marginally small gap, $l_{\text{gap}} < R_{\text{PC}}$, in the region of super-GJ current, due to different electric potentials developing

in these two zones (see Sec. 3.4.1.1)⁸. To study the desynchronization of the discharge, we split the polar cap into 6×6 square patches, with two of them being initially loaded with low-density plasma, which allows the formation of a gap at the beginning of the simulation. One of two “open” patches is located in the zone of super-GJ current, and another one - is in the region of return current. In other patches, up until certain time moments, we inject additional plasma to prevent charge starvation and to keep the accelerating electric field screened.

Overall, the evolution of the discharge is similar to the 2D SCLF simulations discussed in Section 3.4.1.1. The volume rendering of the plasma density for a well-developed discharge is shown in Fig. 3.5a. In the region of the super-GJ current, the transverse size of the patch limits the electric potential, leading to the formation of a large quasi-stationary gap near the stellar surface. The examples of such two isolated gaps can be seen in the early time slice of the electric field, E_z , Fig. 3.5f. The corresponding slice of the plasma density, Fig. 3.5e, shows the production of secondary plasma beyond the gap. At later times, other patches are cleared of the initial plasma and become available for discharge. Similar to 2D large gap simulations in Sec. 3.4.1.4, neighboring discharges synchronize on transverse scales $l_{\perp} \sim l_{\text{gap}}$, leading to a nearly coherent discharge across the super-GJ zone, as clearly visible in a slice of the electric field in Fig. 3.5. We start with only two relatively small patches where discharges are possible and observe that the discharge zone spreads across the polar cap when we stop suppressing pair formation. In a well-developed cascade, discharges produce large clouds of secondary plasma, which spread across all field lines, as shown in Fig. 3.5g,h.

The formation of the gap in the region of the return current starts at high altitudes. Similarly

⁸Further increase of the separation of scales, $R_{\text{PC}}/l_{\text{gap}}$, is challenging because of the computational cost of the simulations.

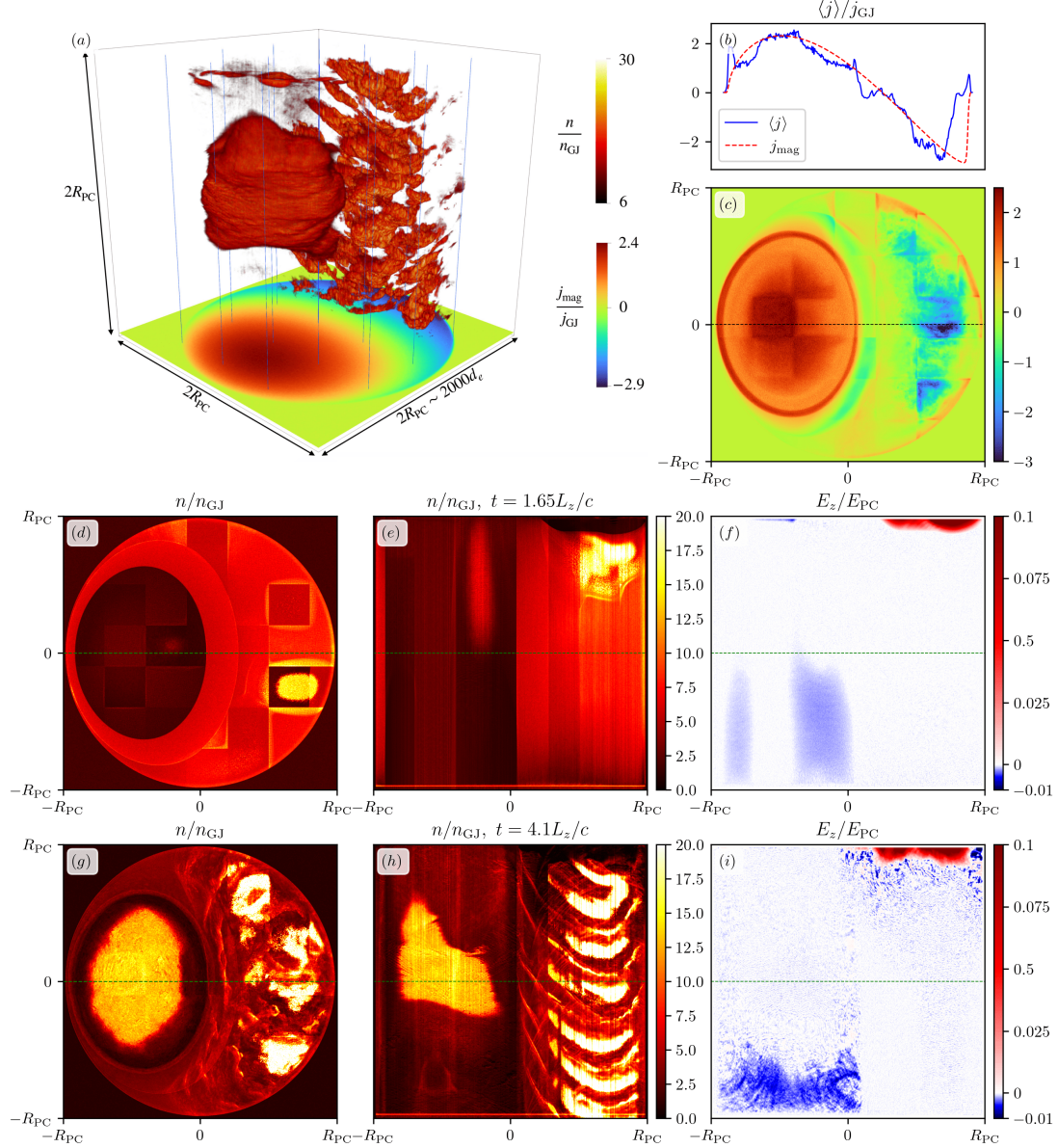


Figure 3.5: 3D SCLF discharge in a polar cap of a 60° -inclined rotator with a marginally small gap, $l_{\text{gap}} < R_{\text{PC}}$. Subplot (a) shows a volume rendering of the density of a secondary plasma. The plane at the bottom shows the position of the stellar surface, and the colormap shows the distribution of the magnetospheric current. The blue lines perpendicular to the stellar surface represent the magnetic field lines. Subplots (b) and (c) show the time-averaged distribution of the plasma current. The black dashed line in the plot (c) shows the position of the 1D slice in (b). Two lower rows show 2D slices of the plasma density (d,e,g,h) and electric field (f, i) through the center of the domain in two different projections. The dashed lines in each of the slices show the position of a slice in a perpendicular plane, e.g., the dashed line in (d) shows the position of slices (e) and (f). Snapshots (d-f) in the middle row show the initial phase of the simulation, when a lot of patches are still filled with initial plasma, and the lower row (g-i) shows a time when all patches are cleared of the initial plasma.

to a 2D discharge, a lack of electrons develops there, leading to a strong accelerating electric field, $E_z \sim E_0 \sim 4\pi|\rho_{\text{GJ}}|l_{\text{gap}}$. It leads to a smaller gap compared to the super-GJ field lines, and hence, a cyclic discharge can start within each initial patch. When all patches are cleared of the initial plasma, significant desynchronization of the discharges across magnetic field lines remains, as seen in the transverse slice of the plasma density, Fig. 3.5g. Here, a few isolated structures are seen, similar to the 3D volume rendering in Fig. 3.5a.

3.4.4 Evolution of the Twist of Magnetic Field Lines

As was discussed in Section 3.3.1, the persistent twist of the magnetic field can be maintained in the plasma-filled, i.e., FFE, state only. This condition is naturally broken in the gap zone. The lack of plasma can have two consequences: the growth of the accelerating electric field and the untwisting of the field lines (increasing $\partial\delta\mathbf{E}/\partial t$ vs. decreasing $\nabla \times \delta\mathbf{B}$ terms in Eq. 3.1). If magnetic field lines become untwisted in a significant fraction of the polar cap, the outflow of plasma along them would be similar to the free outflow into the vacuum, as it was implicitly assumed to be the case in original polar cap models of [Arons & Scharlemann \(1979\)](#) and [Ruderman & Sutherland \(1975\)](#), these models disregarded the presence of the magnetosphere in their estimates of the accelerating electric field. Such an outflow would be “charge-density-driven”. The polar cap and the magnetosphere are parts of the same electric circuit. In this circuit, the magnetosphere, which has a much larger inductance, determines the current density along magnetic field lines. In a real pulsar, if the magnetic field lines become untwisted in the large portion of the polar cap, an “untwisting wave” will propagate into the magnetosphere, where it will be eventually stopped, and a new wave will be sent toward the polar cap, restoring the twist. In our

setup, we model the presence of the magnetosphere by imposing the current j_{mag} in the hope that the system has an inductance high enough to maintain the prescribed current density. A situation where we end with almost a complete untwisting of magnetic field lines would mean that our setup does not model the system polar cap + magnetosphere correctly, which basically means that our simulation domain is not large enough. Therefore, checking for untwisting of magnetic field lines is an important consistency test of our model.

To study the evolution of the twist of magnetic field lines in our simulations, we compare the plasma current, j , with the prescribed magnetospheric current, j_{mag} . In the right column of Fig. 3.3, we compare the plasma and magnetospheric currents for simulations with different ratios of the gap height to the polar cap size. For small gaps, a significant fraction of any field line is loaded with high-multiplicity plasma, leading to the deviations of j from j_{mag} to be localized in a small volume. The magnetospheric twist is easily sustained by the pair production discharge in the polar cap zone, as shown in subplot 3.3e. The twist is preserved particularly well in the “lava-lamp” discharge (simulation `SCLFtails`) since the gaps, where FFE conditions are violated, are small and most of the domain is filled with dense plasma most of the time. In contrast, the simulation `SCLFsmall` starts with initially long, $l_{\perp} \lesssim l_{\text{gap}}$, “charge-density” driven, quasi-stationary gaps that are formed within narrow patches, as described in Section 3.4.1.1. These large gaps lead to a noticeable untwist of the field lines. This effect is particularly strong at the edges of open patches, where the electric potential is smaller due to proximity to the neighboring plasma patch, and, thus, particle acceleration and pair production are less efficient. However, when neighboring patches are cleared of the initial plasma, the electrostatic potential quickly increases, leading to a smaller gap size along the field lines. At this time, the gap becomes effectively shorter along the transverse dimension, $l_{\perp} \gtrsim l_{\text{gap}}$. The domain fills with dense plasma

before the complete untwisting of field lines, the initial magnetospheric twist is partially restored, and discharges enter into the “current-driven” phase. This can be seen in the subplot 3.3e as a decrease in the amplitude of the difference between the plasma current and the magnetospheric one, which is particularly noticeable at the boundaries of initial patches. The evolution of the twist in the polar cap with a large gap, simulation `SCLFlarge`, is overall similar to the evolution of the small gap case. The main difference is in a more significant deviation of the plasma current from the magnetospheric current, as shown in subplot 3.3f, because of the larger region sustaining parallel electric fields. Patches with a low plasma current (centered around the dashed lines in Fig. 3.3d,f) end up as non-sustaining the discharge, as described in Section 3.4.1.4. The field lines within these patches partially restore their twist due to the supply of plasma from the atmosphere.

Similar conclusions hold in the case of 3D discharges, as is shown in Figs 3.5b, c, where it is clear that the distribution of the plasma current overall follows the magnetospheric current. The deviations are localized to the boundaries of the patches, as clearly visible in Fig. 3.5c. Following the results of the twist evolution over a longer time in 2D simulations, we assume that the twist will be gradually restored, and the clearly visible patch boundaries will disappear over longer times (see Section 3.4.1.1).

Overall our numerical setup is capable of dealing with the untwisting of magnetic field lines. It never comes to complete untwisting on field lines, and the average plasma current is reasonably close to the prescribed magnetospheric current, j_{mag} . The size of the simulation domain is adequately large to simulate the presence of a large magnetosphere above the cascade zone. On the other hand, modeling of discharges with gap heights larger than the size of the polar cap, $l_{\text{gap}} \gtrsim R_{\text{pc}}$, would require a substantially larger computational domain to prevent the complete untwisting of magnetic field lines and will be addressed in future works.

3.5 Conclusions and Discussion

In this Letter, we explored the behavior of multidimensional pair-producing discharges within the polar zone of pulsar magnetospheres. Our main goal was to study the transverse coherence of these discharges, i.e., to understand what sets the smallest size of clusters of magnetic field lines that simultaneously possess an unscreened electric field necessary to trigger QED pair production. We found that current-driven discharges function in a cyclic pattern similar to the 1D model described by [Timokhin \(2010\)](#); [Timokhin & Arons \(2013\)](#). Moreover, we demonstrated that the synchronization between localized discharges across magnetic field lines occurs during plasma-starved episodes when the unscreened electric field on different lines can communicate as it would happen in a vacuum. Our main conclusion is that the transverse coherence scale of a discharge zone is comparable to the longitudinal gap size, l_{gap} .

