Particle Acceleration by Pickup Process Upstream of Relativistic Shocks

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Abstract

Particle acceleration at magnetized purely perpendicular relativistic shocks in electron–ion plasmas is studied by means of two-dimensional particle-in-cell simulations. Magnetized shocks with the upstream bulk Lorentz factor \( \gamma_{\text{i}} \gg 1 \) are known to emit intense electromagnetic waves from the shock front, which induce electrostatic plasma waves (wakefield) and transverse filamentary structures in the upstream region via stimulated/induced Raman scattering and filamentation instability, respectively. The wakefield and filaments inject a fraction of the incoming particles into a particle acceleration process, in which particles are once decoupled from the upstream bulk flow by the wakefield, and are picked up again by the flow. The picked-up particles are accelerated by the motional electric field. The maximum attainable Lorentz factor is estimated as \( \gamma_{\text{max, e}} \sim \alpha \gamma_{\text{i}}^3 \) for electrons and \( \gamma_{\text{max, i}} \sim (1 + m_e \gamma_{\text{i}}^2/m_i) \gamma_{\text{i}}^2 \) for ions, where \( \alpha \sim 10 \) is determined from our simulation results, \( \alpha \) can increase up to \( \gamma_{\text{i}} \) for a weakly magnetized shock if \( \gamma_{\text{i}} \) is sufficiently large. This result indicates that highly relativistic astrophysical shocks such as external shockss of gamma-ray bursts can be an efficient particle accelerator.

Unified Astronomy Thesaurus concepts: Plasma astrophysics (1261); Shocks (2086); High energy astrophysics (739); Cosmic rays (329)

1. Introduction

Particle acceleration is a ubiquitous physical process in the universe. The nonthermal emission spectra of high-energy astrophysical objects such as jets from active galactic nuclei (AGNs) and gamma-ray bursts (GRBs) are generally modeled as the synchrotron and inverse Compton emission of relativistic electrons (see, e.g., Piran 2005; Blandford et al. 2019). AGN jets and GRBs are often invoked for the source of ultra-high-energy cosmic rays (UHECRs) with energies beyond \( 10^{18} \) eV (e.g., Hillas 1984). Although the origin of UHECRs is still unknown, recent observations favor an extragalactic origin (Aab et al. 2018; Aartsen et al. 2018). Such astrophysical objects are usually associated with relativistic shocks as a consequence of the interaction between jets and the interstellar medium. Relativistic shocks are assumed to be an efficient particle accelerator.

Coherent emission of electromagnetic waves from the shock front is intrinsic to relativistic shocks, which has been confirmed by one-dimensional (1D; Langdon et al. 1988; Gallant et al. 1992; Hoshino et al. 1992; Amato & Arons 2006; Plotnikov & Sironi 2019), 2D (Iwamoto et al. 2017, 2018; Plotnikov et al. 2018; Babul & Sironi 2020), and 3D (Sironi et al. 2021) particle-in-cell (PIC) simulations in pair plasmas. It results from the synchrotron maser instability (SMI) in the shock transition, which is driven by electrons reflected off the shock-compressed magnetic field (Hoshino & Arons 1991). The excited electromagnetic waves whose group velocities are faster than the shock can propagate thorough the upstream plasmas as precursor waves. The SMI, which is called cyclotron maser instability in the weakly relativistic context, is also known as the emission mechanism of coherent radio sources such as auroral kilometric radiation at Earth and Jovian decametric radiation (see, e.g., Melrose 2017). Recently, some models of fast radio burst based on the coherent emission from relativistic shock via the SMI have been proposed (e.g., Lyubarsky 2014; Beloborodov 2017; Metzger et al. 2019; Beloborodov 2020; Margalit et al. 2020), and the SMI in the context of relativistic shocks attracts more attention from astrophysics. The precursor waves excited by the SMI can be strong enough to induce filamentation instability (FI), which is a transverse self-modulation of an intense electromagnetic wave (Kaw et al. 1973; Drake et al. 1974; Max et al. 1974; Sobacchi et al. 2020). Previous multidimensional simulations indeed demonstrated that the intense electromagnetic waves propagating through upstream plasma induce the transverse filamentary structures. The nonlinear effect of the electromagnetic waves plays a significant role in astrophysical plasmas (see also Lyubarsky 2018, 2019).

In electron–ion plasmas, the nonlinear interactions between the precursor waves and the upstream plasmas become more complicated. The stimulated/induced Raman scattering (SRS), which is the parametric decay of an intense electromagnetic wave into an electrostatic plasma wave such as a Langmuir wave, can work in addition to the FI (see, e.g., Krue 1988; Lyubarsky 2008). This plasma wave is conventionally called wakefield, and the concept of direct particle acceleration by the wakefield via the Landau resonance, which is the so-called wakefield acceleration (WFA), was first proposed in the study of laboratory plasmas (Tajima & Dawson 1979). The application of the WFA to UHECR acceleration was discussed later in the context of astrophysics (e.g., Chen et al. 2002; Arons 2003; Murase et al. 2009; Ebisuzaki & Tajima 2014, 2021) and laboratory plasmas (e.g., Kuramitsu et al. 2008, 2011a, 2011b, 2012; Liu et al. 2017, 2018, 2019). In relativistic shocks, Lyubarsky (2006) first found the wakefield excitation via the SRS using 1D PIC simulations, and recent 2D PIC simulations (Sironi & Spitkovsky 2011;
Ligorini et al. (2021a, 2021b) confirmed this. Furthermore, 1D PIC simulations by Hoshino (2008) demonstrated that nonthermal electrons and ions are generated in the upstream. Although our high-resolution 2D PIC simulations (Iwamoto et al. 2019) showed that this particle acceleration associated with the wakefield operates even in a 2D system, the detailed acceleration mechanism was not fully understood because some aspects are different from the standard WFA in laboratory plasmas.

In this work, we investigate the acceleration mechanism in more detail and show that the energetic particles are generated by a pickup process, where some incoming particles are decoupled once from the upstream bulk flow by the wakefield and are then accelerated by the motional electric field after they are picked up by the flow. In an ideal case, the maximum Lorentz factor may reach $\gamma_{\text{max},e} \sim \alpha \gamma_1^3$ for electrons and $\gamma_{\text{max},i} \sim (1 + m_e \gamma_1/m_i) \gamma_1^2$ for ions, where $\alpha \sim 10$ is the factor determined from our simulation results and $\gamma_1$ is the upstream bulk Lorentz factor. Although the observed Lorentz factor is smaller than this theoretical estimate due to the limitation of the simulation time, the partial trajectories of energetic particles are well described by the pickup process. This efficient acceleration process may operate in highly relativistic astrophysical shocks.

