INJECTION TO RAPID DIFFUSIVE SHOCK ACCELERATION AT PERPENDICULAR SHOCKS IN PARTIALLY IONIZED PLASMAS

YUTAKA OHIRA
Department of Physics and Mathematics, Aoyama Gakuin University, 5-10-1 Fuchinobe, Sagamihara 252-5258, Japan; ohira@phys.aoyama.ac.jp

Received 2016 March 28; revised 2016 June 9; accepted 2016 June 9; published 2016 August 5

ABSTRACT

We present a three-dimensional hybrid simulation of a collisionless perpendicular shock in a partially ionized plasma for the first time. In this simulation, the shock velocity and upstream ionization fraction are \( v_{sh} \approx 1333 \, \text{km s}^{-1} \) and \( f_i \approx 0.5 \), which are typical values for isolated young supernova remnants (SNRs) in the interstellar medium. We confirm previous two-dimensional simulation results showing that downstream hydrogen atoms leak into the upstream region and are accelerated by the pickup process in the upstream region, and large magnetic field fluctuations are generated both in the upstream and downstream regions. In addition, we find that the magnetic field fluctuations have three-dimensional structures and the leaking hydrogen atoms are injected into the diffusive shock acceleration (DSA) at the perpendicular shock after the pickup process. The observed DSA can be interpreted as shock drift acceleration with scattering. In this simulation, particles are accelerated to \( v \sim 100 \, v_{sh} \sim 0.3 \, c \) within \( \sim 100 \) gyroperiods. The acceleration timescale is faster than that of DSA in parallel shocks. Our simulation results suggest that SNRs can accelerate cosmic rays to \( 10^{15.5} \, \text{eV} \) (the knee) during the Sedov phase.

Key words: acceleration of particles – ISM: supernova remnants – shock waves

1. INTRODUCTION

The origin and acceleration mechanism of Galactic cosmic rays (GCRs) are long standing problems in astrophysics. Supernova remnants (SNRs) are believed to be the source of GCRs, which is supported by recent gamma-ray observations (Ohira et al. 2011; Ackermann et al. 2013). One famous acceleration mechanism is diffusive shock acceleration (DSA; Axford et al. 1977; Krymsky 1977; Bell 1978; Blandford & Ostriker 1978). The DSA theory assumes that particles move around a shock diffusively. Then, some of the diffusive particles move back and forth between the upstream and downstream regions many times. As a result, the diffusive particles are accelerated by the shock. The DSA theory predicts that the momentum spectrum of the accelerated particles has a power-law form that is almost consistent with the GCR spectrum once the escape from the sources and the propagation effects are accounted for (Ohira et al. 2010). Therefore, DSA appears to be the most plausible acceleration mechanism of GCRs.

However, there are several problems for DSA. One is the injection problem. The DSA theory does not tell us how many particles are injected into DSA from thermal particles. The angle between the magnetic field and the shock normal, \( \theta_{Bnr} \), is thought to be the most important parameter for injection into DSA. The injection into DSA is expected to be more efficient for parallel shocks (\( \theta_{Bnr} = 0^{\circ} \)) than for perpendicular shocks (\( \theta_{Bnr} = 90^{\circ} \)) because the magnetic field perpendicular to the shock normal prevents charged particles from returning back to the upstream region (Ellison et al. 1995). Using hybrid simulations, Caprioli & Spitkovsky (2014) showed that the injection efficiency markedly decreases with \( \theta_{Bnr} \) for \( \theta_{Bnr} \gtrsim 50^{\circ} \) (see Figure 3 of Caprioli & Spitkovsky 2014).

