MAGNETIC FIELD DECAY DUE TO THE WAVE–PARTICLE RESONANCES IN THE OUTER CRUST OF NEUTRON STARS

HIROYUKI R. TAKAHASHI1, KEI KOTAKE1,2, AND NOBUTOshi YASUKATE2

1 Center for Computational Astrophysics, National Astronomical Observatory of Japan, Mitaka, Tokyo 181-8588, Japan; takahashi@cfca.jp
2 Division of Theoretical Astronomy, National Astronomical Observatory of Japan, Mitaka, Tokyo 181-8588, Japan

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ABSTRACT

Bearing in mind the application to the outer crust of neutron stars (NSs), we investigate the magnetic field decay by means of the fully relativistic Particle-In-Cell simulations. Numerical computations are carried out in two dimensions, in which the initial magnetic fields are set to be composed both of the uniform magnetic fields that model the global fields penetrating the NS and of the turbulent magnetic fields that would originate from the Hall cascade of the large-scale turbulence. Our results show that the whistler cascade of the turbulence transports the magnetic energy preferentially in the direction perpendicular to the uniform magnetic fields. It is also found that the distribution function of electrons becomes anisotropic because electrons with lower energies are predominantly heated in the direction parallel to the uniform magnetic fields due to the Landau resonance, while electrons with higher energies are heated mainly by the cyclotron resonance that makes the distribution function isotropic for the high energy tails. Furthermore, we point out that the degree of anisotropy takes on the maximum value as a function of the initial turbulent magnetic energy. As an alternative to the conventional Ohmic dissipation, we propose that the magnetic fields in the outer crust of NSs, cascading down to the electron inertial scale via the whistler turbulence, would decay predominantly by the dissipation processes through the Landau damping and the cyclotron resonance.

Key words: plasmas – stars: neutron – turbulence

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1. INTRODUCTION

Pushed by the accumulating observations of radio pulsars, accreting neutron stars (NSs), and most recently magnetars, extensive studies have been carried out to understand the evolution of magnetic fields in NSs (e.g., Bhattacharya & Srinivasan 1995; Harding & Lai 2006; Reisenegger 2009). The known radio pulsars so far are generally categorized into two classes: young (≤107 yr) pulsars with spin periods \( P \approx 10^{-1–1} \) s and the magnetic field strength \( B \approx 10^{10–15} \) G, and the old millisecond radio pulsars, which have magnetic field strength as low as \( 10^{-8–9} \) G. While most radio pulsars are isolated objects, the millisecond pulsars are predominantly in binaries, suggesting that the magnetic fields decay with time, perhaps by an accretion of matter from the binary companion (e.g., Harding & Lai 2006 and references therein).

For the isolated radio pulsars, it is still an open question whether or not the NS magnetic fields decay with time. For example, Narayan & Ostriker (1990) argued that the field should decay exponentially on a few Myr timescale (see also Cordes & Chernoff 1998; Gontiher et al. 2004), while similar studies implied the decay time \( \gtrsim 100 \) Myr (e.g., Bhattacharya et al. 1992; Faucher-Giguère & Kaspi 2006). Such divergent results may come from the uncertainties inherent to those population synthesis studies, such as the selection effects, the luminosity evolution, and the dependence of beaming fraction on period.

On the other hand, the discovery of magnetars has provided several important evidences to favor the magnetic field decay on relatively short timescales (e.g., Arras et al. 2004; Harding & Lai 2006). In fact, the soft gamma-ray repeaters and anomalous X-ray pulsars, observed as young (≤107 yr) NSs with strong (\( 10^{14–10} \) G) magnetic fields, are considered to be powered by the decay of their magnetic fields (e.g., Woods & Thompson 2006). Giant flaring activities are proposed to be the rapid release of the magnetic stress building in the NS crust (Thompson & Duncan 1995, 1996; Lyutikov 2006). More recently, Pons et al. (2007) presented an evidence that the magnetic field decay of \( \sim 10 \) Myr can explain the thermal evolution from magnetars continuously to ordinary radio pulsars. At present, these ideas and new observations seem to favor the existence of the magnetic field decay in isolated NSs.