### 2. Simulation Setup

We perform 2D simulations of perpendicular relativistic shocks in electron-ion plasmas by using a fully kinetic electromagnetic PIC code (Matsumoto et al. 2013, 2015), which suppresses the numerical Cherenkov instability by choosing a magic CFL number and enables accurate and stable calculation (Ikeya & Matsumoto 2015). Our basic configuration is illustrated in Figure 1. We consider a rectangular computational domain in the $x$-$y$ plane with the periodic boundary condition applied in the $y$ direction. The number of grids in each direction is $N_x \times N_y = 80,000 \times 1536$. The number of particles per cell per species in the upstream and the grid size are set as $N_s \Delta x^2 = 64$ and $\Delta x/(c/\omega_{pe}) = 1/40$, respectively, which is motivated by the numerical convergence study of our previous simulations (see Iwamoto et al. 2017, Appendix A). $c$ is the speed of light, and $\omega_{pe}$ is the proper electron plasma frequency,

$$\omega_{pe} = \frac{4 \pi N_s e^2}{\gamma_1 m_e}.$$  \hspace{1cm} (1)

The time step is automatically determined as $\omega_{pe} \Delta t = 1/40$ because the magic CFL number is $c \Delta t / \Delta x = 1$ for our implicit Maxwell solver. Note that the implicit Maxwell solver is not restricted by the CFL number, and thus $c \Delta t / \Delta x = 1$ is numerically stable. A cold ion–electron flow with an ion-to-electron mass ratio $m_i/m_e = 50$ is injected from the right-hand boundary and propagating $-x$ direction with the bulk Lorentz factor $\gamma_1$. Our shock simulations are performed for the two cases $\gamma_1 = 40$ and 100. The incoming particles are reflected off at the left-hand conducting-wall boundary and trigger the shock propagating $+x$ direction. Our simulation frame corresponds to the downstream rest frame. We focus on purely perpendicular shocks, and the upstream ambient magnetic field $B_1$ is in the $z$ direction.

The basic structure and coherent emission of relativistic magnetized shocks are well characterized by the ratio of the Poynting flux to the upstream bulk kinetic energy flux,

$$\sigma_t = \frac{B_1^2}{4 \pi \gamma_1 N_s m_e c^2} = \frac{\omega_{pe}^2}{\omega_{ps}^2},$$  \hspace{1cm} (2)

where subscript $s = e, i$ represents the particle species, and $\omega_{cs}$ is the relativistic cyclotron frequency,

$$\omega_{cs} = \frac{eB_1}{\gamma_1 m_e}.$$  \hspace{1cm} (3)

We use fixed values of $\sigma_t = 0.1$ and $\sigma_e = (m_i/m_e) \sigma_t = 5$ throughout this study.

### 3. Shock Structure

Figure 2 shows the global shock structures at $\omega_{pe} \Delta t = 2000$ in the case of $\gamma_1 = 40$ (left) and 100 (right). The electron number density $N_e$, the ion number density $N_i$, the out-of-plane magnetic field $B_z$, the longitudinal electric field $E_z$, the longitudinal electric field averaged over the $y$ direction $\langle E_z \rangle$, and the phase-space densities in $x-u_{xs}$ and $x-u_{ys}$ for both electrons and ions are shown from top to bottom. All quantities are normalized by the corresponding upstream values, and the electron four velocity $u_{se} = \gamma_0 \beta_e$ is scaled by the mass ratio $m_e/m_i$. The global structures are similar to each other and show no clear $\gamma_1$ dependence.

Clear transverse density filaments are observed in the upstream region as in the case of pair plasmas (Iwamoto et al. 2017, 2018; Plotnikov et al. 2018; Babul & Sironi 2020; Sironi et al. 2021). One might think that the filaments are
attributed to the Weibel instability (Weibel 1959; Fried 1959). However, this cannot arise for such a high magnetization \( \sigma_i = 0.1 \) (see, e.g., Spitkovsky 2005; Sironi & Spitkovsky 2011; Sironi et al. 2013). We think that the large-amplitude precursor waves, which are clearly seen in \( B_z \), induce the FI (Kaw et al. 1973; Drake et al. 1974; Max et al. 1974; Sobacchi et al. 2020) and create the filaments. This requires intense electromagnetic pump waves. The precursor wave amplitude is large in the sense that the wave strength parameter \( a \) is greater than unity (Iwamoto et al. 2017; Plotnikov & Sironi 2019),

\[
a = \frac{\delta E}{m_e c \omega} \sim \gamma_1 \sqrt{\varepsilon_p} > 1, \tag{4}
\]

for \( \gamma_1 \gg 1 \). Here \( \delta E \) is the wave amplitude, \( \omega \) is the wave frequency, and \( \varepsilon_p = b B^2 / 4 \pi \gamma N n_e c^2 \) is the normalized precursor wave energy. \( \varepsilon_p \sim 1 \) for \( \sigma_i = 0.1 \) (Iwamoto et al. 2019) and thus \( a \sim \gamma_1 \gg 1 \). Therefore, the precursor waves are subject to the FI. The FI is a nonlinear wave-wave interaction and a kind of the parametric decay instability. An intense electromagnetic pump wave with the wavenumber \( \vec{k}_0 \) parametrically decays into the wakefield with \( \vec{k}_1 \) and a scattered electromagnetic wave with \( \vec{k}_2 = \vec{k}_0 - \vec{k}_1 \) in the linear phase of the SRS (Hoshino 2008). The sinusoidal wakefield in the far upstream region is induced in this linear phase, and particles are merely oscillating, as can be seen in the phase-space plots. The wakefield gradually becomes turbulent, indicating that the SRS enters the nonlinear phase. Both electrons and ions are strongly accelerated/heated near the shock front. Note that the Lorentz transformation from the proper frame into the simulation frame increases the thermal spread only in the \( u_x \) direction. The phase-space plots \( x-u_x \) show that electrons (ions) are preferentially accelerated in the \( +y \) (\(-y\)) direction. This is because particles are mainly

![Figure 2. Shock structures at \( \omega_p t = 2000 \) with \( \gamma_1 = 40 \) (left) and 100 (right). The electron number density \( N_e \), the ion number density \( N_i \), the \( z \) component of the magnetic field \( B_z \), the \( x \) component of the electric field \( E_x \), the transversely averaged electric field \( \langle E_x \rangle \), and the phase-space densities in the \( x-u_x \) and \( x-u_y \) are shown.](image_url)
accelerated by the motional electric field $E_y = -\beta_y B_1$, as discussed in Section 5. These features are consistent with the previous 1D simulation by Hoshino (2008).

4. Particle Energy Spectra

Figure 3 shows energy spectra of electrons (blue) and ions (red) for $\gamma = 40$ (top panels) and $\gamma = 100$ (bottom panels). The downstream spectra (left panels), the upstream spectra (middle panels), and the upstream spectra measured in the proper frame (right panels) are shown.