Another problem is the knee problem. Since the CR spectrum observed at Earth has a spectral break at an energy of \( 10^{15.5} \, \text{eV} \) (this is the so-called knee), the energy of the knee is thought to be the maximum energy of CR protons accelerated by DSA in SNR shocks. However, if we assume a parallel shock with a typical magnetic field strength (a few \( \mu \text{G} \)) for the interstellar medium (ISM), then the maximum energy of DSA in SNRs cannot reach the knee energy because the acceleration timescale of DSA is too large at the parallel shocks (Lagage & Cesarsky 1983). If magnetic fields are amplified to a few hundred \( \mu \text{G} \), then the acceleration timescale becomes small, so that SNRs can accelerate CRs to the knee energy during the free expansion phase of SNRs. X-ray observations of young SNRs suggest magnetic field amplification in downstream regions (Berezhko et al. 2003; Vink & Laming 2003; Bamba et al. 2005; Uchiyama et al. 2007). So far, some magnetic field amplification mechanisms have been proposed and investigated using simulations. However, the dominant mechanism behind magnetic field amplification remains undetermined.

Another proposed solution for the knee problem is DSA at perpendicular shocks (Jokipii 1987). Accelerated particles propagate a short distance from the shock surface for the perpendicular shock compared with the parallel shock because particles easily propagate to the shock normal direction for the parallel shock. As a result, the timescale of the back and forth motion around the shock front, that is, the acceleration timescale, becomes small for the perpendicular shock. Fast acceleration at the perpendicular shock was confirmed by test-particle simulations (Decker & Vlahos 1986; Takahara & Terasawa 1991; Ellison et al. 1995; Naito & Takahara 1995; Giacalone 2005b; Takamoto & Kirk 2015). However, as mentioned above, injection into DSA is thought to be difficult for the perpendicular shock. In fact, using two-dimensional (2D) and three-dimensional (3D) hybrid simulations, Caprioli & Spitkovsky (2014) demonstrated that only acceleration via shock drift acceleration (SDA) lasts for a very short time in a pure perpendicular shock because particles are advected downstream, while some particles are injected into DSA and accelerated for a long time in a parallel shock. Using 2D hybrid simulations, Giacalone (2005a) showed that particles are...
injected into DSA (or multiple SDA) if there is magnetic turbulence in the upstream region of a globally perpendicular shock, but the upstream magnetic turbulence is not generated in a self-consistent manner and the acceleration timescale is not so small. Therefore, the upstream turbulence, the injection into DSA, and the fast acceleration of particles have never been simultaneously demonstrated in a self-consistent simulation.

Early studies of DSA at the perpendicular shock implicitly assumed that upstream plasmas are fully ionized. However, it is well known that the ISM and ejecta of SNRs are not always fully ionized. In fact, the observed line profile of Hα emission from SNRs supports some SNR shocks propagating to the partially ionized ISM (Chevalier & Raymond 1978; Katsuda et al. 2016). It was proposed that many plasma instabilities are excited by ionization around a collisionless shock in partially ionized plasma (Raymond et al. 2008; Ohira et al. 2009; Ohira & Takahara 2010; Ohira 2014). Ohira (2013) performed the first 2D hybrid simulation of a collisionless shock wave propagating into a partially ionized plasma and showed that the ionization of hydrogen atoms excites plasma instabilities both in the upstream and downstream regions. In addition, the simulation showed that downstream hydrogen atoms leak into the upstream region. In the upstream region, density fluctuations are excited by the ionization of leaking neutral particles from the downstream region (Ohira 2013, 2014). Moreover, Ohira (2016) performed a simulation similar to that of Ohira (2013), but the Alfvén Mach number and the shock velocity are higher and smaller, respectively. Then, very large fluctuations of magnetic fields and density are generated in the upstream and downstream regions. Although the leaking neutral particles are accelerated by the pickup process in the upstream region, no particles are injected into DSA in the 2D hybrid simulations of Ohira (2013, 2016). Since particle diffusion perpendicular to the magnetic field line is suppressed in the 2D system (Jokipii et al. 1993; Giacalone & Jokipii 1994), a 3D simulation is needed to understand whether or not injection into DSA is possible in perpendicular shocks.