3.5.1 Non-existence of sparks

We studied the possibility of spark formation, namely the existence of long-lived isolated discharge columns surrounded by plasma-depleted regions where particle acceleration and pair formation are suppressed due to the accelerating electric field being screened by the presence of sparks. In a purely electrostatic problem, like in a discharge between the plates of a flat capacitor, the accelerating potential along the magnetic field lines between spark columns is limited by the potential across the magnetic field between the sparks, the latter, being filled with plasma at the latter stage of the discharge, acting essentially as conducting boundaries for inter-spark space ([Ruderman & Sutherland, 1975](#)). This would lead to the polar cap filled with sparks separated by the distance $\sim l_{\text{gap}}$. However, the problem of polar cap discharges at the base of the FFE

magnetosphere is not an electrostatic one; the magnetosphere requires a certain twist of magnetic field line and hence a certain current density j_{mag} to support it. If the actual current density deviates from j_{mag} , a strong displacement current, $\partial E/\partial t$, appears – the accelerating electric field is current-driven. Hence, a strong electric field should appear between discharge columns, regardless of how small the separation between them is, if the twist of magnetic field lines requires a current that can not be sustained by plasma without pair production. So, the polar cap should be filled by discharges without empty zones if the current density supports pair formation. Our numerical simulations confirm this expectation.

In our setup, initially, we let a few discharge columns develop while keeping the accelerating electric field screened by dense plasma in the space between them. When these inter-discharge spaces are let to evolve freely, an accelerating electric field appears in them and drives pair formation. The whole polar cap becomes filled with discharges, with no plasma-depleted space left between them. This is true for both RS and SCLF regimes, as is evident in Figs. 3.3-3.4. If the width of the initial discharge column is smaller than the gap height, these columns expand until their width becomes comparable to the gap height; otherwise, a new discharge column will appear in the initial inter-discharge space. As in the 1D case, the new discharge is started by particles from the tail of the previous bursts of pair formation and not by particles coming from the NS surface, thus making the specific conditions for particle escape at the NS surface irrelevant for the start of the next discharge. Our code does not prevent the $\mathbf{E} \times \mathbf{B}$ drift, but we do not see any noticeable pair plasma drift in our simulations. This is to be expected as the accelerating potential variation across magnetic field lines is not large enough to make plasma move with respect to the NS with the velocity large enough to be visible at times scales of pair discharges (van Leeuwen & Timokhin, 2012).

In the lava-lamp type cascades, individual discharges are densely packed with no spaces between them. Even if separate emission regions are directly associated with the plasma flows within individual discharges, the resulting emission pattern would not result in individual subpulses. On the timescale of a single pulse, which is much longer than the discharge repetition rate, each discharge would contribute comparable flux to the overall broad emission profile. In principle, one could think of a possibility for subpulse-like features to appear due to individual discharges if the emission is generated (or decouples from the plasma) at the boundaries of individual discharges. This can be possible for emission mechanisms recently proposed by [Philippov et al. \(2020\)](#); [Melrose et al. \(2021a\)](#). The resulting emission pattern might look like a distorted “honeycomb”, rather than a bunch of spots within the polar cap. However, even in the latter case, drifting subpulses seen in some pulsars can not be caused by individual discharge columns. Pulsars exhibiting subpulse drift phenomena are biased closer to the death line in the $P - \dot{P}$ diagram, which means that they have large gaps comparable to or exceeding the radius of the polar cap. We demonstrated that for these conditions, the discharges are proceeding coherently across magnetic field lines, ruling out the presence of many independent discharge zones in old-ish pulsars. Such pulsars can not have distinct emission zones associated with separate discharge columns.

In our simulations, we start with the polar cap filled by plasma and find that in the well-developed phase, the discharges operate on all magnetic field lines where pair formation is possible, regardless of the conditions on particle supply from the NS surface. In other words, an existing FFE magnetospheric configuration can always support itself by generating plasma along all magnetic field lines, with no plasma-depleted regions between discharges. We conclude that a spark-filled polar cap is incompatible with the FFE pulsar magnetosphere.

3.5.2 Implication for discharges in old pulsars

In our paper, we did not consider gaps exceeding the size of the polar cap, as would be appropriate for old, less energetic pulsars. Their magnetospheres can differ from the fully FFE configurations assumed in this work. While in FFE solutions, the twist of the field lines is dictated by the conditions at the light cylinder, our local gap simulations show significant untwisting of the field lines that are unable to fully adjust to the magnetospheric current. This is particularly prominent in the initial phases of our simulations when extended thin patches are unable to ignite the “current-driven” discharge because of the voltage limitation caused by the transverse extent. Such discrepancies between the local current produced in the discharge zone and the global FFE magnetospheric solution will likely result in large-scale “breathing” of the magnetosphere (see Sect. 3.4.4) on timescales comparable to the light-crossing time of the light-cylinder, $\sim P/2\pi$, with potential implications for the short-term radio nulling. Understanding these issues requires global simulations that include accurate modeling of pair production in the polar discharge for conditions appropriate for old pulsars.

3.5.3 On repetition rate of the cascades

A significant issue beyond this Letter is the repetition rate of the pair cascades – the time interval between two successive bursts of pair formation along the same magnetic field lines, $\Delta\tau_{\text{rep}}$. As we conjectured in Sec. 3.4.1.2, it is likely that under realistic conditions, the small repetition rate (large $\Delta\tau_{\text{rep}}$) is caused by a continuous flux of particles reversed from the pair cloud propagating away from the discharge zone. Significant improvements, even in 1D simulations, are required to accurately capture the long-term evolution of electric fields in pair clouds and

resolve this issue. A particularly stunning example of the importance of the long-term evolution of discharges is the repetition rate on field lines carrying the return current. In local simulations, the discharge starts at the upper boundary of the simulation domain, however, it is unclear how far the plasma-starved zone extends into the magnetosphere. It is possible that the region with the unscreened electric field can first form at the null point, where the GJ charge density is zero (e.g., [Bransgrove et al., 2023](#)). If this assumption holds, it is likely that electric fields could easily synchronize across the extended return current zone, significantly enhancing the coherence of the discharge. This is another question that should be addressed using next-generation global simulations that include an accurate discharge model.

Chapter 4: Origin of the Pulsed 20 TeV Emission in the Vela Pulsar

4.1 Abstract

We present a semi-analytical description of the formation of very-high-energy pulsed emission in the Vela pulsar recently observed by H.E.S.S. Our model is motivated by recent particle-in-cell simulations of relativistic magnetic reconnection with strong synchrotron cooling. We show that high-energy synchrotron photons emitted by reconnection-accelerated particles efficiently produce electron-positron pairs via the Breit-Wheeler process. The density of secondary pairs exceeds the supply from the polar cap and results in a self-regulated plasma magnetization parameter of $\sigma \approx \text{few} \times 10^7$. Electrons and positrons accelerate to Lorentz factors comparable to σ and emit the observed GeV radiation via synchrotron process and ~ 10 TeV photons by up-scattering the soft synchrotron photons emitted by secondary pairs. Our model self-consistently accounts for the luminosity of the very-high-energy radiation, $L_{\text{TeV}} \approx 10^{-4} L_{\text{GeV}}$.

4.2 Introduction

Pulsars are broad-band emitters of non-thermal radiation covering more than 20 decades in photon energy (Philippov & Kramer, 2022). The studies of pulsars, discovered as radio sources, were revolutionized by the detection of their γ -ray emission using *Fermi*, with more than 200

γ -ray pulsars known at the moment (Abdollahi et al., 2020). An absence of the super-exponential cutoff in the spectrum of the γ -ray emission (Abdo et al., 2013) indicates the high-altitude location of the emitting region; otherwise, the γ -rays would be absorbed in a strong magnetic field near the surface. The mismatch of rotational phases at which the radio and γ -ray lightcurves peak (e.g., Kuiper & Hermsen, 2015) additionally disfavors the inner magnetosphere as the source of the observed high-energy radiation. The lightcurves of the γ -ray emission are strongly peaked, pointing to their caustic nature as demonstrated in multiple theoretical studies (e.g., Dyks & Rudak, 2003; Bai & Spitkovsky, 2010; Kalapotharakos et al., 2014).

Recent global three-dimensional particle-in-cell simulations of pulsar magnetospheres conclusively identify the zone of particle acceleration responsible for the high-energy emission to be located in the current sheet near and beyond the light cylinder (Cerutti et al., 2016; Philippov & Spitkovsky, 2018; Kalapotharakos et al., 2018; Hakobyan et al., 2023a). The lightcurves extracted from these simulations are formed by photons strongly beamed along the direction of the magnetic field upstream to the current sheet, resulting in a double-peaked caustic structure broadly resembling *Fermi* observations. Simulations also measure the efficiency of converting the spin-down power into the high-energy radiation to be of the order of 1 – 10%, which is consistent with observations (e.g., Marelli et al., 2011). Whether the γ -rays are powered by synchrotron (Lyubarskii, 1996; Uzdensky & Spitkovsky, 2014; Cerutti et al., 2016) or curvature (Kalapotharakos et al., 2018) radiation process is, however, still debated, with important implications for identifying the largest energies of emitting particles.

Recent observations by the H.E.S.S. collaboration revealed pulsed emission with photon energies up to 20 TeV from the Vela pulsar (H. E. S. S. Collaboration et al., 2023). The light curve of the detected emission is aligned in phase with one of the γ -ray peaks in the 1 – 10 GeV

energy interval, indicating a similar emission site for GeV- and TeV- photons. This discovery is significant since it serves as a direct probe of the maximum Lorentz factor of the particles in the zone where the high (HE) and very-high-energy (VHE) emission is produced. This is because the inverse Compton (IC) mechanism responsible for VHE emission is capable of up-scattering photons up to the energies equal to the energy of the target electron (Blumenthal & Gould, 1970). In particular, for producing the 20 TeV photon, the Lorentz factors of the emitting electrons/positrons have to be in excess of 4×10^7 .

In this chapter, we consider the possible emission scenario powering 20 TeV photons from the magnetospheric current sheet. We find that for the conditions of the Vela pulsar, the plasma supply of the current sheet occurs mainly by converting synchrotron photons into electron-positron pairs via the Breit-Wheeler process, $\gamma\gamma \rightarrow e^-e^+$ (Breit & Wheeler, 1934). These secondary pairs then emit soft, 0.1 – 1eV, synchrotron photons that serve as targets for the IC scattering by particles accelerated in the current sheet. We demonstrate that the density of produced secondary pairs and the energy density of soft photons result in the VHE signal consistent with the H.E.S.S. detection. The paper is organized as follows. In Section 4.3, we review the model of the pulsar magnetosphere and summarize the properties of the synchrotron emission from the current sheet; in Section 4.4, we describe a semi-analytical model of the HE and VHE spectrum of the Vela pulsar.

4.3 Magnetospheric Current Sheet

4.3.1 Plasma Feeding from the Polar Cap

Global solutions for the structure of the pulsar magnetosphere demonstrate the presence of the global twist of the open magnetic field lines (Contopoulos et al., 1999; Spitkovsky, 2006; Timokhin, 2006). The twist sets the global distribution of the magnetospheric current, j_{mag} , which is comparable, but not exactly equal, to the characteristic current, $j_{\text{GJ}} = \rho_{\text{GJ}}c$, and ρ_{GJ} is the Goldreich-Julian (GJ) charge density (Goldreich & Julian, 1969). Particle-in-cell simulations of the space-charge-limited flow in the polar cap show that only field lines with super-GJ current, $j_{\text{mag}}/j_{\text{GJ}} > 1$, and volume return current, $j_{\text{mag}}/j_{\text{GJ}} < 0$, can develop accelerating electric fields and drive pair cascades (Timokhin & Arons, 2013). At the same time, global simulations (e.g., Philippov et al., 2015; Philippov & Spitkovsky, 2018) show that the current sheet near the light cylinder is supplied by plasma produced in discharges on the field lines that carry the return current, $j_{\text{mag}}/j_{\text{GJ}} < 0$, for any obliquities that are not too close to orthogonal (see sketch of the magnetosphere in Fig. 4.1a). As we discuss below, most of the plasma produced in these discharges propagates toward the stellar surface, and uncertainty arises as to what is the density of the plasma that feeds the current sheet.