The electron Lorentz factor is scaled by the mass ratio $m_e/m_i$. The spectra (a), (b), (d), and (e) are measured in the simulation frame. No clear nonthermal tail is seen in the downstream energy spectra (a) and (d) in the range $700 \leq x/(c/\omega_{pe}) \leq 1050$. As can be seen in the near-upstream energy spectra (b) and (e) in the range $1120 \leq x/(c/\omega_{pe}) \leq 1200$, an energy equipartition between electrons and ions is already achieved in the upstream region due to the electron–ion coupling via the wakefield (Lyubarsky 2006; Hoshino 2008; Iwamoto et al. 2012). Typical nonthermal and thermal electrons are shown as red and gray lines, respectively. The thermal electron is merely oscillating within the wakefield. It finally enters the shock at $x/(c/\omega_{pe}) \sim 1080$ and gyrates in the downstream region. The bulk Lorentz factor of the thermal electrons has its maximum when energy equipartition between ions and electrons is achieved (Lyubarsky 2006; Hoshino 2008; Iwamoto et al. 2019), and thus it is written as

$$\gamma_{\text{max},e} \sim \frac{1}{2} \frac{m_i}{m_e} \gamma_1.$$  

As shown by the gray line in Figure 4, the normalized maximum Lorentz factor is $m_e \gamma_e/m_i \gamma_1 \sim 1$. The difference is at most a factor of 2, and this result is roughly consistent with the above estimate. However, the maximum Lorentz factor of the nonthermal electron is much larger and cannot be explained by the ion–electron coupling. Although the nonthermal electron is also oscillating in the region $1230 \leq x/(c/\omega_{pe}) \leq 1400$, it begins to travel along the $+y$ direction at $x/(c/\omega_{pe}) \sim 1230$ and seems to gain energy from the motional electric field.
$E_y = -\beta_1 B_1$. This particle acceleration continues until the electron enters the shock at $x/(c/\omega_{pe}) \sim 1080$. Its maximum Lorentz factor $\gamma_{\text{max},e}$ can be calculated as

$$\gamma_{\text{max},e} = \frac{e \beta_1 B_1 \Delta y}{m_e c^2} = \gamma_1 \beta_1 \sqrt{\omega_{ce}} \frac{\Delta y}{c/\omega_{pe}}. \quad (7)$$

We measured $\Delta y/(c/\omega_{pe}) \sim 85$ in the region $1080 \lesssim x/(c/\omega_{pe}) \lesssim 1230$, and thus $m_e \gamma_{\text{max},e}/m_1 \gamma_1 \sim 4$, showing a good agreement with our simulation result. This acceleration process is identical to the pickup process in space physics (e.g., Möbius et al. 1985; Oka et al. 2002). In the heliosphere, some neutrals are ionized mainly through charge exchange with solar wind protons and are picked up by the solar wind electric field. The pickup ions are efficiently accelerated by the motional electric field. The pickup process including the relativistic effect can be theoretically analyzed using the relativistic equations of motion. We have neglected the precursor waves, wakefields, and filaments and only considered the ambient magnetic field $B_z = B_1$ and the motional electric field $E_y = -\beta_1 B_1$. The analytical solution in the simulation frame is written (see Appendix A for the detailed calculations)

$$\gamma_1 = \gamma_{01}^2 \left[ 1 + \beta_1 \beta_0 - \beta_1 (1 + \beta_0) \cos \theta_1 \right], \quad (8)$$

$$x_s = x_0 - c \beta_1 t + \gamma_{01}(\beta_1 + \beta_0) \frac{c}{\omega_{ce}} \sin \theta_1, \quad (9)$$

$$y_s = y_0 + \gamma_{01} \left[ \beta_1 + \beta_0 \right] \frac{c}{\omega_{ce}} (1 - \cos \theta_1), \quad (10)$$

$$\theta_1 = \omega_{ce} \left[ t + \beta_1 (x_0 - x_0)/c \right] \gamma_{01}(1 + \beta_1 \beta_0), \quad (11)$$

where the subscript 0 indicates the initial quantities at the time when the particles are picked up by the upstream bulk flow. The positive (negative) sign in Equation (10) corresponds to electrons (ions). Here we assume $\beta_{0i} = +\beta_{0b} \hat{x}$. Note that the upstream bulk flow propagates toward the $-x$ direction and $\omega_{ce}$ is the unsigned cyclotron frequency. We determined the initial quantities $x_{0i}$, $y_{0i}$, and $\beta_{0i}$ from our simulations, and the theoretical solutions are shown in Figure 4 by the dashed black lines. The simulation results agree well with the theoretical trajectories for the region $1080 \lesssim x/(c/\omega_{pe}) \lesssim 1230$, indicating that the nonthermal electron is picked up by the bulk flow. The electron reaches the shock front at $x/(c/\omega_{pe}) \sim 1080$ and gradually deviates from the theoretical trajectories. According to Equation (8), the maximum Lorentz factor is estimated as

$$\gamma_{\text{max},i} \sim \gamma_{01}^2 \gamma_{0i}. \quad (12)$$

Here we have neglected the factors on the order of unity.

The above analytical solution indicates that the nonthermal particles must travel in the opposite direction of the bulk flow before entering into the pickup process. Sironi & Spitkovsky (2011) reported the same acceleration process and pointed out that these particles are decoupled from the bulk flow, and thus they can feel the motional electric field due to the velocity difference. Equation (8) indeed demonstrates that the particle acceleration does not occur ($\gamma_1 \sim \gamma_1$) for a particle moving with the same velocity as the bulk flow $\beta_{0i} \sim -\beta_1$. Figure 5 shows the trajectories of nonthermal electrons measured in the simulation frame as gray lines. The color map represents the magnetic field $\vec{B}$ at $\omega_{pe} t = 1800$, which satisfies

$$\vec{E}_y = -\beta_1 B_1 - \frac{\partial \phi}{\partial y}, \quad (13)$$

$$\vec{E}_y + \beta_1 \vec{B}_z = 0. \quad (14)$$

The electrostatic potential $\phi$ is calculated from the snapshot at $\omega_{pe} t = 1800$ by performing the Helmholtz decomposition. We have removed the electromagnetic fields arising from the
precursor waves because these superluminal waves are not responsible for the resonant wave-particle interaction and do not directly contribute to the particle acceleration. The trajectories in Figure 5 indeed demonstrate that nonthermal electrons propagate toward the +x direction before they are picked up by the bulk flow. The nonthermal electrons enter the weakly magnetized region resulting from the FI and then enter the acceleration phase, indicating that the filaments trigger the pickup process.