In this paper, we perform the first 3D hybrid simulation of collisionless shocks propagating into the partially ionized ISM. We then show that the upstream turbulence is self-consistently generated, leaking hydrogen atoms are injected into the DSA, and the particles are rapidly accelerated by DSA at the perpendicular shock. In Section 2, we first briefly summarize our hybrid simulation. Then, we show the results of the 3D hybrid simulation in Section 3. Finally, we discuss the acceleration timescale and spectral index of DSA at the perpendicular shock in Section 4 and summarize in Section 5.

2. HYBRID SIMULATIONS

Hybrid simulations compute the induction equation for the magnetic field with a generalized Ohm’s law and a non-relativistic equation of motion for many protons, but electrons are treated as a massless fluid. In addition, our hybrid simulation solves the ionization and free-streaming motion of hydrogen atoms (Ohira 2013, 2016). We consider charge exchange and collisional ionization with protons, electrons, and hydrogen atoms as the ionization of hydrogen atoms. To compute collisional ionization with electrons, we need to assume the velocity distribution of electrons because the hybrid simulation does not solve for the dynamics of electrons. In this simulation, we assume a Maxwell distribution with a temperature $T_e$ (see Ohira 2016, for details). The ionization timescale for hydrogen atoms in typical isolated young SNRs is about $10^5$ times longer than the gyro period of protons. Since the huge gap in the timescale is still challenging for current supercomputers, all cross-sections are boosted by a factor of $2 \times 10^3$ in this paper. Then, the ionization timescale in this simulation is about $50 \Omega_{cp}^{-1}$, where $\Omega_{cp}$ is the proton cyclotron frequency. Furthermore, in order to reduce the simulation box, we set an upstream boundary in a region sufficiently far upstream where leaking neutral hydrogen atoms are artificially ionized.

Simulation particles are injected with a drift velocity parallel to the $x$ direction at the left boundary, $x = 0$, and are then specularly reflected at the right boundary $x = L_x$. As time goes on, a perpendicular shock propagates to the left boundary and the mean velocity becomes zero in the downstream region. Accordingly, the simulation frame is the downstream rest frame. We impose the periodic boundary condition in both the $y$ and $z$ directions. The simulation box size is $L_x \times L_y \times L_z = 8400 \ c/\omega_{pp} \times 200 \ c/\omega_{pp} \times 200 \ c/\omega_{pp}$, where $c$ and $\omega_{pp}$ are the speed of light and the upstream plasma frequency of the protons, respectively. The cell size and time step are set to $\Delta x = \Delta y = \Delta z = 1 \ c/\omega_{pp}$ and $\Delta t = 10^{-3} \Omega_{cp}^{-1}$.

To describe the upstream plasma, eight protons and eight hydrogen atoms per cell are used and the uniform magnetic field parallel to the $y$ direction is imposed, $B_0 = B_0 e_y$. The parameters of the upstream plasma are as follows. The ionization degree is 0.5. The ratio of the particle pressure to the magnetic pressure is $\beta_p = \beta_H = 0.5$ for protons and hydrogen atoms. The drift velocity parallel to the $x$ direction is $v_d = 20 \ v_A = 1000 \ km \ s^{-1}$, where $v_A = B_0/\sqrt{4\pi \rho_p0}$ is the Alfvén velocity and $\rho_p0$ is the proton mass density in the far upstream region. For the electron temperature, we assume $T_e = 0$ and $0.01 m_p v_d^2/3 = (0.01 \times (3/16) m_p v_A^2)$ in the upstream and downstream regions, respectively, where $m_p$ is the proton mass. The factor of 0.01 is based on several studies of electron heating in collisionless shocks (Ghavamian et al. 2007; Ohira & Takahara 2007, 2008; Rakowski et al. 2008; van Adelsberg et al. 2008; Laming et al. 2014; Vink et al. 2015). We have not investigated the electron temperature dependence, but rather leave that for future work. The number of leaking hydrogen atoms depends on the electron temperature (e.g., Morlino et al. 2012).