Simulations show cyclic behavior of the discharges, with periods of the active production of high-density pair plasma with multiplicities $\mathcal{M} \equiv \rho/\rho_{\text{GJ}} \gg 1$ (e.g., Timokhin & Harding, 2019) followed by quiet periods when electric fields are screened and no particle acceleration occurs (Timokhin, 2010; Timokhin & Arons, 2013). The example of a simulation of a discharge on field lines carrying a return current is shown in Fig. 4.1 (panels b-e). It shows that the plasma-starved

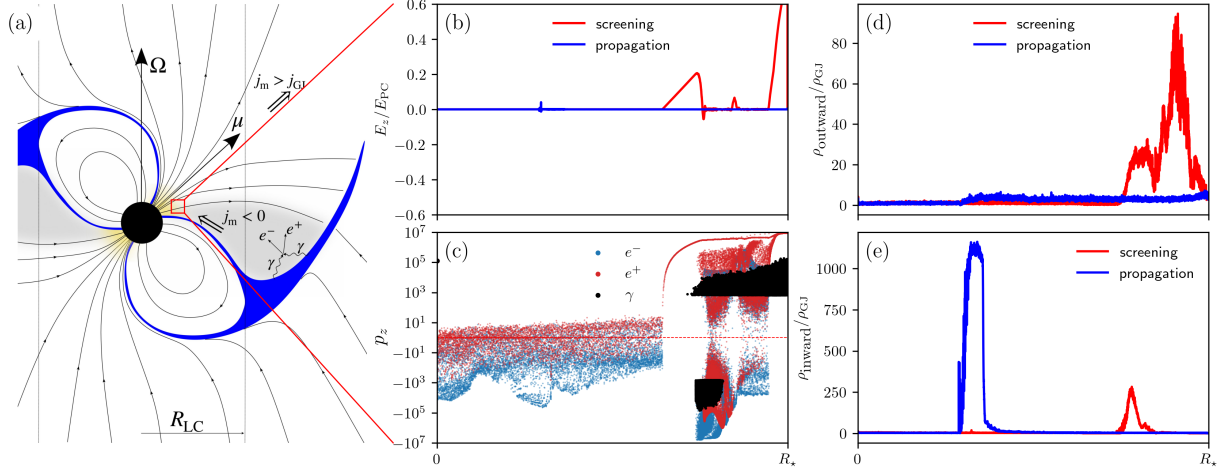


Figure 4.1: Model of current sheet plasma supply from polar pair production discharges. (a) Schematic structure of the global magnetosphere of the inclined pulsar. The shadowed grey region possesses a return volume current, $j_{\text{mag}} < 0$, the blue region shows the current sheet and the yellow zone in the polar region highlights the region of the pair-production discharge. The red inset shows the location of the discharge on field lines with return magnetospheric current. Panels (b)-(e) show snapshots of a particle-in-cell simulation of the discharge in the return current zone. (b) Accelerating electric field during the burst of pair production (red) and the following quiet phase (blue). (c) The phase space, $z - p_z$, of photons and pairs during the burst of pair production. Blue, red, and black dots represent electrons, positrons, and photons, correspondingly. (d) The density of the plasma propagating towards the magnetosphere, ρ_{outward} , at the same moments in time as panel (b). (e) Density of the plasma propagating towards the stellar surface, ρ_{inward} .

zone forms at high altitudes, $r \sim R_*$, above the stellar surface (see also [Timokhin & Arons, 2013](#)). Here, the pair cascade is initiated by an outward-propagating beam of positrons extracted from the cloud produced in the previous discharge episode (visible in the particle/photon phase space, $z - p_z$, shown in panel c). This leads to the creation of a short burst of relatively high-density plasma, $\mathcal{M} \sim 80$, propagating away from the star (see red lines in subplot d). When the electric field is fully screened, only a small fraction of the particles, $\mathcal{M} \sim \text{few}$, propagates outward¹ (see blue lines in subplots b, d, e).

The low-density plasma propagates along the magnetic field lines into the magnetosphere, reaching the upstream of the current sheet at $R \sim R_{\text{LC}}$. The plasma magnetization parameter,

¹This is because the electric field oscillations that are reversing the particles in the inward-propagating cloud are continuously screened by ongoing pair production ([Tolman et al., 2022](#)).

$\sigma_{\text{LC}} = B_{\text{LC}}^2 / (4\pi \mathcal{M} n_{\text{GJ}} m_e c^2)$, estimated for parameters of Vela, will be extremely high, $\sigma_{\text{LC}} \sim 10^{10} / \mathcal{M} \sim 10^9$. As we show below, in situ pair production near the current sheet can provide higher plasma density, limiting the magnetization parameter.

4.3.2 Particle Acceleration and Synchrotron Radiation from the Current Sheet

Relativistic magnetic reconnection is an efficient source of accelerated particles (e.g., [Sironi & Spitkovsky, 2014](#)). The key parameter that characterizes particle acceleration efficiency is the magnetization parameter in the upstream region, $\sigma = B_{\text{up}}^2 / (4\pi n_{\text{up}} m_e c^2)$, where B_{up} and n_{up} are the magnetic field strength and plasma density in the upstream of the current sheet. Reconnection in pulsar magnetospheres is strongly affected by synchrotron radiative cooling (e.g., [Lyubarskii, 1996](#); [Uzdensky & Spitkovsky, 2014](#); [Hakobyan et al., 2019](#); [Chernoglazov et al., 2023](#)). Radiation reaction becomes dynamically important for particles with Lorentz factors comparable to the radiation-reaction-limit, γ_{rad} , which can be estimated by balancing the accelerating electric force and radiation drag:

$$eE_{\text{rec}} = \frac{\sigma_T}{4\pi} B_{\text{up}}^2 \gamma_{\text{rad}}^2, \quad (4.1)$$

where σ_T is the Thomson cross-section, $E_{\text{rec}} \approx \beta_{\text{rec}} B_{\text{up}}$ is the reconnection-driven electric field, and $\beta_{\text{rec}} \sim 0.1$ is the reconnection rate. In Vela, at the light cylinder, $B_{\text{LC}} \approx 7 \times 10^4 \text{G}$, and $\gamma_{\text{rad}} \approx 10^5$, leading to $\sigma_{\text{LC}} \gtrsim \gamma_{\text{rad}}$.

If acceleration and radiation occur in the same physical region with $B \approx B_{\text{up}}$, particles with Lorentz factors comparable to γ_{rad} emit synchrotron photons of universal energy, the “burnoff” limit:

$$\varepsilon_{\text{bo}} = \frac{3}{2} \hbar \frac{e \tilde{B}_{\perp}}{m_e c} \gamma_{\text{rad}}^2 = \frac{9}{4\alpha_F} \frac{E_{\text{rec}}}{B_{\text{up}}} m_e c^2 \approx 16 \text{MeV} \left(\frac{\beta_{\text{rec}}}{0.1} \right), \quad (4.2)$$

where α_F is the fine-structure constant, and $\tilde{\mathbf{B}}_{\perp} = \mathbf{E} + \mathbf{v} \times \mathbf{B}/c$ is the effective perpendicular magnetic field. Reconnecting current sheets, however, are highly inhomogeneous, and physical locations of particle acceleration and radiation do not coincide. In particular, as shown by 3D simulations of radiative reconnection in Chapter 2, particles can be efficiently accelerated along X-lines and reconnection outflows, where the effective perpendicular magnetic field is small, and synchrotron losses are negligible. When accelerated particles encounter stronger fields, they radiate photons of energy

$$\varepsilon_p = 16 \text{ MeV} \frac{\tilde{B}_{\perp}}{B_{\text{up}}} \left(\frac{\gamma}{\gamma_{\text{rad}}} \right)^2. \quad (4.3)$$

In the regime of strong cooling, $\sigma \gtrsim \gamma_{\text{rad}}$, applicable to pulsar current sheets, simulations show that the Lorentz factor of accelerated particles is limited by the magnetization parameter, $\gamma_{\text{max}} \sim \sigma$. Particles of the highest energies start efficiently radiating when they encounter even weak magnetic fields, $\tilde{B}_{\perp} \sim B_{\text{up}}(\gamma_{\text{rad}}/\sigma)$. As a result, the spectrum of synchrotron photons extends beyond the burnoff limit (see Chapter 2):

$$\varepsilon_{\text{cut}} \approx 16 \text{ MeV} \left(\frac{\sigma}{\gamma_{\text{rad}}} \right). \quad (4.4)$$

Simulations also show that highest energy synchrotron photons are strongly beamed along the reconnecting magnetic field, $\mathbf{B}_{\text{up}}$, which is consistent with phenomenological assumptions of γ -ray lightcurve modeling (e.g., [Bai & Spitkovsky, 2010](#)).

To further examine the scaling in Eqn. (4.4), we perform a 3D simulation of a relativistic magnetic reconnection with no guide field and strong synchrotron cooling $\gamma_{\text{rad}}/\sigma = 0.1$, using the setup described in Chapter 2. For this simulation, we resolve the skin depth of the upstream

plasma by 2 numerical cells, which corresponds to twice the resolution of the simulation presented in Chapter 2 for the same cooling. Our plasma magnetization parameter, σ , is equal to 50, and the box size is 2000 cells in both directions in the plane of the current sheet and 1200 cells in the perpendicular direction.

In Fig. 4.2, we show the time-averaged spectra of pairs and photons, resulting in compensated spectra of the pairs and the photons. For small Lorentz-factors, $\gamma < \gamma_{\text{rad}}$, when the radiation cooling is negligible, the power-law index of the particle distribution is -1 , similar to non-radiative simulations (e.g., Sironi & Spitkovsky, 2014). At higher energies, synchrotron cooling becomes efficient, leading to the steepening of the spectrum. As shown in subplot (a), for $\gamma > \gamma_{\text{rad}}$, the spectral index changes to ≈ -2 . The exponential cutoff is located at $\gamma \approx \sigma$. These results are consistent with the simulation shown in Chapter 2.

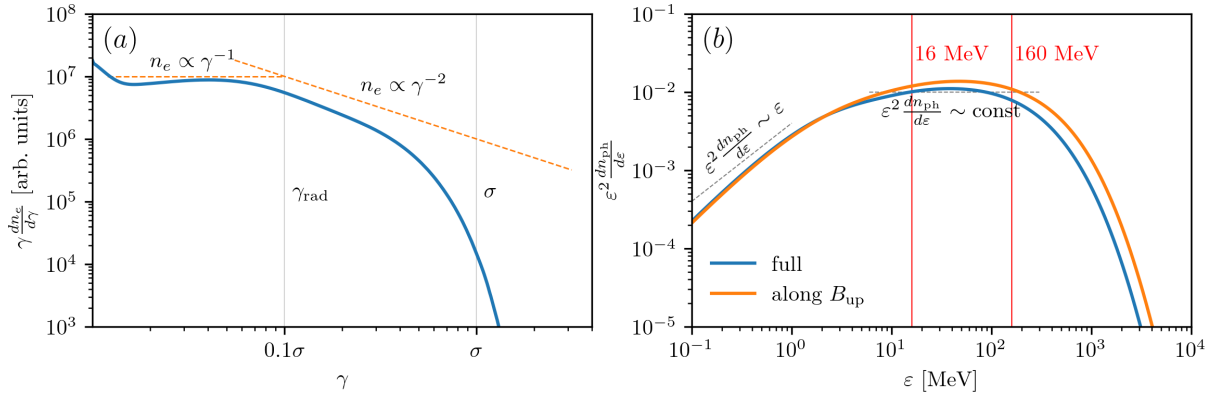


Figure 4.2: The time-averaged pair (a) and photon (b) spectra from the simulation of a 3D current sheet with strong synchrotron cooling, $\gamma_{\text{rad}}/\sigma = 0.1$. Dynamically strong cooling limits the maximum Lorentz factor of the particles at $\gamma_{\text{max}} \sim \sigma$. The resulting photon spectrum has an exponential cutoff at energy exceeding the burnoff limit, $\varepsilon_{\text{cut}} \approx 16\text{MeV}(\sigma/\gamma_{\text{rad}}) \gg 16\text{MeV}$. The blue line shows the normalized spectrum of photons in all directions, and the orange line represents a spectrum of photons propagating nearly along the upstream magnetic field.

For a power-law particle distribution function, $dn_e(\gamma)/d\gamma \equiv f_e(\gamma) \propto \gamma^{-p}$, the corresponding synchrotron spectrum is $\varepsilon^2 dn_{\text{ph}}(\varepsilon)/d\varepsilon \equiv \varepsilon^2 f_{\text{ph}} \sim \varepsilon^{(3-p)/2}$ (Rybicki & Lightman, 1979). For

$p = 1$, such a scaling gives a rising synchrotron spectrum at low energies, $\varepsilon^2 f_{\text{ph}} \propto \varepsilon$, for photon energies below the burnoff limit, ε_{bo} , given by Eqn. (4.2). At higher energies, the scaling is violated, as the spectrum and pitch-angles of emitting particles become time and energy-dependent. The photon spectrum measured in the simulations, and shown in Fig. 4.2b, has a plateau at energies $\varepsilon > 16\text{MeV}$, followed by an exponential cutoff at $\approx 160\text{ MeV}$ consistent with the scaling $\varepsilon_{\text{cut}} \approx 16\text{ MeV } (\sigma/\gamma_{\text{rad}})$ for the explored case $\gamma_{\text{rad}}/\sigma = 0.1$. In Vela, the observed γ -ray spectrum experiences a cutoff at $\sim 4\text{GeV}$, necessitating $\sigma \sim \text{few} \times 10^7$. Below, we show that similar values of the magnetization parameter can be explained by self-regulated pair production near the current sheet.