The pioneering study by Goldreich & Reisenegger (1992) identified the dissipation processes of the magnetic energy in the crust of the isolated NS during its evolution. They first showed that magnetized turbulence in the crust of NSs can be described by the electron magnetohydrodynamic (EMHD) equation, in which the time evolution of the magnetic field is governed by the advection of the field by the Hall drift, the Ohmic dissipation, and the ambipolar diffusion. On top of the Ohmic decay and the ambipolar diffusion, they first proposed that the Hall drift, though non-dissipative itself, could be an important ingredient for the field decay because it can lead to dissipation through the whistler cascade of the turbulence. As the eddy size of the turbulence becomes smaller, the magnetic energy was transported from large to small scales, leading to the dissipation of the magnetic energy finally via the Ohmic dissipation.

In order to confirm their prediction, numerical simulations in EMHD are required because such a cascading process is essentially a nonlinear process. In the two- and three-dimensional (2D and 3D) simulations of low \( \beta \) plasma, Biskamp et al. (1999) showed a clear cascade of the energy due to the whistler turbulence over more than an order of magnitude in the length scale. The 3D EMHD simulations by Cho & Lazarian (2004, 2009) confirmed the scale-dependent anisotropy in the EMHD turbulence. They also showed that the anisotropic cascading processes in the EMHD proceed via the propagation of the whistler waves along the global magnetic fields (see...
also Narita & Gary 2010). This is in contrast to conventional MHD turbulence, in which the Alfvén waves play a major role in transporting turbulent magnetic energy (e.g., Goldreich & Sridhar 1995; Biskamp & Welcher 1989; Cho et al. 2002). Reflecting the fact that the propagation of the whistler waves are more dispersive than the Alfvén waves, the power spectral density (PSD) in the EMHD turbulence was shown to become steeper compared to the MHD turbulence. Here, it should be noted that in the EMHD simulations one can precisely follow the cascading process; however, the dissipation process should be treated phenomenologically, i.e., via the resistivity or the so-called hyperdiffusivity (Cho & Lazarian 2004, 2009).

In this paper, we perform Particle-In-Cell (PIC) simulations, aiming to understand the decay process of the magnetic fields in NSs. As is well known, the PIC simulations, which have often been performed to study the turbulent cascade in solar wind (Saito et al. 2008; Gary et al. 2008, 2010), can precisely take into account the electromagnetic modes in plasma by solving the full Maxwell equations coupling with all the species of plasma particles. The ordinary PIC simulation can treat the collisionless plasma in which the collisional frequency is smaller than the typical frequencies of plasma particles, such as the plasma frequency and the gyrofrequency. Shtrum & Yakovlev (2006) have estimated the collisional frequency between the particles in NSs. From their results, the electron collisional frequency in the outer crust of the NS is of the order of $10^{-2} \omega_{pe}$ with $\omega_{pe}$ being the electron plasma frequency, which indicates that the electrons are marginally collisionless in the electron inertial scale. Since the gyrofrequency is of the order of the plasma frequency, we approximately treat the plasma as collisionless to study the decay process of the magnetic fields in the outer crust of the NS. As a complement to the foregoing EMHD simulations (e.g., Biskamp et al. 1999; Cho & Lazarian 2004, 2009) that focus on the whistler cascade of the turbulence, the PIC simulations that we will present in this paper are for understanding the mechanism of the magnetic field decay, cascading further down to the electron inertial scale.

This paper is organized as follows. In Section 2, we show the numerical setup of the PIC simulations. The numerical results of the turbulent cascade and the dissipation process are shown in Section 3. We summarize our results and discuss their implications in Section 4.

2. NUMERICAL METHODS AND INITIAL CONDITIONS

We carry out 2D fully relativistic PIC simulations (e.g., Birdsell & Langdon 1985). The basic equations are described as

$$\frac{dp_i}{dt} = q_i(E + \beta_i \times B), \quad \beta_i \equiv v_i/c$$

(1)

$$\nabla \cdot B = 0,$$

(2)

$$\nabla \cdot E = 4\pi \rho_e,$$

(3)

$$\frac{\partial B}{\partial t} + c \nabla \times E = 0,$$

(4)

$$\frac{\partial E}{\partial t} - c \nabla \times B = -4\pi j,$$

(5)

$$\rho_e = \sum_i q_i S(x - x_i),$$

(6)

$$j = \sum_i q_i v_i S(x - x_i).$$

(7)

where $p_i$, $q_i$, $E$, $B$, $\rho_e$, and $j$ are the momentum, the particle charge, the electric field, the magnetic field, the charge density, and the current density, respectively. $\beta_i \equiv v_i/c$ is the three velocity normalized by the speed of light. The subscript $i$ denotes the particle species. Here, we ignore the quantum effects such as the Landau level, for simplicity. The particle motion is determined by solving the special relativistic equation of motion.