3. SIMULATION RESULTS

3.1. Phase Space of Protons

We show the phase space of protons at $t = 400 \Omega_{cp}^{-1}$ in Figure 1, where the vertical and horizontal axes show the kinetic energy, $E$, normalized by the upstream kinetic energy, $E_0 = m_p v_d^2/2$, and the $x$ coordinate, respectively. The shock front is propagating to the left and is located at $x \approx 5700 \ c/\omega_{pp}$. The shock velocity is $v_A \approx 26.7 \ v_A = 37.7 \ v_{A,\text{tot}} = 1333 \ km \ s^{-1}$ in the upstream rest frame, where $v_{A,\text{tot}} = B_0/\sqrt{4\pi \rho_p0 + \rho_H0}$ is the Alfvén velocity defined by the total mass density. The upstream cold component with $E \sim E_0$ and the upstream hot component with $E \sim 10E_0$ are the upstream protons and pickup ions generated by the ionization of hydrogen atoms leaking from the downstream region. For shocks propagating to partially ionized plasmas, the hot hydrogen atoms that can leak into the upstream region are generated by the charge exchange process in the downstream region (Lim & Raga 1996; Blasi et al. 2012; Ohira 2012, 2013, 2016). In this simulation, the flux of leaking
hydrogen atoms is a few percent of the upstream hydrogen flux in the shock rest frame. In addition, there are protons accelerated to $E \sim 10^4 E_0$ around the shock front that were not observed in previous 2D simulations (Ohira 2013, 2016). Since 2D simulations cannot accurately follow diffusion perpendicular to the magnetic field line but 3D simulations can (Jokipii et al. 1993; Giacalone & Jokipii 1994), we can observe particle acceleration around the shock front for the first time in this simulation. We discuss the acceleration mechanism in Section 3.3.

3.2. Three-dimensional Structures

Figures 2 and 3 show 3D structures at $t = 400 \Omega_{sp}^{-1}$ for the proton density, $\rho_p/\rho_{p,0}$, and the magnetic field strength, $|\mathbf{B}|/B_0$, respectively. In Figures 2 and 3, panels (a) and (b) show 2D slices of the 3D simulation box at $z = 100 c/\omega_{pp}$ and $y = 100 c/\omega_{pp}$, respectively. Panels (c)–(g) show slices through the $y$-z plane at $x = 5300$, 5600, 5700, 5800, and 6100 $c/\omega_{pp}$, respectively. The shock structures averaged over the $y$ and $z$ directions are very similar to Figure 2 of Ohira (2016).

In the upstream region ($x = 5300$ $c/\omega_{pp}$), the variations are almost isotropic and the density is well correlated with the magnetic field strength. Around the shock front ($x = 5600$, 5700, and 5800 $c/\omega_{pp}$), the proton density and magnetic field strength are very anisotropic and the density is anticorrelated with the magnetic field strength. In the far upstream region, the fast magnetosonic mode is excited by leaking neutral particles (Ohira 2014), so that there is a correlation between the density and magnetic field strength. However, the Alfvén ion cyclotron instability around the shock front becomes the fastest-growing instability because there are many pickup ions around the shock front compared with the far upstream region (Ohira 2016). The Alfvén ion cyclotron instability is driven by the temperature anisotropy of the pickup ions and excites Alfvén waves propagating along the magnetic field line. Alfvén waves with large amplitude push plasma along the magnetic field line, so that the high-density sheets are generated at regions where the amplitude of the Alfvén waves is small. This is why the anticorrelation appears. The strong turbulence is generated by interactions between the high-density sheets and the shock front. As a result, magnetic fields are strongly amplified to the equipartition level by the turbulence in the downstream region. In the far downstream region ($x = 6100 c/\omega_{pp}$), the high-density sheets are disrupted by the downstream turbulence, and so variations become almost isotropic.