Figure 6 shows the time evolution of the typical nonthermal electron measured in the simulation frame. We take the moving average for the time period $\omega_{pe}\Delta t = 5$, which is motivated by the typical frequency of the precursor waves $\omega/\omega_{pe} \sim 2-5$ (Iwamoto et al. 2017, 2018), to remove the effect of the precursor waves. The top panel displays the energy gain $\Delta \gamma_e = \gamma_e - \gamma_1$ (red) and work done by $E_x$ (green) and $E_y$ (blue) normalized by $m_1 \gamma_1/m_e$.

$$\Delta \gamma_k = -\frac{e}{m_e c^2} \int_0^t E_k v_k dt,$$

(15)

where $k=x, y$ and $\omega_{pe}t_0 = 1500$. The incoming electron is decelerated by the wakefield and begins to gyrate. Then it loses its energy due to both the wakefield and the motional electric field for $1750 \lesssim \omega_{pe}t \lesssim 1795$. $\Delta \gamma_x$ increases in time for $1795 \lesssim \omega_{pe}t \lesssim 1810$, whereas $\Delta \gamma_y$ is almost constant, showing that the energy gain during the corresponding time period originates from the wakefield. After $\omega_{pe}t \approx 1810$, $\Delta \gamma_y$ becomes dominant, and thus the nonthermal electron enters the pickup process. The $x$ and $y$ components of the electron velocity normalized by the speed of light, $\beta_x$ (green) and $\beta_y$ (blue), are shown in the middle panel. $\beta_y$ is positive for $1795 \lesssim \omega_{pe}t \lesssim 1810$, and the nonthermal electron moves with the relativistic velocity $\beta_y \sim 1$ in the same direction as the wakefield propagation. Note that the phase velocity of the wakefield is almost equal to the speed of light (Hoshino 2008). $B_z \sim 0$ is satisfied inside the filaments, and the electrostatic force $-eE_x$ easily overcomes the Lorentz force $-e\beta_z B_z$ despite $E_{wake}/B_1 < 1$. Therefore, the electron is trapped by the wakefield and accelerated via the Landau resonance. The acceleration continues until the Lorentz force exceeds the electrostatic force. The bottom panel of Figure 6 shows the total force $-e(E_x + \beta_z B_z)$ (red), electrostatic force $-eE_x$ (green), and Lorentz force $-e\beta_z B_z$ (blue) normalized by $eB_1$ at the electron position. The electrostatic force indeed dominates the Lorentz force for $1795 \lesssim \omega_{pe}t \lesssim 1810$, and then the total force is controlled by the Lorentz force after $\omega_{pe}t \approx 1810$. One may think that this process is similar to the shock surfing acceleration (SSA; Shimada & Hoshino 2000; Hoshino & Shimada 2002), in which the electrostatic waves trap the incoming electrons and the motional electric field accelerates them during the multiple reflection within the electrostatic waves. The essential difference is that the electrons gain the energies from the wakefield during the trapping because the motional electric field vanishes within the filaments $B_z \sim 0$. This acceleration mechanism is analogous to the standard WFA in laboratory plasmas rather than the SSA.

We have confirmed our idea that the electrons that are preaccelerated by the wakefield inside the filaments are further accelerated via the pickup process by performing the test-particle simulations (see Appendix B).

We evaluate the initial Lorentz factor $\gamma_0$. Since the wakefield directly accelerates the nonthermal electron within the filaments, the energy gain is expressed as

$$\Delta \gamma_x = \frac{eE_{wake}L_{acc,e}}{m_e c^2} = \gamma_1 \sqrt{\sigma_e} \frac{E_{wake} L_{acc,e}}{B_1} c/\omega_{pe},$$

(16)

where $E_{wake}$ is the wakefield amplitude and $L_{acc,e}$ is the acceleration length. $E_{wake}$ can be estimated as (Hoshino 2008)

$$\frac{E_{wake}}{B_1} = \frac{1}{\gamma_1 \sqrt{\sigma_e}} \frac{\eta \mu^2}{\sqrt{1 + \eta^2}} \sim \sqrt{\sigma_e},$$

(17)

where $\eta$ represents the wave polarization: $\eta = 1$ for circular polarization and $\eta = 1/2$ for linear polarization. Here we have used Equation (4) and $\eta = 1/2$, and neglected factors on the order of unity. By substituting $\sigma_e \sim 1$ and $\sigma_e = 5$ into Equation (17), we obtain $E_{wake}/B_1 \sim O(10^{-1})$, which agrees with our simulation results (see Figure 2). The acceleration
length corresponds to the wakefield wavelength for the WFA in laboratory plasmas. In our shock simulations, however, the acceleration length is limited by the size of the unmagnetized region, which is much smaller than the wakefield wavelength. We evaluate $L_{\text{acc},e}$ from the nonthermal electron trajectories in Figure 5.

$$L_{\text{acc},e} \sim 10.$$  \hfill (18)

By substituting Equations (17) and (18) into Equation (16), we have the estimate of the initial Lorentz factor $\gamma_{0e}$,

$$\gamma_{0e} \approx 1 + \Delta \gamma_e \sim \alpha \gamma_1 \sqrt{c_p},$$  \hfill (19)

where

$$L_{\text{acc},e} \sim 10.$$  \hfill (20)

This shows that a highly relativistic shock $\gamma_1 \gg 1$ can be an efficient particle accelerator.

Figure 6. Time evolution of a nonthermal electron. Top panel: energy gain $\Delta \gamma_e = \gamma_e - \gamma_1$ (red) and work done by $E_x$ (green) and $E_y$ (blue). Middle panel: three velocity $\beta_x$ (green) and $\beta_y$ (blue). Bottom panel: total force $-e(E_x + \beta_y B_z)$ (red), electrostatic force $-eE_x$ (green), and Lorentz force $-e\beta_y B_z$ (blue) at the electron position.

Equation (21) shows the electron maximum Lorentz factor is $\gamma_{\text{max},e} \sim 10^4$ for $\gamma_1 = 40$ and $\gamma_{\text{max},e} \sim 10^6$ for $\gamma_1 = 100$. However, the electron energy spectra in Figure 3 demonstrate that the maximum Lorentz factor in our simulations is much smaller than we expect. Although the maximum Lorentz factor for $\gamma_1 = 100$ is larger than that for $\gamma_1 = 40$, the difference is at most a factor of 2. The deviation from the analytical estimate can be explained as follows. In the simulation frame, the picked-up electrons propagate toward the $-x$ direction while accelerated by the motional electric field. If they are picked up near the shock front, they enter the shock soon and the acceleration ceases before they take the maximum Lorentz factor. Equation (8) shows that the Lorentz factor takes its maximum when $\theta_e = \pi$. Equation (9) reduces to

$$x_e = x_{0e} - c \beta_{\text{acc},e}.$$  \hfill (22)
This estimate shows $\Delta t_{acc,e}$ is the acceleration timescale. By substituting this and $\theta_c = \pi$ into Equation (11), we can evaluate $t_{acc,e}$:

$$\omega_{ce} t_{acc,e} \sim 2\pi \gamma_1^3 \gamma_{0e} \sim 2\pi \alpha \gamma_1^3.$$  

(23)

We finally obtain the moving distance of the electron in the $x$ direction $\Delta x_e = |x_e - x_{0e}|$ during the time period $\Delta t = t_{acc,e}$.