Then, the electric charge and the current density are obtained from Equations (6) and (7), in which the shape factor $S$ can extrapolate the physical quantities on discretized grids at $x$ from particles. By using the electric charge and the current density, the electromagnetic fields are updated by solving Maxwell equations, so that the system can be solved self-consistently. We solve these equations in the rectangular coordinates by assuming $\partial / (\partial z) = 0$ (the so-called 2.5 dimension). The number of grid points is $(N_x, N_y) = (1024, 1024)$ and the corresponding system size is $L_x = L_y = 102.4c/\omega_{pe}$, where $\omega_{pe} \equiv \sqrt{4\pi ne^2/m_e}$ is the electron plasma frequency. Each cell contains 80 particles and the total number of particles in the simulation box is $8.4 \times 10^7$ for each species. The boundary conditions are periodic in both directions. We assume that ions are immovable, which is reasonable in the outer crust of the NS (e.g., Goldreich & Reisenegger 1992). We confirmed that the numerical results are qualitatively and quantitatively consistent with those when we take into account the ion motion.

In the initial state, the plasma is distributed uniformly in space and the electron distribution function obeys the relativistic Maxwellian with the temperature $T_e = 10^8$ K (e.g., Aguclera et al. 2008). We assume that there are two components of the magnetic fields in the NS. The first one is the uniform magnetic field that would globally penetrate the NS, which should be necessary to explain the spin-down of the NS. The global field is set to be uniform in space, $B = B_0 e_z$, where $e_z$ is the unit vector in the $z$-direction. The corresponding ratio of the gyrofrequency evaluated from the uniform magnetic fields $\omega_{ge} = eB_0/(m_e c)$ and the plasma frequency $\omega_{pe}$ is 0.5. The second one is the turbulent magnetic fields, which would be cascaded down from large to small scales due to the Hall turbulence. The initial turbulent magnetic fields are modeled to take the form of $B_i(k) \propto \exp(-k^2/k_0^2 + i \alpha_0)$, where $k$ is the wave number, $k_0 = 0.31 \omega_{pe}/c$ is the typical wave number of the initial fluctuations, and $\alpha_0$ is the random phase. The hat denotes the variable in the Fourier space. We change the amplitude of the turbulent magnetic fields $B_i \equiv \tilde{B}_i - B_0 e_z$, by introducing a parameter, $\epsilon \equiv \int dV B_i(t = 0)^2 / \int dV B_0^2 = 0.1, 0.4, 0.8$. We also carry out the simulation with $\epsilon = 0.0$ to assess the validity of our simulation codes.

3. RESULTS

3.1. Energy Exchange

Figure 1 shows the turbulent magnetic field energy density (color) and the magnetic field lines (white curves) at the initial state $t\omega_{pe} = 0$ (left) and at the end of the simulation $t\omega_{pe} = 1350$ (right) for $\epsilon = 0.1$. It is clearly shown that the turbulent magnetic field energy given at the initial state (bright spots in the left panel) decreases at the end of the simulations (right panel). The pattern size of the turbulent magnetic fields becomes relatively smaller with time. These results imply that the magnetic energy is converted to the plasma energy.

To see clearly how the energy conversion proceeds, we calculate deviations of each energy from the initial state, such as the electron kinetic energy $\delta E_{kin} \equiv \sum_i m_i (\gamma_i(t) - \gamma(t = 0))/E_{0i}$.
and the turbulent magnetic field energies $\delta E_{tB,xy} = \int dV \left[ (B_x^2 + B_y^2) - |B_z^2(\tau = 0) + B_z^2(\tau = 0)|/(8\pi E_0) \right]$ and $\delta E_{tB,y} = \int dV \left[ B_y^2(\tau = 0)/\left(8\pi E_0\right) - B_y^2(\tau = 0)/\left(8\pi E_0\right)\right]$, where $E_0 = \int dV B_z^2/(8\pi)$ is the magnetic energy of uniform fields. Figure 2 depicts the time evolution of the energy deviation of the electrons $\delta E_{\text{kin}}$ (solid curves) and of the negative of the turbulent magnetic energies, $-\delta E_{tB,xy}$ (dashed curves) and $-\delta E_{tB,y}$ (dot-dashed curves). It can be seen that the turbulent magnetic energy in the perpendicular direction ($\delta E_{tB,xy}$) decreases with time (note the minus sign), while the particle kinetic energy ($\delta E_{\text{kin}}$) increases with it. The energy deviation of the particles $\delta E_{\text{kin}}$ is almost comparable to $-\delta E_{tB,xy}$. This means that the perpendicular component of the magnetic energy is preferentially converted to the particle energy. Here it should be noted that the small difference between $\delta E_{\text{kin}}$ and $-\delta E_{tB,xy}$ comes from the electric field energy generated by turbulent motions. The particle energy distribution function obtained in this simulation can be well fitted by the Maxwellian distribution. Therefore, the plasma is not accelerated, but is heated up through the particle–wave interaction. The resulting heating is expected to become larger for models with larger initial turbulent magnetic energy (see models with different $\epsilon$ in Figure 2). The efficient energy conversion to the perpendicular component also suggests that the energy dissipation will proceed anisotropically, which acts to make the configuration of the magnetic-field lines approach the uniform one, i.e., the potential field.