4.4 Model of the HE and VHE Emission

In this Section, we formulate a semi-analytical model of the emission mechanism behind the observed TeV pulsed emission in the Vela pulsar. We assume that the radiation is produced in the reconnecting current sheet near the light cylinder. For Vela, the rotational period is $P_{\star} = 0.089\text{s}$, and the strength of the magnetic field at the light cylinder $B_{\text{LC}} = 7.5 \times 10^4\text{ G}$. For these parameters, the radiation reaction limit is $\gamma_{\text{rad}} = 1.2 \times 10^5$. As we discussed in Section 4.3.1, the field lines carrying volume return current can have a small plasma supply from the polar cap, resulting in a very high magnetization at the light cylinder, $\sigma_{\text{LC}} \approx 10^9$. At the same time, the high-energy synchrotron photons emitted by particles accelerated in the current sheet can produce additional pair plasma through the Breit-Wheeler process, $\gamma\gamma \rightarrow e^-e^+$ (Breit & Wheeler, 1934). Hereafter, we assume that two-photon pair production proceeds in the optically thin regime, i.e., the characteristic optical depth $\tau_{\gamma\gamma} \ll 1$ ². The process of TeV photon emission can be split into

²This assumption is confirmed later.

a few main steps.

1. Pairs accelerated in the current sheet emit synchrotron photons (Current Sheet Synchrotron, CSS) as described in Section 4.3.2.
2. Most of the photons escape to infinity, with a small fraction of them colliding and producing electron-positron pairs upstream of the current sheet.
3. Produced secondary pairs have non-zero pitch angles and emit low-energy secondary synchrotron photons (Secondary Pair Synchrotron, SPS). Potentially, these photons can serve as additional targets for the $\gamma - \gamma$ pair production.
4. Simultaneously, the density of the secondary plasma alters the value of the plasma magnetization parameter, σ_{LC} , and, hence, the spectrum of CSS. The combined process of photon emission and pair production results in self-regulation of the plasma magnetization.
5. Inverse Compton (IC) scattering of soft SPS photons by pairs accelerated in the current sheet results in the production of VHE photons.

For detailed analysis, we use the findings from the local simulations of the relativistic reconnection with strong synchrotron cooling (Section 4.3.2) to prescribe the distribution function

of particles and photons in a form

$$\frac{dn_e(\gamma)}{d\gamma} = \begin{cases} \gamma^{-1}, & \gamma < \gamma_{\text{rad}} \\ \gamma_{\text{rad}} \gamma^{-2} \exp\left(-\frac{\gamma}{\sigma}\right), & \gamma > \gamma_{\text{rad}} \end{cases} \quad (4.5)$$

$$\varepsilon^2 \frac{dn_{\text{ph}}(\varepsilon)}{d\varepsilon} = \begin{cases} \varepsilon, & \varepsilon < \varepsilon_{\text{bo}} \\ \varepsilon_{\text{bo}} \exp\left(-\frac{\varepsilon}{\varepsilon_{\text{cut}}}\right), & \varepsilon > \varepsilon_{\text{bo}} \end{cases} \quad (4.6)$$

The corresponding spectra are shown in blue in subplots (a) and (b) of Fig. 4.3, correspondingly³.

4.4.1 Pair Production

The strength of radiative effects in reconnection can be parameterized using the dimensionless parameter of magnetic compactness parameter,

$$l_B = \frac{2\sigma_T U_B L}{m_e c^2}, \quad (4.7)$$

where σ_T is Thomson cross-section, and $L \sim R_{\text{LC}}$ is the length of the current sheet, and $U_B = B^2/8\pi$ is the magnetic field energy density. The synchrotron cooling time, i.e., the timescale when a particle radiates away most of its energy, can be expressed as $t_{\text{cool}} \approx 1/(\gamma l_B)$. For the parameters at the light cylinder of the Vela pulsar, the magnetic compactness is $l_B \approx 0.1$, so that particles with Lorentz factors $\gamma \gtrsim 10$ radiate most of their energy before escaping from the

³The observed GeV spectrum of Vela is harder at most phases, $\varepsilon^2 dn_{\text{ph}}/d\varepsilon \sim \varepsilon^{2-\Gamma}$, with $\Gamma \approx 1.6 - 1.8$ for the two peaks in the lightcurve (Abdo et al., 2010b). We expect this effect to be mainly due to anisotropy of high-energy photons expected for $\sigma_{\text{LC}}/\gamma_{\text{rad}} \approx 100$ to be more prominent compared to our simulated case $\sigma_{\text{LC}}/\gamma_{\text{rad}} \approx 10$ (see hint of this effect in Fig. 4.2). Observers seeing photons accumulated in a caustic nearly aligned with the field line will naturally see an excess of high-energy photons, emitted nearly along the magnetic field line, with respect to the lower-energy, less anisotropic ones.

light cylinder region. Thus, most of the dissipated magnetic energy, which is first converted into plasma energy, is immediately converted into photons, which escape with the speed of light as long as the optical depth is small. Then, the rate of magnetic energy supply, $v_{\text{rec}}U_B$, where v_{rec} is the reconnection velocity that determines the rate of supply of fresh magnetic flux, is balanced by the radiation escape, $U_{\text{rad}}c$. This results in $U_{\text{rad}} \sim \beta_{\text{rec}}U_B$. For a photon spectrum in the form of Eq. (4.6), this means that the energy density of $\sim 16\text{MeV}$ and higher energy photons is comparable to $\beta_{\text{rec}}U_B$.

While most of the emitted synchrotron photons escape the system, some of them interact and create electron-positron pairs, $\gamma\gamma \rightarrow e^-e^+$. The approximate expression for the corresponding pair-production rate in a unit volume can be obtained by assuming that pair production occurs only for pairs of photons with energies ε_1 and ε_2 , such that $\varepsilon_1\varepsilon_2 = 2m_e^2c^4$. In this approximation, the $\gamma - \gamma$ cross-section is then given (Gould & Schröder, 1967; Akhiezer & Berestetskij, 1985):

$$\sigma_{\gamma\gamma} \approx 0.1\sigma_T\varepsilon_2\delta\left(\varepsilon_2 - \frac{2m_e^2c^4}{\varepsilon_1}\right). \quad (4.8)$$

The Lorentz factor of the produced pairs is then $\gamma \approx \max(\varepsilon_1, \varepsilon_2)/2m_e c^2$. Equation (4.8) allows to estimate the optical depth for a photon with energy $\varepsilon_1 = \varepsilon_{\text{bo}}$:

$$\tau_{\gamma\gamma} \approx 0.1\sigma_T R_{\text{LC}} \frac{U_{\text{CSS}}(\varepsilon_2)}{\varepsilon_2} \approx 0.1\sigma_T R_{\text{LC}} \frac{U_{\text{CSS}}}{\varepsilon_{\text{bo}}}, \quad (4.9)$$

where $\varepsilon_2 = 2m_e^2c^4/\varepsilon_{\text{bo}}$, and $U(\varepsilon)/\varepsilon \sim \varepsilon f_{\text{ph}} = \text{const}$ for $\varepsilon < \varepsilon_{\text{bo}}$ for the CSS spectrum given by Eqn. (4.6). For parameters of the Vela pulsar, $\tau_{\gamma\gamma} \sim 5 \times 10^{-6} \ll 1$, justifying assumption of the small optical depth. For the delta-function assumption for the interaction, the pair production

rate per unit volume for the power-law distributed photons is given by (Svensson, 1987):

$$Q_{\gamma\gamma} \approx 0.5\sigma_T c n_{\text{ph}}^2 \int_{2m_e c^2}^{\infty} d\varepsilon \frac{1}{\varepsilon} \tilde{f}(\varepsilon) \tilde{f}\left(\frac{m_e^2 c^4}{\varepsilon}\right), \quad (4.10)$$

where $\tilde{f}_{\text{ph}}$ is the normalized distribution function of the photons given by Eqn. (4.6), $\int d\varepsilon \tilde{f}_{\text{ph}}(\varepsilon) = 1$, and the photon number density, $n_{\text{ph}} \approx \beta_{\text{rec}} U_B / \varepsilon_{\text{bo}}$. Since the photon energy distribution steepens at the burnoff limit, the integral is accumulated only at $\varepsilon \lesssim \varepsilon_{\text{bo}}$. The result is given by (e.g., Hakobyan et al., 2023b): $Q_{\gamma\gamma} \approx 2\sigma_T c n_{\text{ph}}^2 \approx 10^{10} \text{cm}^{-3}/\text{s}$.

In order to find an energy-dependent injection rate of the secondary pairs, we use a more accurate approximation of the $\gamma - \gamma$ interaction rate proposed by Aharonyan et al. (1983); Boettcher & Schlickeiser (1997):

$$Q_{\gamma\gamma}(\gamma) \approx \frac{3}{32} c \sigma_T n_{\text{ph}}^2 \int_{\gamma}^{\infty} d\tilde{\varepsilon}_1 \frac{\tilde{f}_{\text{ph}}(\tilde{\varepsilon}_1)}{\tilde{\varepsilon}_1^3} \int_{\frac{\tilde{\varepsilon}_1}{4\gamma(\tilde{\varepsilon}_1 - \gamma)}}^{\infty} d\tilde{\varepsilon}_2 \frac{\tilde{f}_{\text{ph}}(\tilde{\varepsilon}_2)}{\tilde{\varepsilon}_2^2} \cdot \left\{ \frac{4\tilde{\varepsilon}_1^2}{\gamma(\tilde{\varepsilon}_1 - \gamma)} \ln \left(\frac{4\tilde{\varepsilon}_2 \gamma (\tilde{\varepsilon}_1 - \gamma)}{\tilde{\varepsilon}_1} \right) - 8\tilde{\varepsilon}_1 \tilde{\varepsilon}_2 + \frac{2(2\tilde{\varepsilon}_1 \tilde{\varepsilon}_2 - 1)\tilde{\varepsilon}_1^2}{\gamma(\tilde{\varepsilon}_1 - \gamma)} - \left(1 - \frac{1}{\tilde{\varepsilon}_1 \tilde{\varepsilon}_2} \right) \frac{\tilde{\varepsilon}_1^4}{\gamma^2 (\tilde{\varepsilon}_1 - \gamma)^2} \right\}, \quad (4.11)$$

where $\tilde{\varepsilon} \equiv \varepsilon / m_e c^2$ is the dimensionless photon energy, γ is the Lorentz factor of produced pairs. The normalization of the photon distribution function, n_{ph} , can be found more accurately as $\beta_{\text{rec}} U_B / \langle \varepsilon \rangle$, where averaging is done over the photon spectrum given by Eqn. (4.6). In order to compute the integral, one needs to assume the value of σ_{LC} , which itself depends on $Q_{\gamma\gamma}$ and sets the high-energy cutoff of the CSS spectrum. However, as mentioned above, the integral is not sensitive to the exact value of the high-energy cutoff and is accumulated at energies below the

burnoff limit. The total injected number density over a light-crossing time of the light cylinder is given by: $n_e^{\text{sec}} = \tau_{\text{LC}} \int Q_{\gamma\gamma} d\gamma \approx 10^7 \text{cm}^{-3}$. For the parameters of Vela, this density translates into the value of the magnetization parameter, $\sigma_{\text{LC}} = 4 \times 10^7$, which is consistent with the γ -ray cutoff observed by *Fermi*, $\varepsilon_{\text{cut}} = 16 \text{MeV} (\sigma_{\text{LC}} / \gamma_{\text{rad}}) \approx 4 \text{GeV}$.

The resulting energy-dependent injected distribution function is shown by a dashed line in Fig. 4.3a. The low-energy part has a spectral index of -1 , as follows from the delta-function approximation for a single power-law distribution function of photons (Boettcher & Schlickeiser, 1997). Pairs of high energies are produced by the interaction of photons with energies $\varepsilon_{\text{high}} > \varepsilon_{\text{bo}}$ with photons of energies $\varepsilon_{\text{low}} \approx m_e^2 c^4 / \varepsilon_{\text{high}}$. Steepening of the photon energy distribution function for $\varepsilon > \varepsilon_{\text{bo}}$ leads to the steepening of the injected function to γ^{-2} .