These behaviors were observed in previous 2D simulations for the high Alfvén Mach number shock (Ohira 2016). Therefore, we confirm that 2D simulations can approximately follow the actual behavior of the density and magnetic fields on the $xy$ plane. However, Figures 2 and 3 show that there are 3D structures. This is crucial for the injection into DSA. In fact, we observe particle acceleration up to $10^4 E_0$ in the 3D simulation (see Figures 1 and 4) but not in the 2D simulation. In the next subsection, we show the energy spectrum of the protons and the trajectories of the accelerated particles for the 3D simulation.

3.3. Energy Spectrum and Particle Trajectory

Figure 4 shows the energy spectra of protons in the region of $2000 c/\omega_{pp} \leq x \leq 7000 c/\omega_{pp}$ at $t = 400 \Omega_{sp}^{-1}$. For the 2D simulation (the black histogram), some protons are typically accelerated to about $10 E_0$, where $E_0 = m_p v_0^2 / 2$ is the initial kinetic energy. When leaking neutral particles are ionized in the upstream region, they are picked up by the upstream flow and accelerated to a few times $E_0$. Then, they are advected to the shock and accelerated by the second shock heating (Ohira 2013). Even though there is strong magnetic turbulence, the diffusion perpendicular to the magnetic field line is strongly suppressed in the 2D system. Therefore, the pickup ions are not injected into DSA in the 2D simulation. The second shock heating can be interpreted as SDA in the scatter-free system and accelerates particles only to about four times its energy. Hence, the typical energy of the accelerated particles becomes about $10 E_0$ for the 2D simulation.

On the other hand, for the 3D simulation (the red histograms), particles are very efficiently accelerated to $E \sim 10^4 E_0$ ($v \sim 10^2 v_g \sim 0.3 c$) at $t = 400 \Omega_{sp}^{-1}$ and the maximum energy is still increasing, where $v$ is the velocity. In the downstream region, the energy fractions of the nonthermal particles with $E > 4 E_0$ and $E > 10^4 E_0$ are about 30% and 1% of the total kinetic energy, respectively.

In order to understand the acceleration mechanism in the 3D simulation, we show the trajectories of an accelerated particle in Figure 5 where panel (a) shows the time evolution of the kinetic energy and panels (b)–(d) show the particle trajectories by assuming no periodic boundaries. The black curve in panel (b) represents the mean position of the shock front. As one can see in panel (b) of Figure 5, the particle interacts with the shock at $t \sim 10 \Omega_{sp}^{-1}$, becomes a neutral hydrogen atom via the charge exchange process, and leaks into the upstream region. At $t \sim 60 \Omega_{sp}^{-1}$, the leaking hydrogen atom is ionized and picked up by the upstream flow. The pickup ion interacts with the shock front again at $t \sim 100 \Omega_{sp}^{-1}$. For the 2D simulation, the particle is not accelerated further and is simply advected to the downstream region, while the particle is trapped around the shock front and accelerated for the 3D simulation. After $t \sim 100 \Omega_{sp}^{-1}$, the orbit in the upstream region is almost a simple gyro motion in the uniform magnetic field and the particle crosses the shock front at every gyro motion. The energy gain is correlated with the motion in the $-z$ direction which is the same as the direction of the motional electric field, $E_m = -(q_0/c) \times B_0$. Therefore, the acceleration mechanism
is SDA. Contrary to the 2D system, particles can diffuse perpendicular to the magnetic field line in the 3D system. As a result, particles can go back to the upstream region and SDA with scattering can accelerate particles for a long time, which can be interpreted as DSA at the perpendicular shock. We track the trajectories of the 500 most energetic particles and confirm that all of the energetic particles are accelerated by DSA at the perpendicular shock after they leak into the upstream region through the charge exchange process. Hence, our simulation shows that the mechanism of injection into DSA at the perpendicular shock in partially ionized plasma is the leakage of hydrogen atoms through the charge exchange process.