$$\frac{\Delta x_e}{c/\omega_{pe}} = \beta_1 \omega_{pe} t_{acc,e} \sim \frac{2\pi \alpha \gamma_1^3}{\sqrt{\varepsilon_e}}.$$  

(24)

This estimate shows $\Delta x_e/(c/\omega_{pe}) \sim 10^6$ for $\gamma_1 = 40$ and $\Delta x_e/(c/\omega_{pe}) \sim 10^7$ for $\gamma_1 = 100$, which are much larger than the precursor wave region: $1100 \lesssim x/(c/\omega_{pe}) \lesssim 1800$ in the final state of our simulations. Therefore, the picked-up electrons enter the shock before they obtain the theoretical maximum Lorentz factor. Since the group velocity of the precursor wave $v_g \sim c$ is faster than the shock propagation velocity $v_{sh} \sim (1/2 + 3\sigma_z/4)c \sim 0.575c$, the precursor wave region becomes larger as time passes. In the later phase, the incoming electrons can be picked up far away from the shock front and are sufficiently accelerated by the motional electric field before entering the shock. We thus think that the observed Lorentz factor will be closer to the theoretical one if we follow the long-term evolution. Our test-particle simulations indeed demonstrate that the maximum Lorentz factor is consistent with the above estimate (see Appendix B). These results indicate that the particle energy spectra do not yet reach steady state. Nonthermal tails might be observed downstream in the later phase.

5.2. Ion Acceleration

The ion acceleration can be explained by the pickup process as well. Figure 7 shows the nonthermal and thermal ion trajectories in the same format as Figure 4. The thermal ion is oscillating inside the wakefield, whereas the nonthermal ion is accelerated by the motional electric field. The analytical solutions of the pickup process (Equations (8), (9), (10), and (11)) are shown in black and agree well with the simulation results.

We here discuss how ions are injected into the pickup process. Figure 8 shows the trajectories of nonthermal (gray) and thermal (green) ions with the magnetic field $B_z$, which is determined from the snapshot at $\omega_{pe} t = 1660$ in the same manner as for the electron. The ion injection occurs in the highly magnetized region $B_z/B_1 > 1$, unlike for the electron. The incoming cold ions are gradually thermalized by the SRS and/or the FI, which is clearly seen in the phase-space density plots of Figure 2. The thermalized ions can slightly deviate from the bulk motion, and an $E \times B$ drift can be induced. The ion trajectory is thus given by the cycloid, which is the case for the thermal ions (green lines). On the other hand, the nonthermal ions (gray lines) are suddenly reflected toward the $+x$ direction during the cycloid motion and are then picked up by the bulk flow. This kick toward upstream seems to trigger the pickup process.

Figure 9 displays the time evolution of the typical nonthermal (left) and thermal (right) ions in the same format as Figure 6. We take the moving average for the time period $\omega_{pe} \Delta t = 5$ as well. The energy gain $\Delta \gamma_i = \gamma_i - \gamma_1$ (red) and work done by $E_z$ (green) and $E_y$ (blue) normalized by $\gamma_1$ are shown in the top panels. In the case of the nonthermal ions (left), both the wakefield and the motional electric field contribute to the energy loss for $\omega_{pe} \Delta t \lesssim 1669$. $\Delta \gamma_i$ increases in time after $\omega_{pe} \Delta t \sim 1669$ and $\Delta \gamma_y$ exhibits the same tendency, indicating that the ion enters the pickup process at $\omega_{pe} \Delta t \sim 1669$. As can be seen in the left middle panel, at $\omega_{pe} \Delta t \sim 1669$, $\beta_z$ becomes positive and the ion decoupling occurs. The bottom left panel of Figure 6 demonstrates that the electrostatic force $eE_z$ (green) exceeds the Lorentz force $\gamma_i \beta_i B_z$ (blue) for $1668 \lesssim \omega_{pe} \Delta t \lesssim 1669$, and the wakefield can reflect the incoming
ion. We think that the kick imparted by the wakefield determines whether the ion enters the pickup process or not. The drifting ions can satisfy $\beta_y \sim 0$ at some point on the way to the shock. If the wakefield kicks them at the time when $\beta_y \sim 0$ is satisfied, the electrostatic force $eE_x$ can easily overcome the Lorentz force $e\beta_y B_z \sim 0$. Furthermore, the Lorentz factor of the drifting ions has the minimum value when $\beta_y \sim 0$ and they are relatively subject to the wakefield. In fact, the thermal ion (right) shows that the $eE_x$ is negative at the time $\omega_{pe}t \sim 1670$ when $\beta_y \sim 0$ is satisfied and the wakefield cannot reflect it. Although the increase in $\Delta \gamma_x$ is barely visible for $1668 \lesssim \omega_{pe}t \lesssim 1669$ in the top left panel, we think that the finite kick imparted by the wakefield is responsible for the decoupling.

The ion injection into the pickup process seems to be different from the electron. This may be attributed to the mass difference like the standard WFA in laboratory plasmas. Electrons are relatively easily accelerated by the wakefield, whose phase velocity is almost equal to the speed of light via the Landau resonance, whereas ions have difficulty with the resonance due to the high mass. Our simulations indeed demonstrate the efficient electron WFA inside the filaments. Although the ion WFA is transient and inefficient, we think that this finite preacceleration injects ions into the pickup process.

We here evaluate the ion maximum Lorentz factor $\gamma_{\text{max},i}$ in the same manner as for the electron. Since the incoming ion is reflected by the wakefield, $\Delta \gamma_x$ can be written as

$$
\Delta \gamma_x = \frac{eE_{\text{wake}} \omega_{\text{acc},i}}{m_i c^2} \sim \frac{m_i}{m_e} \frac{\sqrt{\rho}}{c/\omega_{pe}} L_{\text{acc},i}.
$$

Figure 8. Trajectories of nonthermal (gray) and thermal (green) ions with $\vec{B}_i$ at $\omega_{pe}t = 1660$.