4.4.2 Steady-state Pair Spectrum

Injected pairs experience synchrotron cooling in the upstream magnetic field, B_{up} , leading to modification of the injected distribution function. The resulting steady-state distribution function of the secondary pairs is given by the time-independent Boltzmann equation (e.g., Blumenthal & Gould, 1970):

$$\frac{d}{d\gamma} (\dot{\gamma} f_e^{\text{sec}}) = Q_{\gamma\gamma}(\gamma) - \frac{f_e^{\text{sec}}}{\tau_{\text{LC}}}, \quad (4.12)$$

where the synchrotron cooling rate is $\dot{\gamma} = -l_B \gamma^2 / \tau_{\text{LC}}$, the first term in the R.H.S. is the source of secondary pairs due to $\gamma - \gamma$ pair production given by Eqn. (4.11), and the last term describes the escape of particles from the light cylinder region, $\tau_{\text{LC}} \equiv R_{\text{LC}} / c = P_{\star} / 2\pi$. Equation (4.12) has

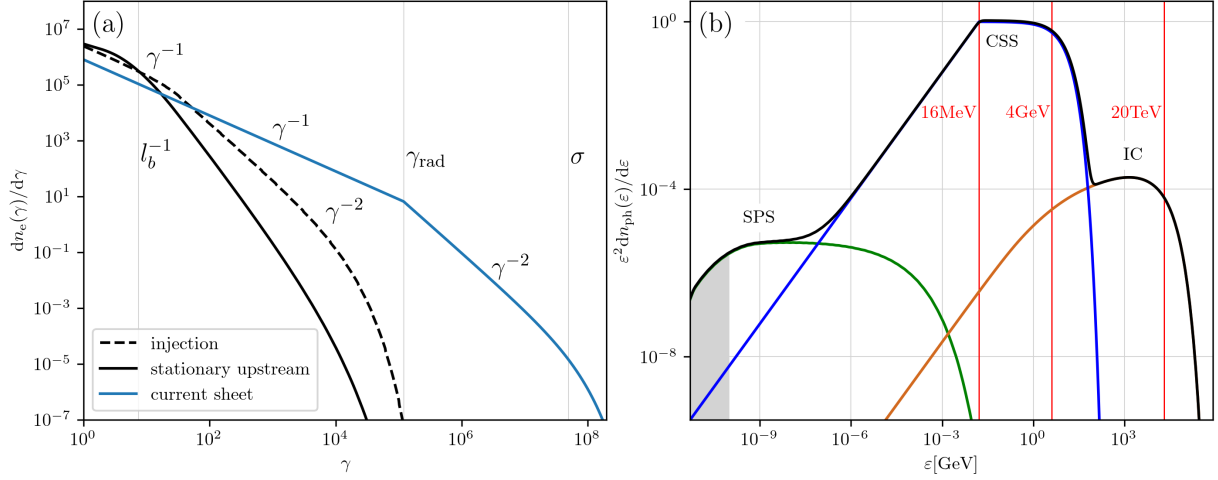


Figure 4.3: Particle distribution function and resulting photon spectrum. (a) The spectrum of pairs injected upstream of the current sheet through $\gamma-\gamma$ pair production (dashed line) and the resulting steady-state spectrum of the injected particles that takes into account their synchrotron cooling (solid black line). The blue line shows the spectrum of pairs accelerated in the current sheet. (b) Resulting photon spectrum: the blue line shows the synchrotron spectrum from particles accelerated in the current sheet (CSS), the green line shows the spectrum of the synchrotron photons emitted by secondary pairs (SPS), and the orange line is the spectrum of the inverse-Compton upscattered photons. The shadowed region at low photon energies shows the photons interacting with the highest energy pairs in the Thomson regime.

an analytical solution:

$$f_e^{\text{sec}}(\gamma) = \frac{1}{\gamma^2} \exp\left(-\frac{1}{C\gamma\tau_{\text{LC}}}\right) \int_{\gamma}^{\infty} \frac{Q_{\gamma\gamma}(x)}{C} \exp\left(\frac{1}{Cx\tau_{\text{LC}}}\right) dx, \quad (4.13)$$

where we used a boundary condition, $\gamma^2 f_e^{\text{sec}}(\gamma) \rightarrow 0$ when $\gamma \rightarrow \infty$. The resulting steady-state distribution function of secondary pairs, $f_e^{\text{sec}}(\gamma)$, is shown in Fig. 4.3a. Synchrotron cooling of the injected particles leads to steepening of the steady-state distribution function for Lorentz factors where the cooling is efficient on a timescale τ_{LC} . The cooling break is located at $\gamma \sim l_B^{-1} \sim \text{few}^4$.

⁴A noticeable pileup at energies $\gamma \lesssim 1/l_B$ is a feature of cooling breaks for distribution functions with hard power-laws, γ^{-p} , $p < 2$, as originally discussed by Kardashev (1962).

4.4.3 Secondary Pair Synchrotron Spectrum

We calculate the spectrum of SPS photons under a simplified assumption of the uniform pitch angle distribution of secondary pairs. The resulting spectrum, $f_{\text{ph}}(\varepsilon) = Q_{\text{syn}}\tau_{\text{LC}}$, is shown by a green line in Fig. 4.3b. This spectrum has a rising component at low energies, below 1eV, where photons are emitted by the low-energy part of the stationary distribution function of secondary pairs, $\gamma \lesssim 1/l_B$, where cooling is negligible. The more energetic part of the SPS photons in the energy range 10eV $\div$ 1keV is emitted by the higher energy particles, $\gamma > 1/l_B$, where the synchrotron cooling is efficient.

The energy density of the synchrotron emission of secondary pairs (SPS), U_{SPS} , can be calculated as $\varepsilon^2 f_{\text{ph}}(\varepsilon) = n_{\text{sec}} m_e c^2 l_B \tilde{f}_e(\gamma) (\varepsilon / \hbar \omega_B)^{3/2}$. The energy density of soft radiation, $\lesssim 1\text{eV}$, is dominated by photons emitted by particles with Lorentz factors $\gamma \lesssim 1/l_B$. Not surprisingly, it results in $U_{\text{SPS}} \sim U_p = n_{\text{sec}} m_e c^2 / l_B$ that corresponds to the energy density of particles with $\gamma \sim 1/l_B$ that radiate all of their energy away as synchrotron radiation. Another way to estimate the energy density of these particles can be obtained by considering the fact that a $\tau_{\gamma\gamma}$ fraction of CSS photons produce pairs so that $U_p \sim \tau_{\gamma\gamma} U_{\text{CSS}} \sim \tau_{\gamma\gamma} \beta_{\text{rec}} U_B$. For an optically thin system, the total energy of SPS photons is small, $U_{\text{SPS}} \ll U_{\text{CSS}}$.

In principle, the highest end of the SPS spectrum, $\varepsilon \gtrsim 1\text{keV}$ can serve as targets for additional $\gamma - \gamma$ collisions with $\sim 1\text{GeV}$ photons. However, for the parameters of Vela, the magnetospheric non-thermal keV radiation is dominated by the current sheet emission, so the contribution of the SPS to pair production is negligible. Note that the energy density of the thermal X-ray radiation, $\sim L_X^{\text{th}} / (4\pi c R_{\text{LC}}^2)$, at the light cylinder, is comparable to the non-thermal X-rays, which do not affect pair production rate because of the steepening of the photon distribu-

tion function at energies beyond the burnoff limit. Thus, there is no additional photon source in the magnetosphere that can affect the pair production rate estimated in Section 4.4.1.

4.4.4 Inverse Compton Signal

Soft photons with energies $\varepsilon_{\text{soft}}$ can be upscattered to multi-TeV energies by the leptons with Lorentz factors $\gtrsim 10\text{TeV}/m_e c^2 = 2 \times 10^7$. The most efficient interaction occurs in the Thomson regime when $4\gamma\varepsilon_{\text{soft}}/m_e c^2 \lesssim 1$. For the required energies of the target pairs, the energy of the soft photons is $\varepsilon_{\text{soft}} \sim m_e c^2/\sigma \sim 0.01 - 0.1\text{eV}$. Synchrotron photons of these energies can be emitted by particles with Lorentz-factors $\gamma \sim \text{few}$ as follows from Eqn. (4.3), under the assumption $\tilde{B}_\perp \approx B_{\text{up}}$ ⁵. At low energies, the total photon spectrum is dominated by the SPS photons, and hence, most of the IC power is produced by upscattering of the SPS photons, not the initial CSS: the shadowed region in Fig. 4.3b shows the part of the photon spectrum which is most efficiently upscattered by very energetic pairs with $\gamma \sim \sigma$.

The power density of the IC emission can be found as $P_{\text{IC}} \approx (4/3)\sigma_T c n_e \langle \gamma^2 \rangle U_{\text{soft}}$ (Rybicki & Lightman, 1979), where $U_{\text{soft}} \approx U_{\text{SPS}}$. At the same time, the power density of the CSS emission is $P_{\text{CSS}} \approx 2\sigma_T c n_e \langle \gamma^2 \tilde{B}_\perp^2 / 8\pi \rangle$. As was shown by Hakobyan et al. (2023b), for the current sheet emission, $\langle \gamma^2 \tilde{B}_\perp^2 / 8\pi \rangle \approx \gamma_{\text{rad}}^2 U_B$. Thus, the ratio of the IC and CSS powers is

$$\frac{P_{\text{IC}}}{P_{\text{CSS}}} \approx \frac{\langle \gamma^2 \rangle U_{\text{SPS}}}{\gamma_{\text{rad}}^2 U_B} \approx \frac{\sigma}{\gamma_{\text{rad}}} \frac{U_{\text{SPS}}}{U_B} \approx \frac{\sigma}{\gamma_{\text{rad}}} \tau_{\gamma\gamma} \beta_{\text{rec}}. \quad (4.14)$$

Here, $\langle \gamma^2 \rangle \equiv \int \gamma^2 f_e(\gamma) d\gamma / \int f_e d\gamma \approx \sigma \gamma_{\text{rad}}$, for the distribution function given by Eqn. (4.5).

⁵In models by Harding & Kalapotharakos (2015); Harding et al. (2018), small particle pitch angle is excited by the absorption of radio waves, and particles of larger energies are producing the required soft synchrotron signal. We think these models are less robust as they require assumptions about the location and geometry of the radio emission region.

For the optical thickness $\tau_{\gamma\gamma}$ given by Eqn (4.9), the ratio of luminosities of SCC and IC photons can then be estimated as:

$$\frac{L_{\text{IC}}}{L_{\text{syn}}} = \frac{P_{\text{IC}}}{P_{\text{syn}}} \sim 10^{-4}. \quad (4.15)$$

This result nicely matches with the observed GeV to TeV luminosity ratio, $L_{\text{TeV}}/L_{\text{GeV}} \sim 10^{-4}$.

To compute the spectrum of the IC photons, we use the full expression for the inverse Compton scattering probability, which includes both the Thomson and Klein-Nishina regimes (Blumenthal & Gould, 1970; Lefa et al., 2012). The resulting photon injection rate into the unit volume is:

$$Q_{\text{IC}}(\varepsilon_{\text{IC}}) = 3\sigma_T c \frac{n_e n_{\text{ph}}}{m_e c^2} \int_0^\infty \int_{\gamma_{\min}}^\infty W(\gamma, \varepsilon, \varepsilon_{\text{IC}}) \tilde{f}_{\text{ph}}(\varepsilon) \tilde{f}_e(\varepsilon) d\gamma d\varepsilon, \quad (4.16)$$

where

$$\begin{aligned} \gamma_{\min} &= 0.5 \frac{\varepsilon_{\text{IC}}}{m_e c^2} \left(1 + \sqrt{1 + \frac{m_e^2 c^4}{\varepsilon_{\text{IC}} \varepsilon}} \right), \\ W(\gamma, \varepsilon, \varepsilon_{\text{IC}}) &= \frac{1}{\gamma \eta} \left[2q \ln q + (1 - q) \left(1 + 2q + \frac{\eta^2 q^2}{2(1 + \eta q)} \right) \right], \\ \eta &= \frac{4\gamma \varepsilon}{m_e c^2}, \quad q = \frac{\varepsilon_{\text{IC}}}{\eta(m_e c^2 \gamma - \varepsilon)}. \end{aligned} \quad (4.17)$$

Here, $\tilde{f}_{\text{ph}}$ and $\tilde{f}_e$ are normalized distribution functions, ε is the energy of the soft photon, and ε_{IC} is the energy of the upscattered photon. The normalization coefficient n_e is the number density of electrons, i.e., $n_e \approx 10^7 \text{cm}^{-3}$, and n_{ph} now is the density of SPS and CSS photons. The distribution function of IC photons is given again $f_{\text{IC}} = \tau_{\text{LC}} Q_{\text{IC}}$. The resulting spectrum of the IC photons is shown by the orange line in Fig. 4.3b. The more exact calculation confirms the ratio of the IC and CSS luminosities to be $\sim 10^{-4}$, in agreement with the observations. The peak of the IC spectrum is located at slightly lower energies than appears in observation. However, given the simplicity of our model, we consider the agreement to be quite good.