For SDA, during one gyro motion, $\Delta t \sim \Omega_{cp}^{-1}$, the accelerated particles drift a distance of about $r_g = v/\Omega_{cp}$ in the direction of the motional electric field, so that the energy gain per gyro period is $\Delta E \sim |eE_m|r_g$. Then, the acceleration rate is given by

$$\frac{dE}{dt} \approx \frac{\Delta E}{\Delta t} = \alpha p_0 v \Omega_{cp},$$

(1)

where $p_0 = m_e v_0$ and $\alpha$ are the initial upstream momentum and a numerical factor of the order of unity, respectively. From the above equation, the momentum increasing rate is represented by

$$\frac{dp}{dt} = \alpha p_0 \Omega_{cp},$$

(2)

and the solution is given by

$$p = \alpha p_0 \Omega_{cp} (t - t_s) + p_s,$$

(3)

where $p_s$ is the momentum at the start time of SDA, $t_s$. Figure 6 shows the time evolution of the momentum for the 10 most energetic particles (gray curves) and their average (red curve). After the initial pickup process ($t > 100 \Omega_{cp}^{-1}$), the particles are accelerated and their momentum increases with time linearly as predicted by Equation (3). Hence, the momentum evolution is consistent with SDA. $\alpha = 0.39$ is derived from the fit of Equation (3) to the average curve, and the fit line is shown by the blue dashed curve in Figure 6.

It should be noted that the ionization timescale is artificially lowered in this simulation. For actual young SNRs, the timescale of the initial pickup process typically becomes $10^5 \Omega_{cp}^{-1} \approx 10^7$ s.
Figure 3. Same as Figure 2, but for the magnetic field strength normalized by the upstream value, $|B|/B_0$.

Figure 4. Energy spectrum of protons in the region $2000 \, c/\omega_{pp} < x < 7000 \, c/\omega_{pp}$ at $t = 400 \, \Omega_{cp}^{-1}$. The black and red histograms show the energy spectra for 2D and 3D simulations, respectively.

Figure 5. Trajectory of an accelerated particle. Panel (a) shows the time evolution of the energy, and panels (b)–(d) show the particle trajectories in the $x$, $y$, and $z$ coordinates, respectively. The back curve in panel (b) shows the mean position of the shock front.
Since the 10 most energetic particles in this simulation are still
the knee
shock
minimum acceleration timescale of DSA at the perpendicular
shock. As already mentioned, SDA with scattering can be
interpreted as DSA at the perpendicular shock. In fact, the
shock front is strongly modified by the leaking neutral
particles (Blasi et al. 2012; Ohira 2012, 2013, 2016)
and the spectral index around $E \approx 10^4 E_0$ due to a finite time. In order to
determine the spectral index more precisely, we require a
longer simulation with a more realistic ionization timescale.

Finally, we discuss some effects caused by the limitations of
this simulation. In this simulation, we use artificially large
cross-sections of several ionization processes in order to reduce
the ionization length scale and timescale. For actual SNR
shocks, the length scale of the neutral precursor region is about
$10^3 (v_{th}/\gamma_0c/\omega_{pp})$ in the typical ISM, which is much larger than
that of this simulation, $\sim 10^4 c/\omega_{pp}$. Since the gyroradius of
protons with $E = 10^4 E_0$ is also $r_g \sim 10^3 c/\omega_{pp}$ in the upstream
region of this simulation, the spectrum of accelerated particles
is affected by the artificially small precursor. Furthermore, the
smaller precursor in this simulation might result in an
artificially stronger turbulence. In addition, the size of the
simulation box transverse to the shock normal, $L_y = L_z = 200 c/\omega_{pp}$, is smaller than the gyroradius of the
accelerated protons in the upstream region. If we use a larger
simulation box, then the turbulence with a larger length scale
would be excited and the accelerated particles would be
efficiently scattered in the upstream region. If so, then the
accelerated particles diffuse to the upstream region and the
acceleration timescale becomes longer than that observed in
this simulation. Furthermore, the injection to DSA and the
acceleration timescale should depend on many parameters
/ionization fraction, electron temperature, and magnetic field
orientation). Since perfect perpendicular shocks are never
realized in astrophysical environments, it would be interesting
to determine whether or not the results in this paper can be
applied to oblique shocks. In order to study the above
problems, we need to perform a more realistic large simulation
and systematic studies. These will be addressed in future work.