3.2. Anisotropy

To see the anisotropy in more detail, we perform the Fourier analysis of the magnetic fields. The top panels of Figure 3 show the PSD of the turbulent magnetic field $|\mathbf{B}_t|/B_0$ at $t\omega_{pe} = 1350$. It can be seen that the anisotropic turbulence preferentially transports the magnetic energy perpendicular to the uniform magnetic fields. The PSD is larger for a larger $\epsilon$, suggesting that the larger amplitudes of the turbulence lead to an efficient electron heating. The lower panels of Figure 3 show one-dimensional cumulative power spectrum densities of the turbulent magnetic fields at $t\omega_{pe} = 1350$. The cumulative spectrum is defined as $|\mathbf{B}_t(k_x)|^2 = \sum_k |\mathbf{B}_t(k_x, k_y)|^2$ (thick solid curves) and $|\mathbf{B}_t(k_y)|^2 = \sum_k |\mathbf{B}_t(k_x, k_y)|^2$ (thin solid curves). The summation in $k$-space is performed in the range of $0 < k_x, k_y < 10$ to reduce thermal noises. Dashed curves show the initial turbulent spectrum, which is isotropic in $k_x$-$k_y$ space. From these panels, the turbulent magnetic fields are shown to proceed via the forward cascade perpendicular to the uniform magnetic fields. The PSD in the perpendicular direction has a clear power-law distribution. Also, the PSD in the parallel direction deviates from the initial (Gaussian) distribution for a larger $k$. It suggests that the turbulent energy is mainly transported in the perpendicular direction, but some part of the energy (a few percent) is transported in the parallel direction.

To visualize the time evolution of the anisotropy, we show the time evolution of a quantity $\theta$ in Figure 4,

$$\tan^2 \theta = \frac{\int d k_x d k_y k_x^2 |\mathbf{B}_t|^2}{\int d k_x d k_y k_y^2 |\mathbf{B}_t|^2}$$

(see Shebalin et al. 1983; Saito et al. 2008). As already mentioned, the turbulent eddies interact each other with time more frequently for larger $\epsilon$, transferring the turbulent energy to smaller scales. When $t\omega_{pe} \lesssim 100$, the anisotropic cascades proceed faster for larger $\epsilon$ because the cascade rate is an increasing function of the fluctuation. The turbulent energy is then transferred to the smaller scale. After that, the anisotropy $\theta$ saturates almost at a constant level. The saturation level is shown to be highest for $\epsilon = 0.4$ among the computed models. This is because too much initial turbulent energy (like $\epsilon = 0.8$)
Figure 3. Upper panel: power spectral density of the turbulent magnetic fields, log $|\hat{B}_t / B_0|^2$ at $t\omega_{pe} = 1350$ for $\epsilon = 0.1, 0.4, 0.8$ from left to right, respectively. The lower panels show the one-dimensional, cumulative power spectral density. Thick and thin solid curves show $|\hat{B}_t(k_x) / B_0|^2$ and $|\hat{B}_t(k_y) / B_0|^2$, respectively, while dashed curves show the initial spectra of the turbulent magnetic energy. (A color version of this figure is available in the online journal.)

Figure 4. Time evolution of the anisotropy $\tan^2 \theta$ for $\epsilon = 0.8$ (thick solid), $\epsilon = 0.4$ (thin solid), and $\epsilon = 0.1$ (dotted curve).