4.5 Conclusion and Discussion

In this paper, we present a semi-analytical model of the low, high, and very high-energy radiation in the Vela pulsar. Our conclusions are based on recent 3D kinetic simulations of current sheets undergoing relativistic reconnection with strong synchrotron cooling, as well as simulations of pair-production discharges in polar caps of pulsars. We self-consistently evaluate the supply of electron-positron pairs in the current sheet. Based on the local simulations of the pair production discharge in the polar cap, we demonstrate that field lines carrying volume return current can only supply the current sheet with plasma of multiplicity $\mathcal{M} \sim 10$. However, we find that the in-situ pair production via collisions of synchrotron photons emitted by reconnection-accelerated particles is efficient. The density of the produced secondary plasma exceeds the density of the plasma from the polar cap and determines the magnetization of the current sheet. We find that self-regulated value of the magnetization is about $\sigma_{\text{LC}} \approx 4 \times 10^7$, which is sufficient to explain the cutoff measured by *Fermi*, $\sim 16\text{MeV}(\sigma_{\text{LC}}/\gamma_{\text{rad}}) \sim \text{few} \times \text{GeV}$ (Abdo et al., 2013).

Secondary pairs are produced in the upstream with random pitch angles and radiate abundant synchrotron soft photons of energies $0.1 - 10^3\text{eV}$. These photons produce the luminosity from NIR, $L_{\text{NIR}} \sim 10^{-6}L_{\text{GeV}}$, to X-rays, $L_{\text{X}} \sim 10^{-3}L_{\text{GeV}}$, which is consistent with measurements at $\sim 0.5\text{eV}$ (Zyuzin et al., 2013) and $10 - 100\text{keV}$ X-rays (Manzali et al., 2007). Their number density in the X-ray range, however, is subdominant to the low-energy tail of the current sheet synchrotron radiation and does not lead to enhancement of the $\gamma - \gamma$ pair production. Upscattering of soft photons of energies $\sim 0.1 - 1\text{eV}$ by the reconnection-accelerated particles leads to a substantial signal in the TeV range. Our model predicts the TeV luminosity to be $L_{\text{TeV}} \approx 10^{-4}L_{\text{GeV}}$, which agrees with observations (H. E. S. S. Collaboration et al., 2023).

In our model, the dissipated magnetic energy is directly processed into synchrotron photons, implying radiation energy density of $\beta_{\text{rec}} U_B$, where $\beta_{\text{rec}} \approx 0.1$ is the reconnection rate in the collisionless regime, and $U_B = B_{\text{LC}}^2/8\pi$ is magnetic field energy density. This condition also sets the characteristic fraction of the full spin-down luminosity converted into the gamma-ray emission, $L_\gamma/L_{\text{sd}} \sim \beta_{\text{rec}} \sim 0.1$. Observations show lower values, $L_\gamma/L_{\text{sd}} \sim 10^{-2}$ (Smith et al., 2023). A lower efficiency of the γ -ray production can be explained by the fact that in high-obliquity pulsars such as Vela (as inferred by both radio polarization, Johnston et al. 2001, and Chandra X-ray torus modeling, Ng & Romani 2008) dissipation is less efficient, as revealed by global PIC simulations (Philippov et al., 2015; Cerutti et al., 2016). We expect that ratios between luminosities in different bands calculated using our model will be similar in oblique rotators. However, various geometrical and beaming effects need to be self-consistently accounted for using global magnetospheric solutions.

Another pulsar with observed emission in the TeV range is Crab Pulsar (Ansoldi et al., 2016a), and its high-energy spectrum is more peculiar: while it spans up to 1.5 TeV, it does not have easily separable components that can be associated with the synchrotron or inverse Compton emission mechanisms. A possible explanation is a more efficient $\gamma - \gamma$ pair production since the magnetic field at the light cylinder ≈ 20 times larger than in Vela. A stronger magnetic field leads to substantially increased production of pairs both by collisions of the current sheet synchrotron photons with themselves and with substantially boosted X-ray synchrotron photons emitted by secondary pairs. This can be consistent with the fact the total luminosity of the Crab is dominated by the X-rays, $L_X \sim 10 \times L_{\text{GeV}}$. Excess pair production can potentially lead to switching to a “weak” cooling regime, $\gamma_{\text{rad}} \gtrsim \sigma$, and cessation of the pair production until the current sheet ejects the excess plasma. This non-stationary regime cannot be treated by the

steady-state approach developed in this paper and requires numerical simulations of magnetic reconnection with self-consistent pair production, similar to the ones performed by [Hakobyan et al. \(2019\)](#).

Chapter 5: Conclusion and Future Work

5.1 Summary

In this thesis, I performed extensive studies of relativistic plasmas in pulsar magnetospheres. In Chapter 2, I utilized relativistic radiative Particle-In-Cell simulations to study the dynamics and radiation from a relativistic three-dimensional current sheet with efficient synchrotron cooling. For the first time, using three-dimensional large-scale simulations, I demonstrated that strong inhomogeneity of the magnetic field inside the current sheet leads to efficient particle acceleration to the energies significantly above the theoretical radiation-reaction-limited regime. In the work, I demonstrated that these high-energy particles are short-lived and emit high-energy photons with energies far beyond the “burnoff” limit. My results also demonstrate the highly efficient acceleration of heavier ions in the current sheet formed by strongly cooled electron-positron plasma.

In Chapter 3, I explored the key mechanisms of plasma production in the pulsar magnetosphere - the polar discharge. In this Chapter, I constructed a model of a multi-dimensional discharge with accurate incorporation of the quantum electrodynamics (QED) effect. In this work, I waded into the mire of the observed variability of the radio emission (e.g., micro-structure and individual pulses). Here, I demonstrated that the discharge does not have to happen simultaneously at all open field lines. Instead, the polar cap can be split into a “honeycomb” pattern

of separate discharges with self-consistently emerging transverse coherence within each cell. I discovered that the intermittent nature of the multi-dimensional discharge leads to the pair production columns filling the entire polar cap zone, challenging the popular theory of separated “sparks” within the polar cap.

In Chapter 4, I applied the results of two previous Chapters to develop a semi-analytical model of the formation of recently discovered very-high-energy pulsed emission in Vela pulsar. The phase-resolved light curves of Vela demonstrate that most of the high-energy emission comes from a similar location, indicating the magnetospheric current sheet to be the primary radiation site. Here, I explained how the deficit of the plasma supply from the polar discharge into the magnetospheric current sheet makes the current sheet responsible for its plasma feeding. I demonstrated that the current sheet enters into a self-regulatory regime: a lack of upstream plasma leads to more efficient particle acceleration and a higher injection rate of the plasma produced in photon collisions. Remarkably, for the parameters of Vela, my model is capable of explaining the observed luminosities across nearly 15 orders of photon energies: from the near-infrared spectral band with the photon energies of $\sim 0.1 - 1$ eV up to TeV gamma-rays reaching energies of $\sim 10^{13}$ eV.

5.2 Future Work

Observations of the thermal X-ray emission from the polar cap of pulsars demonstrate the lack of heating of the stellar surface by the bombardment by relativistic particles from the discharge. Under the framework of the intermittent discharge model presented in my thesis, this can be naturally explained by the long quiet periods between discharges when all electric fields

are screened, and no particle acceleration happens. Back-of-the-envelope calculations demonstrate that for Vela these quiet periods need to be ~ 100 times longer than the periods of active discharge. At the same time, it remains to be seen how the screened electric field can be maintained for such a long time without the continuous production of plasma. Unfortunately, my attempts to resolve the conundrum using 1D PIC simulations failed. As more pairs were produced, the noise level in the simulations substantially increased, and the evolution of the electric field became unreliable. In my future work, I plan to address this issue by performing a study of 1-dimensional discharge using a relativistic Vlasov code, *GKeyll* (Juno et al., 2018; Hakim & Juno, 2020). *GKeyll* solves the Vlasov equation as a partial differential equation (as opposed to the Monte-Carlo method used in PIC codes), adopting a highly accurate discretization using a discontinuous Galerkin scheme. Figure 5.1 demonstrates the preliminary results of the electric field screening by the continuously injected plasma. As is seen from the density plots (a,b), the low-density tail continuously forms behind the leading plasma cloud. The tail is formed by a continuous reversal of the plasma particles (the reversed particles are seen in the lower half of the phase space in subplot c) by the long-lived oscillations of the electric field within the plasma cloud (subplot d). As shown in Figure 5.1, highly accurate Vlasov simulations allow the resolution of oscillations even for the plasma density increase by a factor of a few thousand. My future steps will involve incorporating self-consistent photon transfer into *GKeyll* and performing first discharge simulations.

Another direction I want to pursue is to adopt the QED pair production technique, described in Chapter 3, to accurately simulate the global pulsar magnetosphere. The motivation for this work is to study the feedback of the magnetosphere on regular twisting/untwisting of the magnetic field lines in the gap zone. As discussed in Chapter 3, it is particularly important

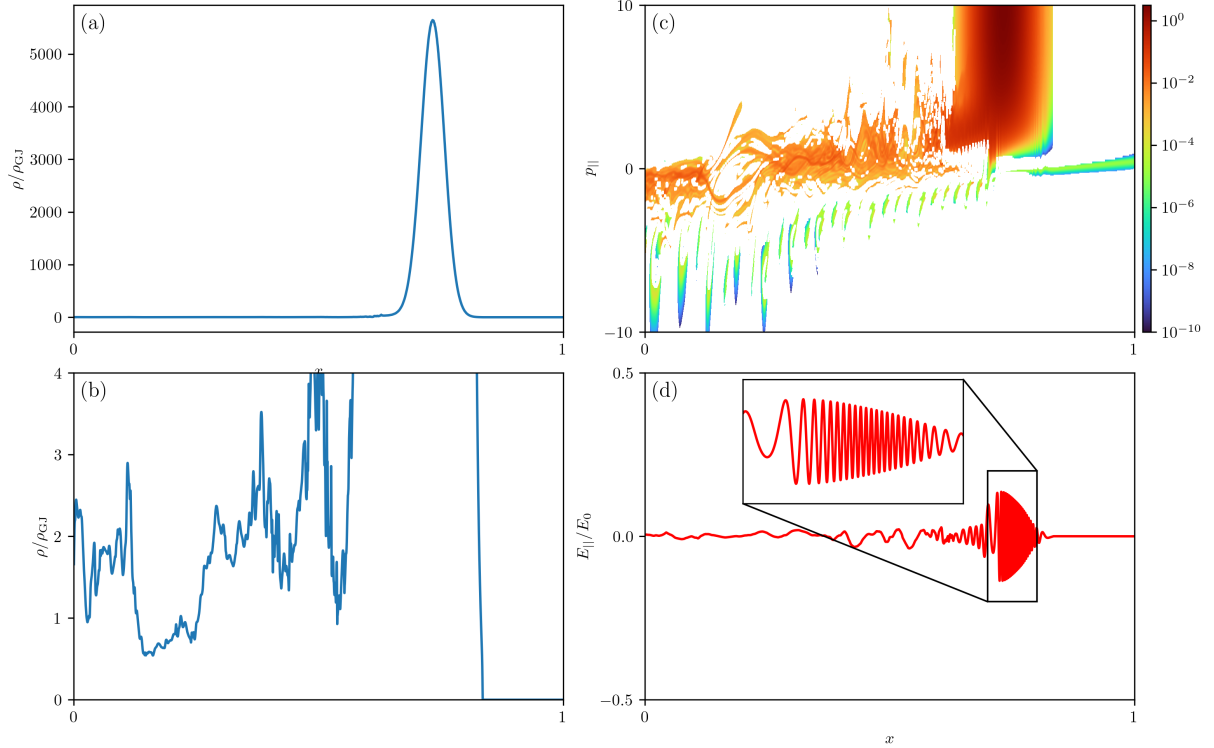


Figure 5.1: Preliminary study of the electric field screening by plasma discharge using *GKeyll*. Subplot (a) shows the total injected density while subplot (b) zooms into the low-density backward propagating plasma that forms a tail of the leading cloud. Subplot (c) depicts the phase space of the distribution function, with most of the plasma moving outwards while some plasma is continuously reversed. Subplot (d) shows the accelerating electric field, with the inset showing the pile-up of the forward-propagating waves.

for understanding the radio variability in old pulsars, which can not sustain a small gap. An intriguing possibility is that such “breathing” behavior can lead to a new state of the magnetosphere different from the one obtained in force-free and PIC simulations with abundant plasma supply (the possibility of this “breathing” was initially demonstrated by [Chen et al. 2020](#)). I hope that my new simulations, performed for the old pulsars with absent pair production at the light cylinder, the accurate modeling of the discharge behavior can uncover the nature of the various observed variabilities in the pulsar radio emission. I have already implemented QED processes into a new GPU-based code *Entity*, which will allow me to perform global QED-PIC simulations of axisymmetric pulsars.