Figure 9. Time evolution of a nonthermal (left) and thermal (right) ion in the same format as Figure 6.
The nonthermal ion trajectories in Figure 8 indicate

\[ \frac{L_{\text{acc},i}}{c/\omega_{pe}} \sim 1. \]  

(26)

\( \gamma_{0i} \) can be estimated as

\[ \gamma_{0i} \sim 1 + \Delta \gamma_{\text{s}} \sim 1 + \frac{m_e}{m_i} \gamma_{\text{1}} \sqrt{\epsilon_{p}} \]  

(27)

The ion maximum Lorentz factor \( \gamma_{\text{max},i} \) can be derived from Equation (12),

\[ \gamma_{\text{max},i} \sim \left( 1 + \frac{m_e}{m_i} \gamma_{\text{1}} \sqrt{\epsilon_{p}} \right)^2 \sim \left( 1 + \frac{m_e}{m_i} \gamma_{\text{1}} \right)^2 \]  

(28)

Here we have used \( \epsilon_{p} \sim 1 \). As can be seen in Figure 3, the ion energy spectra show a smaller maximum Lorentz factor than this estimate. We can estimate the moving distance in the \( x \) direction during the acceleration as in the case of the electron. The acceleration timescale \( t_{\text{acc},i} \) is expressed as

\[ \omega_{ci} t_{\text{acc},i} \sim 2 \pi \left( 1 + \frac{m_e}{m_i} \gamma_{\text{1}} \right)^2. \]  

(29)

We obtain the moving distance of the ion in the \( x \) direction

\[ \Delta x_i = |x_i - x_{0i}|, \]

\[ \frac{\Delta x_i}{c/\omega_{pe}} = \beta_{\text{1}} \omega_{pe} t_{\text{acc},i} \sim \frac{2 \pi (m_i/m_e + \gamma_{\text{1}}) \gamma_{\text{1}}^2}{\sqrt{\sigma_{e}}}. \]  

(30)

This estimate gives \( \Delta x_i/(c/\omega_{pe}) \sim 10^5 \) for \( \gamma_{\text{1}} = 40 \) and \( \sim 10^6 \) for \( \gamma_{\text{1}} = 100 \), and thus the pickup ions enter the shock before they reach the maximum Lorentz factor in our simulations. Nonthermal ions as well as electrons might be seen in the downstream in the later phase.

### 6. Discussion

In this work, we assumed a precursor wave power \( \epsilon_{p} \sim 1 \), which is valid for \( \sigma_{i} \sim 0.1–1 \) (Iwamoto et al. 2019). \( \epsilon_{p} \) is independent of \( \gamma_{\text{1}} \) as long as \( \gamma_{\text{1}} \gg 1 \) (Plotnikov & Sironi 2019), and it is mainly controlled by \( \sigma_{i} \) due to the ion–electron coupling (Lyubarsky 2006; Hoshino 2008; Iwamoto et al. 2019). Although the \( \sigma_{i} \) dependence is not fully understood, previous PIC simulations demonstrated that \( \epsilon_{p} \) is convex upward as a function of \( \sigma_{i} \) and takes the maximum value \( \epsilon_{p} \sim 1 \) at \( \sigma_{i} \sim 0.1 \). The ion acceleration efficiency (Equation (28)) is not strongly dependent on \( \epsilon_{p} \) as long as \( \gamma_{\text{1}} < m_i/m_e \) and...
\[
\gamma_{\text{max, }i} \sim \gamma_1^3
\]
for \( \epsilon_p \ll 1 \). On the other hand, the electron acceleration efficiency (Equation (21)) drastically deteriorates for \( \epsilon_p \ll 1 \) and may be reduced to \( \gamma_{\text{max, }e} \sim \gamma_{\text{max, }i} \sim \gamma_1^2 \). For low \( \sigma_i \), however, the electron acceleration is not necessarily less efficient. The acceleration length \( L_{\text{acc, }e} \) may be much greater because \( E_{\text{wake}}/B_1 \sim \delta B/B_1 > 1 \) can be satisfied for weakly...
magnetized plasmas and the electrostatic force can easily exceed the Lorentz force. The acceleration length may be comparable to the wakefield wavelength (Krueer 1988; Hoshiba 2008),

$$\alpha = \frac{L_{\text{acc},e}}{c/\omega_{pe}} \sim \gamma^4.\text{e}$$

(31)

The maximum Lorentz factor is expressed as

$$\gamma_{\text{max},e} \sim \gamma^4 \sqrt{e^p}.\text{e}$$

(32)

The Weibel instability develops for $\sigma_i \ll 1$, and the wave power declines because the ring-like momentum distribution in the shock transition that is essential for the SMI is strongly modified by the Weibel-generated magnetic field (Sironi & Spitkovsky 2011; Iwamoto et al. 2017, 2018). However, Equation (32) exhibits a weaker dependence on $\epsilon_p$ than $\gamma_1$. We thus speculate that the efficient electron acceleration occurs as long as $\gamma_1 \gg 1$.

The upstream temperature has an influence on the acceleration efficiency as well. Since the accelerated/heated particles enter the shock, $\epsilon_p$ may decrease in time due to the suppression of the higher-order harmonics (Amato & Arons 2006). In pair plasma, Babul & Sironi (2020) reported that the wave emission efficiency declines by almost two orders of magnitude for the thermal spread $k_B T_e/m_e c^2 \gtrsim 10^{-1}$. Although the temperature dependence in ion-electron plasmas remains unsolved, $\epsilon_p$ probably shows a similar tendency. The precursor wave emission might cease and the size of the precursor wave region might be insufficient to accelerate the incoming particles up to the theoretical estimate even if we follow the long-term evolution. The particle energy spectra in the final state are an open question.

Both electrons and ions are accelerated via the pickup process in the upstream. We speculate that the pickup process provides seed particles for other acceleration mechanisms such as Fermi acceleration. The preaccelerated particles in the upstream may be further accelerated, and power-law spectra may be generated in the downstream.

The preexisting cosmic rays may be reaccelerated by the pickup process. These energetic protons can diffuse far upstream. Since they are decoupled from the upstream bulk flow, the pickup process can work. According to Equation (12), they can be reaccelerated by a factor of $\gamma^2$. This reacceleration process may operate repeatedly, and they may be accelerated up to the UHECR energy range.

7. Summary

We investigated the particle acceleration in relativistic ion–electron shocks by 2D PIC simulations. The particle energy spectra in the upstream show nonthermal tails for both electrons and ions. We found that they are mainly accelerated by the motional electric field. This particle acceleration is well described by the pickup process, in which particles are once decoupled from the upstream bulk flow by the wakefield and are picked up again by the flow. We estimated a maximum Lorentz factor $\gamma_{\text{max},e} \sim \alpha \gamma^4$ for the electron and $\gamma_{\text{max},i} \sim (1 + m_e \gamma_1/m_i) \gamma_i^2$ for the ion, where $\alpha \sim 10$ is the normalized acceleration length and is determined from our simulations. Since this acceleration requires a large computational domain, we were unable to follow the whole acceleration process because our computational resources were limited. The accelerated particles might exhibit power-law-like spectra in the downstream at a later phase. The pickup process may play a significant role for particle acceleration in highly relativistic shocks $\gamma_1 \gg 1$, such as external shocks of GRBs.