3.3. Electron Heating

Now we are in a position to evaluate the temperature of electrons heated by the particle–wave interactions mentioned above. In doing so, we utilize the energy–momentum tensor as

$$ T^{\mu \nu} = \frac{1}{m_e} \int d^3 p f(p) p^\mu p^\nu / e, \quad (9) $$

where $e$, $p^\mu$, and $f$ are the electron energy, the electron four momentum, and the electron distribution function, respectively. It should be noted that the temperature obtained from this equation is evaluated in the observer frame, while it should be naturally defined in the comoving frame. However, this simplification is good enough in our case because the averaged velocity of electrons in mostly random motions is much smaller than the speed of light.

Figure 5 shows the time evolution of the electron temperature. It can be seen that the electron temperature parallel to the uniform magnetic field ($T_{yy}$, solid curves) increases with time faster than the perpendicular one ($T_{xx}$, dashed curves), and then $T_{yy}$ becomes larger than $T_{xx}$. It is also shown that the electrons are heated more rapidly for a larger $\epsilon$ because the turbulent energy cascade proceeds faster. Figure 6 shows the time evolution of the ratio of $T_{yy}$ and $T_{xx}$. Before $t\omega_{pe} < 200$, the anisotropy of the electron temperature increases with time faster for a larger
\[ f(t = 1350) - f(t = 0) \]

\[ \epsilon = 0.1 \]

\[ \epsilon = 0.8 \]

\[ B_0 \]

\[ \beta_y \]

\[ \beta_x \]

\[ \beta_y \]

\[ \beta_x \]

\[ f(t = 1350) - f(t = 0) \]

\[ \epsilon = 0.1 \]

\[ \epsilon = 0.8 \]

\[ B_0 \]

\[ \beta_y \]

\[ \beta_x \]

\[ \beta_y \]

\[ \beta_x \]

\[ f(t = 1350) - f(t = 0) \]

\[ \epsilon = 0.1 \]

\[ \epsilon = 0.8 \]

\[ B_0 \]

\[ \beta_y \]

\[ \beta_x \]

\[ \beta_y \]

\[ \beta_x \]

\[ f(t = 1350) - f(t = 0) \]

\[ \epsilon = 0.1 \]

\[ \epsilon = 0.8 \]

\[ B_0 \]

\[ \beta_y \]

\[ \beta_x \]

\[ \beta_y \]

\[ \beta_x \]

\[ f(t = 1350) - f(t = 0) \]

\[ \epsilon = 0.1 \]

\[ \epsilon = 0.8 \]

\[ B_0 \]

\[ \beta_y \]

\[ \beta_x \]

\[ \beta_y \]

\[ \beta_x \]

\[ f(t = 1350) - f(t = 0) \]

\[ \epsilon = 0.1 \]

\[ \epsilon = 0.8 \]

\[ B_0 \]

\[ \beta_y \]

\[ \beta_x \]

\[ \beta_y \]

\[ \beta_x \]

\[ f(t = 1350) - f(t = 0) \]

\[ \epsilon = 0.1 \]

\[ \epsilon = 0.8 \]

\[ B_0 \]

\[ \beta_y \]

\[ \beta_x \]

\[ \beta_y \]

\[ \beta_x \]

\[ f(t = 1350) - f(t = 0) \]

\[ \epsilon = 0.1 \]

\[ \epsilon = 0.8 \]

\[ B_0 \]

\[ \beta_y \]

\[ \beta_x \]

\[ \beta_y \]

\[ \beta_x \]

\[ f(t = 1350) - f(t = 0) \]

\[ \epsilon = 0.1 \]

\[ \epsilon = 0.8 \]

\[ B_0 \]

\[ \beta_y \]

\[ \beta_x \]

\[ \beta_y \]

\[ \beta_x \]
Takahashi, Kotake, & Yasutake

Figure 8. Electron distribution function at $t \omega_{pe} = 1350$ for $\epsilon = 0.1$ (left) and for $\epsilon = 0.8$ (right) on the plane of $\beta_x - \beta_y$.

(A color version of this figure is available in the online journal.)

Figure 9. Dispersion relation of parallel propagating waves of $E_x$. The color shows the numerical result for $\epsilon = 0.1$. The dashed curves represent the dispersion relations for the $R$-mode (top and bottom) and $L$-mode (middle) obtained from linear analysis.

(A color version of this figure is available in the online journal.)