Appendix A: Appendices for Chapter 2

A.1 Varying the fraction of ions and the mass ratio

In this appendix, we compare the efficiency of ion acceleration in the current sheet with a varying number density fraction of ions, $f_i = n_i/n_{\text{up}}$, as well as the mass ratio, m_i/m_e . For this study, we use simulations with a spatial resolution of $d_e = \Delta x$ and a box size $L_x/d_e = 1000$. In Sec. 2.5.1, we have demonstrated that in simulations with no cooling and mass ratio $m_i/m_e = 1$ the evolution of positrons and ions is identical, and, obviously, there is no difference for different fractions of ions in the simulation. The difference is most pronounced when pairs are strongly cooled. In this section, we compare the evolution of the distribution of ions in the strong cooling regime, $\gamma_{\text{syn}}/\sigma = 0.2$.

In Fig. A.1a, b, c we show a comparison of the distribution of ions measured in the steady state of the reconnection, $t = 4L_x/c$, for varying number density fractions, $f_i = 1\%, 10\%, 49\%$, while keeping $m_i/m_e = 1$ (simulation shown in Fig. A.1a is identical to the previously discussed simulation 1dx1kCoo102). From these results, we see that as long as the number density contribution of ions is small, $f_i \ll 1$, their distribution remains very hard, $dn_i/d\gamma \propto \gamma^{-1}$ for energies $\gamma \gtrsim \sigma$. In case when the number of positrons is low, and ions constitute almost half of the particle population (Fig. A.1c), we observe a significant deviation, with the resulting ion distribution, $dn_i/d\gamma \propto \gamma^{-1.7}$, being similar to that observed in the simulations without synchrotron cooling.

In this case, the pressure contribution of hot ions to the dynamics of the reconnection layer and plasmoids is no longer negligible. As a result, plasmoids in this case are much larger than in strongly cooled simulations with a lower number of ions, and this leads to more efficient trapping of the accelerated ions and, ultimately, to a steeper power-law index. In all cases, the spectrum of the strongly cooled pairs is not affected, as their acceleration primarily occurs in the X-point region.

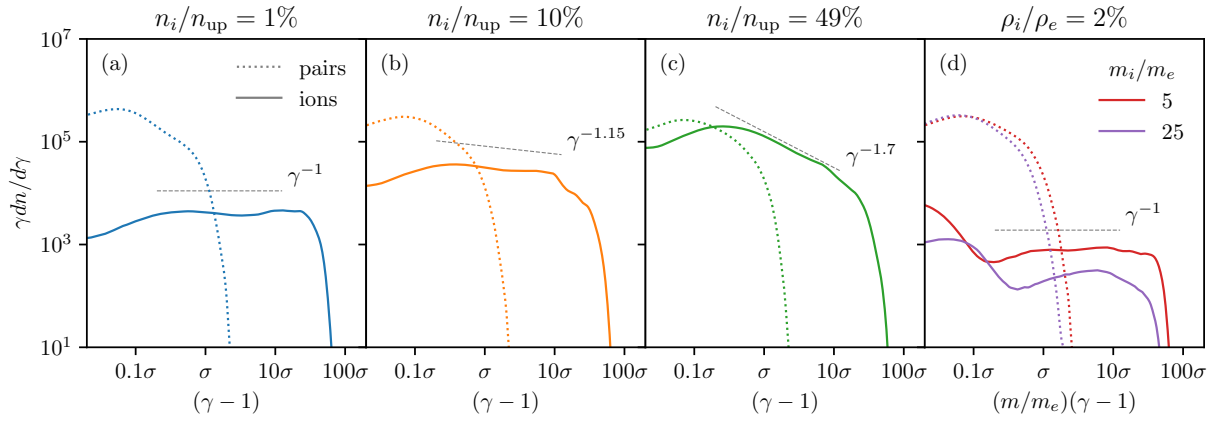


Figure A.1: Distribution functions of pairs (dashed) and ions (solid) in simulations containing a different fraction of ions, $f_i = n_i/n_{\text{up}} = 1\%$, 10% , 49% (panels a...c), and different mass ratio $m_i/m_e = 5, 25$ (panel d). As long as the mass density fraction of ions is relatively small, $n_i m_i / (n_e m_e) \ll 1$ (which is not the case in panel c), the power-law index of their distribution is insensitive to either f_i , or the m_i/m_e . The x -axis of panel (d) is normalized by the mass of the particles.

Another limitation of our simulations is the reduced mass ratio, m_i/m_e , which we have chosen to be 1 for most of the discussions throughout the paper. The reason for this choice is to maximize the separation of scales between the Larmor radius of the most energetic ions, r_L , and the size of the box, and to minimize the effects of boundary conditions on particle acceleration. To confirm that the reduced mass ratio $m_i/m_e = 1$ adequately captures the general properties of the acceleration of ions, we carried out two simulations with the parameters described above, but with larger mass ratios, $m_i/m_e = 5$ and $m_i/m_e = 25$. To keep the mass fraction of ions constant,

$\rho_i/\rho_e = 2\%$, we also decreased their number density fraction, f_i . The results of these simulations are shown in Fig. A.1d. Evidently, the power-law slope of the distribution is only marginally affected by the mass ratio, with the power-law tail extending to energies $\gg m_e\sigma$ (where σ takes into account the corresponding mass ratio as shown in Eq. (2.1)).

A.2 Varying the resolution

To ensure that our results are robust, and do not depend on the resolution, i.e., the number of cells per upstream plasma skin depth, in this section we vary this parameter, $d_e = 1 \dots 3\Delta x$, while keeping $m_i/m_e = 1$, and $f_i = 1\%$. For a direct comparison, we will keep other parameters fixed in the simulation: the ratio of the size of the simulation box to the skin depth, $L_x/d_e = 1000$, and the magnetization parameter, $\sigma = 50$.

We compare two particular cases for the strong cooling, $\gamma_{\text{syn}}/\sigma = 0.2$ (namely, simulations 3dx3kCool02 and 1dx1kCool02), where the results are expected to be the most sensitive to the resolution, as these simulations demonstrate a strong compression of plasmoids, making the local skin depth potentially difficult to resolve. We choose the efficiency of particle acceleration, i.e., the total distribution function as well as the distribution of energy gains in $E > B$ regions, as the main criteria for comparing different numerical resolutions. The results of this comparison are shown in Fig. A.2. Panels (a) and (b) show slices in $y = 0$ plane of the plasma density at time $t = 2.7L_x/c$. Panel (c) shows the normalized distribution of both pairs and ions, as well as the contributions by particles that encountered $E > B$ during their lifetime. Panel (d) shows the normalized distribution functions of the energy gained by both pairs and ions during their flythrough in the $E > B$ region. Evidently, the cross-sections of the current sheet are similar in

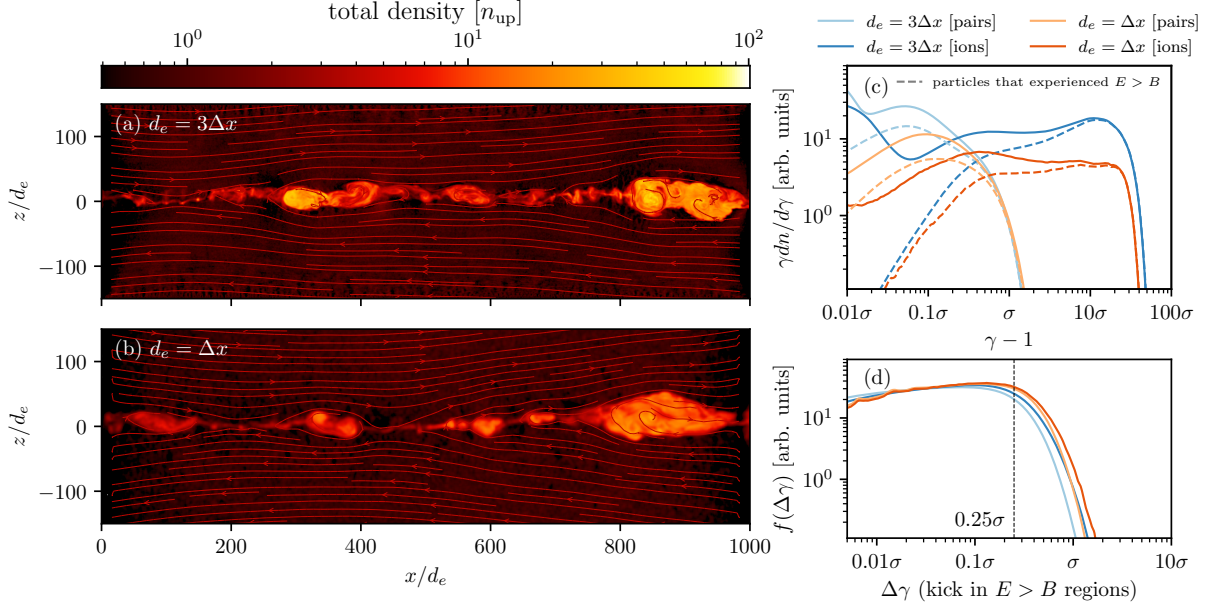


Figure A.2: Comparison of 2D plasma density slices (a, b) and normalized particle distributions (c, d) for different numerical resolutions $d_e/\Delta x = 1, 3$. Dark and bright lines correspond to ions and pairs, while blue and orange correspond to high- and low-resolution runs. The acceleration mechanisms, the statistics of the energy gains in $E > B$ regions, and thus the resulting distributions are largely insensitive to numerical resolution. The compression of plasmoids, on the other hand, is somewhat smaller for the low-resolution run.

both cases (Fig. A.2a, b), with the lower-resolution simulation showing less high-density regions.

The power-law slopes of the distributions, especially at energies $\gamma \gtrsim \sigma$, are also similar both for the high and the low-resolution runs. The energy gain in the regions of non-ideal field $E > B$ is identical for both simulations, showing that the dynamics around X-points is equally resolved for both simulations.

Appendix B: Appendices for Chapter 3

B.1 Additional Details on radiation and QED physics

The radiation reaction force of the curvature radiation and the corresponding energy losses are given by

$$mc \frac{d\mathbf{u}}{dt} = -\frac{2}{3}e^2 \frac{\gamma^4}{\rho^2} \boldsymbol{\beta}, \quad \frac{dE}{dt} \equiv -I = -\frac{2}{3}e^2 c \frac{\gamma^4}{\rho^2} \beta^2, \quad (\text{B.1})$$

where ρ is the curvature radius of the field line. The emitted photons have a peak energy of

$$\varepsilon_{\text{ph}}^* = \frac{3}{2} \hbar \frac{c}{\rho} \gamma^3, \quad (\text{B.2})$$

and the spectral emissivity is given by the synchrotron kernel ([Rybicki & Lightman, 1979](#))

$$\frac{dI(\varepsilon)}{d\varepsilon} = A \frac{\varepsilon}{\varepsilon_{\text{ph}}^*} \int_{\frac{\varepsilon}{\varepsilon_{\text{ph}}^*}}^{\infty} K_{5/3}(x) dx \quad (\text{B.3})$$

where $K_{5/3}(x)$ is the Macdonald function, and $\varepsilon_{\text{ph}}^*$ is given by Equation (B.2). The coefficient A is determined by equating the full emitted power, Eqn. (B.1), with the losses integrated over the

frequency spectrum. For the normalized photon energy, $\tilde{\varepsilon} = \varepsilon/m_e c^2$, the photon emission rate is

$$\frac{dN}{dt d\tilde{\varepsilon}} = \frac{1}{\sqrt{3}\pi} \frac{e^2 m_e c^2}{\hbar^2 c} \frac{1}{\gamma^2} \int_{\frac{\tilde{\varepsilon}}{\varepsilon_{\text{ph}}}}^{\infty} K_{5/3}(x) dx = \frac{1}{\sqrt{3}\pi} \frac{\alpha_F c}{\lambda_c} \frac{1}{\gamma^2} \int_{\frac{\tilde{\varepsilon}}{\varepsilon_{\text{ph}}}}^{\infty} K_{5/3}(x) dx, \quad (\text{B.4})$$

where $\alpha_F = e^2/\hbar c$ is the fine structure constant, and $\lambda_c = \hbar/m_e c$ is the reduced Compton wavelength.