4. DISCUSSION

In this section, we first discuss the acceleration timescale.
Since the 10 most energetic particles in this simulation are still
accelerating at $t = 400 \Theta_{cp}^{-1}$, we can expect further acceleration.
From Equation (2), the acceleration timescale of SDA observed
in this simulation is given by

$$t_{acc} = \frac{p}{\alpha \Omega_{cp}} = \frac{1}{\alpha \Omega_{cp}} \left( \frac{p}{p_0} \right),$$

where $\frac{p}{p_0} \approx \frac{v_{th}}{v_0}$ and $p \approx E/c$ are used in the last equation. This
is about 100 times smaller than that of DSA at the parallel
shock. As already mentioned, SDA with scattering can be
interpreted as DSA at the perpendicular shock. In fact, the
minimum acceleration timescale of DSA at the perpendicular
shock (Jokipii 1987) is the same as Equation (4). Since the
deceleration timescale of SNRs from the free expansion
velocity ($\sim 10^4 \mathrm{km} \mathrm{s}^{-1}$) to $10^3 \mathrm{km} \mathrm{s}^{-1}$ is about 10 kyr and the
SNR size at that time ($\sim 10 \mathrm{pc}$) is larger than the gyroradius of
the knee ($r_g \approx 1 \mu \mathrm{G}$), CRs could be accelerated
to the knee by SDA during the Sedov phase of SNRs. It should
be noted that two conditions have to be fulfilled to extrapolate
Equation (4) to the knee. One condition is that the coherent
length scale of the ordered magnetic field in the ISM must be
larger than the gyro radius of the knee, $r_g$. The other
condition is that particles with $E = 10^{1+5.5} \mathrm{eV}$ have to be
efficiently scattered by magnetic turbulence in the downstream
region in order to return to the upstream region. Since our
simulation is too small, it has not yet shown that the latter
condition is fulfilled.

Next, we discuss the spectral index of the accelerated
particles. The spectral index of the accelerated particles, $s = -d \log (dN/dE)/d \log E$, in the nonrelativistic energy
range is $s = 1.5$ for the standard DSA at the parallel shock
(Bell 1978; Blandford & Ostriker 1978). Figure 7 shows the
spectral index that is obtained from Figure 4. As shown in

\begin{align*}
\frac{d}{dE} = \frac{d}{dE} (dN/dE) \frac{dE}{d \log E} = \frac{d}{d \log E} (dN/\log E) = \frac{d}{d \log E} (dN/dE) \frac{dE}{d \log E}.
\end{align*}
5. SUMMARY

In this paper, we have performed the first 3D hybrid simulation of collisionless shocks generated by SNRs in partially ionized plasma. The simulation results can be summarized as follows.

1. Similar to our previous simulation, some downstream hydrogen atoms leak into the upstream region and modify the shock structures.
2. Large magnetic field and density fluctuations are excited both in the upstream and downstream regions. They have 3D structures.
3. The leaking hydrogen atoms are accelerated by the pickup process in the upstream region. Then, they are injected into DSA (or SDA) at the perpendicular shock.
4. Particles are accelerated to $v \sim 100 v_{th} \sim 0.3 c$ at $t = 400 \Omega^{-1}$. The acceleration rate in this simulation is faster than that of the typical DSA up to the end of the simulation.

Isolated SNRs (e.g., SNRs of SNe Ia) are expected to be generated frequently in the partially ionized ISM. Therefore, our results show that such SNRs can produce GCRs and suggest that such SNRs can accelerate CRs to the knee energy if the shock is a perpendicular shock.

Numerical computations were carried out on the XC30 system at the Center for Computational Astrophysics (CfCA) of the National Astronomical Observatory of Japan. We thank T. Terasawa, T. Inoue, and R. Yamazaki for useful comments. We are also grateful to the anonymous referee for useful suggestions and comments. This work was supported in part by Grants-in-Aid for Scientific Research of the Japanese Ministry of Education, Culture, Sports, Science and Technology No. 16K17702.