4. SUMMARY AND DISCUSSIONS

Bearing in mind the application to the outer crust of NSs, we investigated the turbulent magnetic field decay by means of fully relativistic 2D PIC simulations. In the numerical simulations, the initial magnetic fields were set to be composed of both the uniform magnetic fields that model the global fields penetrating the NS and the turbulent magnetic fields that would originate from the Hall cascade of the large-scale turbulence. We showed that the turbulent whistler cascade transports the magnetic energy preferentially in the direction perpendicular to the uniform magnetic fields. It was also found that the distribution function of electrons becomes anisotropic because electrons with lower energies are predominantly heated in the direction parallel to the uniform magnetic fields due to the Landau resonance, while electrons with higher energies are heated mainly by the cyclotron resonance that makes the distribution function isotropic for the high energy tails. Furthermore, we pointed out that the degree of anisotropy takes on the maximum value as a function of the initial turbulent magnetic energy. This is because too much initial turbulent energy disturbs the direction of the uniform magnetic fields, acting to smear out the anisotropy. The findings of this paper suggest that the particle–wave interactions via the Landau resonance and the cyclotron resonance, an alternative to the conventional particle–particle collisions via the Ohmic dissipation, would be a pivotal dissipation mechanism in the outer crust of NSs.

The initial turbulent magnetic fields assumed in this study are relatively larger compared to the previous work (Saito et al. 2008; Gary et al. 2008). For $\epsilon < 0.4$ in our simulation, the obtained results of the anisotropic cascade and the electron heating are basically consistent with the previous studies. As already mentioned, the anisotropy of the temperature decreases as $\epsilon$ increases furthermore since the turbulent fields disturb the direction of the uniform magnetic fields. Although the temperature becomes nearly isotropic, it should be noted that the cascading process itself is still anisotropic. According to Cho & Lazarian (2004, 2009), the small eddy (with size $\sim l$) of the turbulent magnetic fields interacts not with the global magnetic fields but with the “local” mean magnetic fields $B_L$. Thus, even when the turbulent magnetic field energy is much smaller than that of the uniform magnetic fields (especially $B_0 = 0$), the cascading process should be locally anisotropic along $B_L$. Such diffusion processes lead to the anisotropic heating of electrons; however, the resulting electron temperature is on average isotropic, which should be the case obtained for high $\epsilon$ in our simulation.

In the conventional model of the magnetic field decay, the magnetic fields are considered to be dissipated through the collisional process, namely, via the Ohmic dissipation, while in this paper we proposed that the magnetic fields are dissipated due to the Landau and cyclotron dissipation, leading to the plasma heating in the collisionless regime. Since the wave–particle interactions occur on a very small scale ($\sim$gyroradius), the corresponding timescale ($\sim \omega_{ce}^{-1}$) is much shorter than that of the Ohmic dissipation, and it is instantaneous compared to the expected magnetic field decay timescale in NSs. This means that the decay timescale should be determined by the cascading processes of the magnetic fields above the electron inertia scale. Although the PIC simulations have an advantage in its capability to determine the dissipation processes consistently, it is still computationally too expensive to perform the PIC simulations covering over the wide spatial range required to estimate the timescale. For that purpose, we think it important to perform the EMHD simulations including a phenomenological diffusivity which is adjusted to mimic the dissipation obtained in this study. Although this is apparently beyond the scope of this paper, we regard it as one of the most important tasks, which we plan to investigate as a sequel of this study.
So far there have been extensive work that focuses on the origin of the observed anisotropy on the NSs’ surface temperature (e.g., Geppert et al. 2004; Pons & Geppert 2007; Aguilera et al. 2008; Pons et al. 2009). Unfortunately, it is hard for us to do so immediately because the anisotropy obtained in the current simulation is confined to a very small scale (collisionless scale). To clarify how the particle–particle collisions in the larger scales redistribute the electron distribution function and what the resulting global temperature could be, one may need a new simulation technique which bridges the PIC simulation and the global (E)MHD simulations, which is not an easy job at present. In addition to the NSs’ surface, we speculate that the anisotropy due to the Landau resonance could also be important in the pulsar atmosphere because the plasma density is so small that the plasma there is expected to be globally collisionless. The consideration of the electron heating due to the Landau and cyclotron resonances might make the origin of the observed temperature anisotropy of NSs (Zavlin 2007; Haberl 2007; Nakagawa et al. 2009) less mysterious, which we are going to study one by one as an extension of this study.

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