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Appendix A

Analytical Solutions of Pickup Process

We here derive the analytical solutions of the pickup process. Let us assume a charged particle in the background magnetic field $B_i = B_1$ and a motional electric field $E_i = -\beta_i B_1$. The basic equations are the relativistic equations of motion,

\[ m_i c \frac{d\mathbf{u}_i}{dt} = q_i \beta_i B_1, \quad (A1) \]

\[ m_i c \frac{d\mathbf{u}_i}{dt'} = -q_i (\beta_{s} \mathbf{e} + \beta_i) B_1. \quad (A2) \]

Here $\mathbf{u}_s = \gamma_i \mathbf{\beta}_s$ is the four velocity, $q_e = -e$ is the electron charge, and $q_i = +e$ is the ion charge. By performing a Lorentz transformation from the simulation frame into the plasma rest frame, Equations (A1) and (A2) reduce to

\[ m_i c \frac{d\mathbf{u}_i}{dt'} = q_i \beta_{s} B_1', \quad (A3) \]

\[ m_i c \frac{d\mathbf{u}_i}{dt'} = -q_i \beta_{s} B_1'. \quad (A4) \]

where the prime indicates the physical quantities in the plasma rest frame. The motional electric field vanishes in the plasma rest frame, and thus the kinetic energy is conserved,

\[ \gamma_1' = \text{const.} = \gamma_1 \gamma_0 (1 + \beta_i \beta_{0,\gamma}), \quad (A5) \]

where $\gamma_0 = 1/\sqrt{1 - \beta_{0,\gamma}^2}$ is the initial Lorentz factor at $t = 0$. Here we have assumed the initial three velocity $\beta_{0,\gamma} = \beta_{0,\gamma} \hat{x}$ and performed the Lorentz transformation $\gamma_1' = \gamma_1 \gamma_1 (1 + \beta_i \beta_{0,\gamma})$. Equations (A3) and (A4) describe the gyromotion around the background magnetic field $B_1' = B_1/\gamma_1$ and are easily solved,

\[ u_{s,xx} = u_{0,s} \cos \theta', \quad (A6) \]

\[ u_{s,yy} = \pm u_{0,s} \sin \theta', \quad (A7) \]

\[ x_1' = x_0 + u_{0,s} c \frac{\omega_{\omega_{s}}}{\omega_{\omega_{s}}} \sin \theta', \quad (A8) \]

\[ y_1' = y_0' \pm u_{0,s} c \frac{\omega_{\omega_{s}}}{\omega_{\omega_{s}}} (1 - \cos \theta'), \quad (A9) \]

\[ \theta' = \frac{\omega_{\omega_{s}} t'}{\gamma_1}, \quad (A10) \]
\[ u_{\text{fl}}' = \gamma_1 \gamma_0 (\beta_1 + \beta_0), \]  

where the subscript 0 represents the initial quantities at \( t = 0 \).

The positive (negative) sign corresponds to the electron (ion).

Note that \( \omega_{cs} = e B_1 / \gamma_1 m_e c \) is the unsigned cyclotron frequency.

By performing Lorentz transformation from the plasma rest frame into the simulation frame, we obtain the exact solutions of Equations (A1) and (A2),

\[ \gamma_\alpha = \gamma_1^2 \gamma_0 [ (1 + \beta_1 \beta_0) - \beta_1 (1 + \beta_0) \cos \theta_\alpha], \]  
\[ u_{\text{fl}} = \gamma_1^2 \gamma_0 [ - \beta_1 (1 + \beta_1 \beta_0) + (\beta_1 + \beta_0) \cos \theta_\alpha], \]  
\[ u_{\text{xs}} = \pm \gamma_1 \gamma_0 (\beta_1 + \beta_0) \sin \theta_\alpha, \]  
\[ x = x_0 - c \beta_1 t + \gamma_0 (\beta_1 + \beta_0) c \frac{\omega}{\omega_{cs}} \sin \theta_\alpha, \]  
\[ y = y_0 \pm \gamma_1 \gamma_0 (\beta_1 + \beta_0) c \left( 1 - \cos \theta_\alpha \right), \]  
\[ \theta_\alpha = \frac{\omega_{cs} [ t + \beta_1 (x_\text{w} - x_0)/c]}{\gamma_0 (1 + \beta_1 \beta_0) c}. \]

By numerically solving Equations (A15) and (A17), we can determine \( x_\alpha \) and \( \theta_\alpha \) and finally obtain the theoretical solutions of the pickup process.

**Appendix B**

**Test-particle Simulation**

To confirm that the transverse filamentary structures trigger the pickup process, we perform test-particle simulations. The particle pusher proposed by Vay (2008) is applied to this test-particle code. We consider the ambient magnetic field \( B_1 \), the motional electric field \( E_x = -\beta_1 B_1 \), the wakefield, and the filaments. The wakefield and filaments are modeled as

\[ E_x = \begin{cases} -E_{\text{wake}} \sin \frac{2\pi}{\lambda_{\text{wake}}} (x - ct), & (x < ct), \\ 0, & (x \geq ct) \end{cases} \]  

\[ B_z = B_1 + B_1 \sin \frac{2\pi}{\lambda_y}, \]  

\[ E_y = -\beta_1 B_1, \]

where \( E_{\text{wake}} \) is the wakefield amplitude, \( \lambda_{\text{wake}} \) is the wakefield wavelength, \( B_1 \) is the filament amplitude, and \( \lambda_y \) is the filament wavelength. Particles are injected at \( x = 0 \) toward the \(-x\) direction with the bulk Lorentz factor \( \gamma_1 = 40 \). The thermal velocity of the injected plasma flow in the plasma rest frame is \( \beta_0 = 0.1 \). Based on our PIC simulation results, we determined \( E_{\text{wake}} / B_1 = 0.2 \), \( \lambda_{\text{wake}} (c / \omega_{pe}) = 500 \), \( B_1 / B_1 = 1 \), and \( \lambda_y / (c / \omega_{pe}) = 15 \). The other parameters are identical to our PIC simulations.

Figure 10 shows the trajectories of the energetic electrons in \( x-y \) space (top) and \( x-\gamma \) space (bottom). The color maps represent \( B_1 \) in the case of \( B_1 = 1 \) (left) and \( B_1 = 0 \) (right). The incoming electrons are picked up at \( x / (c / \omega_{pe}) \approx -220 \) for \( B_1 = 1 \), whereas they are merely oscillating inside the wakefield for \( B_1 = 0 \). The trajectories give clear proof that the filaments are essential for entering the pickup process.

Figure 11 displays the enlarged view of the top left panel of Figure 10. The incoming electrons enter the unmagnetized region arising from the filaments, and then they are picked up by the bulk flow. The filaments obviously trigger the pickup process, as in the case of our PIC simulations.

Figure 12 shows the time evolution of a typical energetic electron in the same format as Figure 6. The time evolution in our test-particle simulation exhibits qualitatively the same behavior as that in our PIC simulation. The incoming electron is first decelerated by the wakefield and begins to gyrate. The motional electric field as well as the wakefield then decelerate it for \( 210 \leq \omega_{pe} t \leq 244 \). The increase in \( \Delta \gamma_\alpha \) and the positive velocity \( \beta_1 \) \( \sim 1 \) for \( 244 \leq \omega_{pe} t \leq 256 \) show that the electrons are accelerated by the wakefield via the Landau resonance. The electrostatic force dominates the Lorentz force for \( 244 \leq \omega_{pe} t \leq 256 \), and thus the Landau resonance can work. After \( \omega_{pe} t \approx 256 \), \( \Delta \gamma_\alpha \) becomes dominant, and the electron enters the pickup process. The Lorentz force indeed exceeds the electrostatic force after \( \omega_{pe} t \approx 256 \). The electron is released from the wakefield and then picked up by the bulk flow.