REFERENCES

Ackermann, M., Ajello, M., Allafort, A., et al. 2013, Sci, 339, 807
Axford, W. I., Leer, E., & Skadron, G. 1977, Proc. ICRC (Plovdiv), 11, 132
Bamba, A., Yamazaki, R., Yoshida, T., Terasawa, T., & Koyama, K. 2005, ApJ, 621, 793
Bell, A. R. 1978, MNRAS, 182, 147
Berezhko, E. G., Ksenofontov, L. T., & Völk, H. J. A&A, 412, L11
Blandford, R. D., & Ostriker, J. P. 1978, ApJL, 221, L29
Blasi, P., Morlino, G., Bandiera, R., Amato, E., & Caprioli, D. 2012, ApJ, 755, 121
Caprioli, D., & Spitkovsky, A. 2014, ApJ, 783, 91
Chevalier, R. A., & Raymond, J. C. 1978, ApJL, 225, L27
Decker, R. B., & Vlahos, L. 1986, ApJ, 306, 710
Ellison, D. C., Baring, M. G., & Jones, F. C. 1995, ApJ, 453, 873
Ghavamian, J., Laming, J. M., & Rakowski, C. E. 2007, ApJL, 654, L69
Giacalone, J. 2005a, ApJ, 624, 765
Giacalone, J. 2005b, ApJL, 628, L37
Giacalone, J., & Jokipii, J. R. 1994, ApJL, 430, L137
Jokipii, J. R. 1987, ApJ, 313, 842
Katsuda, S., Maeda, K., Ohira, Y., et al. 2016, ApJL, 819, L32
Krymsky, G. F. 1977, DoSSR, 234, 1306
Laming, J. M., Hwang, U., Ghavamian, P., & Rakowski, C. 2014, ApJ, 799, 11
Lim, A. J., & Raga, A. C. 1996, MNRAS, 280, 103
Morlino, G., Bandiera, R., Blasi, P., & Amato, E. 2012, ApJ, 760, 137
Naito, T., & Takahara, F. 1995, MNRAS, 275, 1077
Ohira, Y. 2012, ApJ, 758, 97
Ohira, Y. 2013, PhtoV, 111, 245002
Ohira, Y. 2014, MNRAS, 440, 414
Ohira, Y. 2016, ApJ, 817, 137
Ohira, Y., Murase, K., & Yamazaki, R. 2010, A&A, 513, A17
Ohira, Y., Murase, K., & Yamazaki, R. 2011, MNRAS, 410, 1577
Ohira, Y., & Takahara, F. 2007, ApJL, 661, L171
Ohira, Y., & Takahara, F. 2008, ApJ, 686, 320
Ohira, Y., & Takahara, F. 2010, ApJL, 721, L43
Ohira, Y., Terasawa, T., & Takahara, F. 2009, ApJL, 703, L59
Rakowski, C. E., Laming, J. M., & Ghavamian, P. 2008, ApJ, 664, 348
Raymond, J. C., Isenberg, P. A., & Laming, J. M. 2008, ApJ, 682, 408
Takahara, F., & Terasawa, T. 1991, in Proc. ICRR Int. Symp. 291, Astrophysical Aspects of the Most Energetic Cosmic Rays, ed. M. Nagano & F. Takahara (Singapore: World Scientific), 291
Takamoto, M., & Kirk, J. G. 2015, ApJ, 809, 29
Uchiyama, Y., Aharonian, F. A., Tanaka, T., Takahashi, T., & Maeda, Y. 2007, Natur, 449, 576
van Adelsberg, M., Heng, K., McCray, R., & Raymond, J. C. 2008, ApJ, 689, 1089
Vink, J., Broersen, S., Bykov, A., & Goubeli, S. 2015, A&A, 579, A13
Vink, J., & Laming, J. M. 2003, ApJ, 584, 758