To rescale the efficiency of curvature cooling, we choose the Lorentz factor of a particle, γ_{rad} , for which curvature radiation reaction force balances acceleration in a vacuum electric field on a scale of the polar cap, E_{PC} :

$$eE_{\text{PC}} = \frac{2}{3} e^2 \frac{\gamma_{\text{rad}}^4}{\rho_0^2}, \quad E_{\text{PC}} = 4\pi \rho_{\text{GJ}} R_{\text{PC}} \quad (\text{B.5})$$

where R_{PC} is the radius of the polar cap; and the vacuum electric field can be found as:

$$\nabla \cdot \mathbf{E} \equiv \frac{dE_z}{dz} = -4\pi e n_{\text{GJ}} \implies E_z(z) = -4\pi e n_{\text{GJ}} z. \quad (\text{B.6})$$

A convenient way to parameterize the size of the polar cap is to express it through the Lorentz factor reached by a particle in the full potential drop:

$$m_e c^2 \gamma_{\text{PC}} = 4\pi e \int_0^{R_{\text{PC}}} e n_{\text{GJ}} z dz = 4\pi e^2 n_{\text{GJ}} R_{\text{PC}}^2 / 2 \longrightarrow R_{\text{PC}}^2 = \frac{m_e c^2}{4\pi n_{\text{GJ}} e^2} 2\gamma_{\text{PC}} \equiv \left(d_e^{\text{GJ}} \sqrt{2\gamma_{\text{PC}}} \right)^2. \quad (\text{B.7})$$

Curvature radius of the field line, ρ_0 , can be defined as:

$$\rho_0 = R_{\text{PC}} \Theta = d_e^{\text{GJ}} \sqrt{2\gamma_{\text{PC}}} \Theta, \quad \Theta \gg 1. \quad (\text{B.8})$$

Another energy scale defines the Lorentz factor of a particle producing a photon of energy equal to the electron rest mass:

$$m_e c^2 \equiv \frac{3}{2} \hbar \frac{c}{\rho_0} \gamma_{\text{emit}}^3. \quad (\text{B.9})$$

These equations allow to parametrize the energies of curvature photons and the strength of the radiation reaction force. To summarize,

$$R_{\text{PC}} = d_{\text{e}}^{\text{GJ}} \sqrt{2\gamma_{\text{PC}}}, \quad F_{\text{curv}} = -e^2 n_{\text{GJ}} R_{\text{PC}} \left(\frac{\gamma}{\gamma_{\text{rad}}} \right)^4 \left(\frac{\rho}{\rho_0} \right)^{-2},$$

$$\rho_0 = R_{\text{PC}} \Theta, \quad \tilde{\varepsilon}_{\text{ph}}^* = \frac{\rho_0}{\rho} \left(\frac{\gamma}{\gamma_{\text{emit}}} \right)^3$$

With these renormalizations, the spectral photon emission rate (B.4) is

$$\begin{aligned} \frac{dN}{dt d\tilde{\varepsilon}} &= \frac{9\sqrt{3}}{8\pi} \frac{1}{\varepsilon_{\text{ph}}^*} I \frac{1}{\tilde{\varepsilon}_{\text{ph}}^*} \int_{\frac{\tilde{\varepsilon}}{\tilde{\varepsilon}_{\text{ph}}^*}}^{\infty} K_{5/3}(x) dx = \\ &= \frac{9\sqrt{3}}{8\pi} c \frac{e^2 n_{\text{GJ}} R_{\text{PC}}}{m_e c^2} \left(\frac{\gamma}{\gamma_{\text{rad}}} \right)^4 \left(\frac{\rho}{\rho_0} \right)^{-2} \frac{1}{\left(\frac{\gamma}{\gamma_{\text{emit}}} \right)^6 \left(\frac{\rho}{\rho_0} \right)^{-2}} \int_{\frac{\tilde{\varepsilon}}{\tilde{\varepsilon}_{\text{ph}}^*}}^{\infty} K_{5/3}(x) dx = \\ &= \frac{9\sqrt{3}}{4\sqrt{2}\pi} \frac{c R_{\text{PC}}}{(d_{\text{e}}^{\text{GJ}})^2} \frac{\gamma_{\text{emit}}^6}{\gamma_{\text{rad}}^4} \frac{1}{\gamma^2} \int_{\frac{\tilde{\varepsilon}}{\tilde{\varepsilon}_{\text{ph}}^*}}^{\infty} K_{5/3}(x) dx = \frac{9\sqrt{3}}{4\sqrt{2}\pi} \frac{c}{d_{\text{e}}^{\text{GJ}}} \frac{\gamma_{\text{PC}}^{1/2} \gamma_{\text{emit}}^6}{\gamma_{\text{rad}}^4} \frac{1}{\gamma^2} \int_{\frac{\tilde{\varepsilon}}{\tilde{\varepsilon}_{\text{ph}}^*}}^{\infty} K_{5/3}(x) dx \end{aligned} \quad (\text{B.10})$$

In the code, we employ tabulated values of $T(y_j) = \int_{y_j}^{\infty} K_{5/3}(x) dx$ with $y = \tilde{\varepsilon}/\tilde{\varepsilon}_{\text{rad}}^*$.

The differential cross section of the pair production reaction $\gamma + B \rightarrow e^- + e^+$ is given by (Erber, 1966):

$$d\sigma = 0.23 \frac{B}{B_q} \sin \psi \frac{\alpha_F}{\lambda_c} \exp \left(-\frac{8}{3 \frac{B}{B_q} \sin \psi \tilde{\varepsilon}_{\text{ph}}} \right) dl. \quad (\text{B.11})$$

Here, $B_q = m_e^2 c^3 / e \hbar \approx 4.41 \times 10^{13} \text{G}$ is the critical magnetic field, and ψ is the angle between the magnetic field and the direction of photon propagation. Since a photon is emitted along the field line and accumulates angle during the propagation, it can be estimated as $\psi(x) = \int_0^z dz / \rho(z)$,

where $\rho(z)$ is the curvature of the local field line. To renormalize the cross-section in the simulation, we need to express α_F/λ_c in the code units. Direct comparison of Equations (B.4) and (B.10) gives:

$$\begin{aligned} \frac{1}{\sqrt{3}\pi} \frac{\alpha_F c}{\lambda_c} \frac{1}{\gamma^2} \int_{\frac{\tilde{\varepsilon}}{\tilde{\varepsilon}_{\text{ph}}}}^{\infty} K_{5/3}(x) dx &= \frac{9\sqrt{3}}{4\sqrt{2}\pi} \frac{c}{d_e^{\text{GJ}}} \frac{\gamma_{\text{PC}}^{1/2} \gamma_{\text{emit}}^6}{\gamma_{\text{rad}}^4} \frac{1}{\gamma^2} \int_{\frac{\tilde{\varepsilon}}{\tilde{\varepsilon}_{\text{ph}}}}^{\infty} K_{5/3}(x) dx \\ &\Rightarrow \frac{\alpha_F}{\lambda_c} = \frac{27}{4\sqrt{2}} \frac{1}{d_e^{\text{GJ}}} \frac{\gamma_{\text{PC}}^{1/2} \gamma_{\text{emit}}^6}{\gamma_{\text{rad}}^4}. \end{aligned} \quad (\text{B.12})$$

Then, the resulting differential cross-section is

$$d\sigma = 0.23 \frac{27}{4\sqrt{2}} \frac{B}{B_q} \frac{1}{d_e^{\text{GJ}}} \frac{\gamma_{\text{PC}}^{1/2} \gamma_{\text{emit}}^6}{\gamma_{\text{rad}}^4} \sin \psi \exp \left(-\frac{8}{3 \frac{B}{B_q} \sin \psi \tilde{\varepsilon}_{\text{ph}}} \right) dl. \quad (\text{B.13})$$

We compute the total optical depth, $\sigma(l) = \int_0^l d\sigma$, during the propagation of a photon, as the angle between the magnetic field and the photon wave vector increases, $d\psi = c\Delta t/\rho(z)$, at every time step. When the optical thickness reaches unity, we simulate the photon decay into an electron-positron pair by removing the photon and adding an electron-positron pair at the same location, with the 4-velocity along the field line

$$u_{||} = \frac{|\cos \psi| (\tilde{\varepsilon}_{\text{ph}}^2 - 4)^{1/2}}{(\tilde{\varepsilon}_{\text{ph}}^2 \sin^2 \psi + 4 \cos^2 \psi)^{1/2}} \frac{(\mathbf{k} \cdot \mathbf{b})}{|\mathbf{k}| |\mathbf{b}|}, \quad (\text{B.14})$$

where $\mathbf{k}$ and $\mathbf{b}$ are unit vectors in the directions of the photon wave vector and the magnetic field, correspondingly.

In order to test both the plasma supply from the stellar atmosphere and the efficiency of curvature cooling, we employ a 1D charge-separated setup with different values of the magnetospheric currents, following (Timokhin & Arons, 2013). In Fig. B.1, we show tests for two

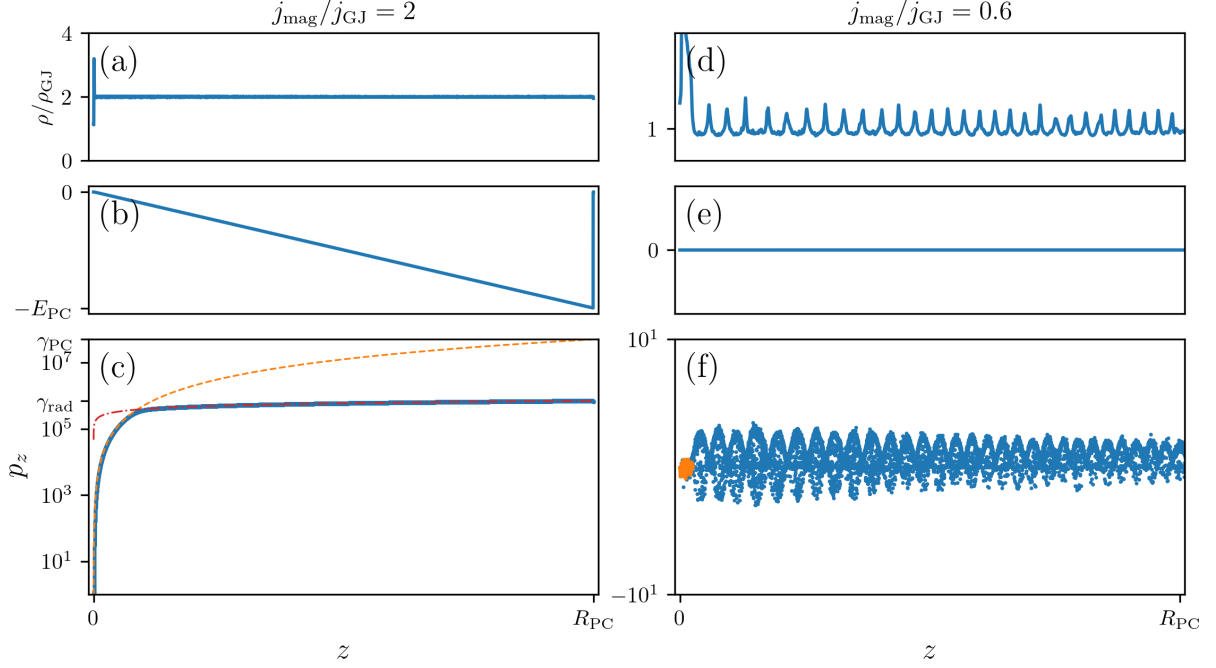


Figure B.1: 1D test of the curvature cooling and atmospheric particle injection in the charge-separated SCLF model: the left column shows (a) plasma density, (b) electric field, and (c) phase space, $z - p_z$ for the super-GJ current; and (d-g) - same quantities for the sub-GJ current. For the super-GJ current, the phase space plot, (c), shows that electrons initially accelerate in almost vacuum potential (the orange dotted line shows analytical prediction) and later enter into the radiation reaction limited regime (indicated by the red dot-dashed line). In the case of the sub-GJ current, extracted electrons travel at mildly relativistic speeds and maintain a state of small electric field.

different magnetospheric currents: sub-GJ, $j_{\text{mag}} = 0.6j_{\text{GJ}}$, and super-GJ, $j_{\text{mag}} = 2j_{\text{GJ}}$. In the case of the sub-GJ current, the magnetospheric current extracts the charge-separated flow of electrons with density $\approx \rho_{\text{GJ}}$, moving with sub-relativistic velocity (subplots (d) and (f), respectively), which leads to an efficient screening of the electric field (subplot (e) of Fig. B.1). In the case of the super-GJ current, the larger electron density is extracted, $n \approx |j_{\text{mag}}|/(ec)$, which leads to a strong electric field, $E_z \approx 4\pi(j_{\text{mag}}/c - \rho_{\text{GJ}})z$ (see subplots (a) and (b) in Fig. B.1) and quick electron acceleration. During the initial phases of acceleration, the cooling to curvature radiation is negligible, and the acceleration is determined by the electric field, $\gamma(z) = 2\pi|\rho_{\text{GJ}}|z^2$ (orange

dashed line in Fig. B.1c). When the Lorentz factor becomes comparable to γ_{rad} , the electron enters a radiation-reaction-limited regime when the radiation losses are balanced by an acceleration in the electric field $\gamma = \gamma_{\text{rad}}(E_z(z)/E_{\text{PC}})^{1/4} = \gamma_{\text{rad}}(z/R_{\text{PC}})^{1/4}$ (shown as red dashed-dotted line). We also verified that the total energy of the emitted photons is equal to the full radiative cooling losses experienced by the particle.

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