The injection into the pickup process is well described by this toy model. We thus think that the filaments trigger the pickup process.

Figure 13 shows the time evolution of the ion Lorentz factor. The maximum Lorentz factor is \( \gamma_{\text{max}, e} \sim \alpha_1^{-1} \) and the acceleration timescale is \( \omega_{pe} \tau_{\text{acc}, e} \sim 2\pi \alpha_1^{-3} \). This result is consistent with the theoretical estimate discussed in the main text, indicating that the pickup particles can be accelerated up to the theoretical maximum Lorentz factor in the shock system if the size of the precursor wave region is sufficiently large.

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**References**

Aab, A., Abreu, P., Aglietta, M., et al. 2018, ApJL, 853, L29  
Aartsen, M. G., Ackermann, M., Adams, J., et al. 2018, Sci, 361, 147  
Amato, E., & Arons, J. 2006, ApJ, 653, 325  
Arons, J. 2003, ApJ, 589, 871  
Chen, P., Tajima, T., & Takahashi, Y. 2002, PRvL, 89, 161101  
Drake, J. F., Kaw, P. K., Lee, Y. C., et al. 1974, PhFl, 14, 778  
Ebisuzaki, T., & Tajima, T. 2021, APh, 128, 102567  
Fried, B. D. 1959, PhFl, 2, 337  
Gallant, Y. A., Hoshino, M., Langdon, A. B., Arons, J., & Max, C. E. 1992, ApJ, 391, 73  
Hillas, A. M. 1984, ARA&A, 22, 425  
Hoshino, M. 2008, ApJL, 672, 940  
Hoshino, M., & Arons, J. 1991, PhFlB, 3, 871  
Iwamoto, M., Amano, T., Hoshino, M., et al. 2019, ApJL, 883, L35  
Kruer, W. L. 1988, in The Physics of Laser Plasma Interactions, ed. D. Pines (Boston, MA: Addison-Wesley)  
Kuramitsu, Y., Sakawa, Y., Hoshino, M., Chen, S. H., & Takabe, H. 2012, HEDP, 8, 266
Kuramitsu, Y., Sakawa, Y., Kato, T., Takabe, H., & Hoshino, M. 2008, ApJ, 682, 113
Kuramitsu, Y., Nakanii, N., Kondo, K., et al. 2011a, PhPl, 18, 010701
Kuramitsu, Y., Nakanii, N., Kondo, K., et al. 2011b, PhRvE, 83, 026401
Langdon, A. B., Arons, J., & Max, C. E. 1988, PhRvL, 61, 779
Ligorini, A., Niemiec, J., Kobzar, O., et al. 2021a, MNRAS, 501, 4837
Ligorini, A., Niemiec, J., Kobzar, O., et al. 2021b, MNRAS, 502, 5065
Liu, Y. L., Isayama, S., Chen, S. H., & Kuramitsu, Y. 2019, HEDP, 31, 64
Liu, Y. L., Kuramitsu, Y., Isayama, S., & Chen, S. H. 2018, PhPl, 25, 013110
Liu, Y. L., Kuramitsu, Y., Moritaka, T., & Chen, S. H. 2017, HEDP, 22, 46
Lyubarsky, Y. 2006, ApJ, 652, 1297
Lyubarsky, Y. 2008, ApJ, 682, 1443
Lyubarsky, Y. 2014, MNRAS, 442, L9
Lyubarsky, Y. 2018, MNRAS, 474, 1135
Lyubarsky, Y. 2019, MNRAS, 490, 1474
Margalit, B., Beniamini, P., Sridhar, N., & Metzger, B. D. 2020, ApJL, 899, L27
Matsumoto, Y., Amano, T., & Hoshino, M. 2013, PhRvL, 111, 215003
Matsumoto, Y., Amano, T., Kato, T. N., & Hoshino, M. 2015, Sci, 347, 974
Max, C. E., Arons, J., & Langdon, A. B. 1974, PhRvL, 33, 209
Melrose, D. B. 2017, RvMPP, 15, 1
Lyubarsky, Y. 2006, ApJ, 652, 1297
Lyubarsky, Y. 2008, ApJ, 682, 1443
Lyubarsky, Y. 2014, MNRAS, 442, L9
Lyubarsky, Y. 2018, MNRAS, 474, 1135
Lyubarsky, Y. 2019, MNRAS, 490, 1474
Margalit, B., Beniamini, P., Sridhar, N., & Metzger, B. D. 2020, ApJL, 899, L27
Matsumoto, Y., Amano, T., & Hoshino, M. 2013, PhRvL, 111, 215003
Matsumoto, Y., Amano, T., Kato, T. N., & Hoshino, M. 2015, Sci, 347, 974
Max, C. E., Arons, J., & Langdon, A. B. 1974, PhRvL, 33, 209
Melrose, D. B. 2017, RvMPP, 1, 5
Metzger, B. D., Margalit, B., & Sironi, L. 2019, MNRAS, 485, 4091
Möbius, E., Hovestadt, D., Klecker, B., et al. 1985, Natur, 318, 426
Murase, K., Mészáros, P., & Zhang, B. 2009, PhRvD, 79, 103001
Oka, M., Terasawa, T., Noda, H., Saito, Y., & Makai, T. 2002, GeoRL, 29, 1612
Piran, T. 2005, RvMP, 76, 1143
Plotnikov, I., Grassi, A., & Grech, M. 2018, MNRAS, 477, 5238
Plotnikov, I., & Sironi, L. 2019, MNRAS, 485, 3816
Shimada, N., & Hoshino, M. 2000, ApJL, 543, L67
Sironi, L., Plotnikov, I., Nättilä, J., & Beloborodov, A. M. 2021, PhRvL, 127, 035101
Sironi, L., & Spitkovsky, A. 2011, ApJ, 726, 75
Sironi, L., Spitkovsky, A., & Arons, J. 2013, ApJ, 771, 54
Sobacchi, E., Lyubarsky, Y., Beloborodov, A. M., & Sironi, L. 2020, MNRAS, 500, 272
Spitkovsky, A. 2005, in AIP Conf. Proc. 801, Astrophysical Source of High Energy Particles and Radiation, ed. T. Bulik, B. Rudak, & G. Madejski (Melville, NY: AIP), 345
Tajima, T., & Dawson, J. M. 1979, PhRvL, 43, 267
Vay, J.-L. 2008, PhPl, 15, 056701
Weibel, E. S. 1959, PhRvL, 2, 83