ELECTRON ACCELERATION AT A CORONAL SHOCK PROPAGATING THROUGH A LARGE-SCALE STREAMER-LIKE MAGNETIC FIELD

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Received 2015 December 14; accepted 2016 February 25; published 2016 April 5

ABSTRACT

Using a test-particle simulation, we investigate the effect of large-scale coronal magnetic fields on electron acceleration at an outward-propagating coronal shock with a circular front. The coronal field is approximated by an analytical solution with a streamer-like magnetic field featuring a partially open magnetic field and a current sheet at the equator atop the closed region. We show that the large-scale shock-field configuration, especially the relative curvature of the shock and the magnetic field line across which the shock is sweeping, plays an important role in the efficiency of electron acceleration. At low shock altitudes, when the shock curvature is larger than that of the magnetic field lines, the electrons are mainly accelerated at the shock flanks; at higher altitudes, when the shock curvature is smaller, the electrons are mainly accelerated at the shock nose around the top of closed field lines. The above process reveals the shift of the efficient electron acceleration region along the shock front during its propagation. We also find that, in general, the electron acceleration at the shock flank is not as efficient as that at the top of the closed field because a collapsing magnetic trap can be formed at the top. In addition, we find that the energy spectra of electrons are power-law-like, first hardening then softening with the spectral index varying in a range of \(-3\) to \(-6\). Physical interpretations of the results and implications for the study of solar radio bursts are discussed.

Key words: acceleration of particles – shock waves – Sun: coronal mass ejections (CMEs) – Sun: radio radiation

1. INTRODUCTION

Shock waves are believed to be efficient particle accelerators in the universe. They are the primary source of solar energetic particles (SEPs), although the nature of the driver of the shocks in the solar corona is still under debate. The coronal shock could be a piston shock driven by a coronal mass ejection (CME) or a blast wave ignited by a flare (see the recent review by Vršnak & Cliver 2008). Shock-induced energetic electrons can excite electromagnetic radiation in the radio wavelength via a plasma emission mechanism (Ginzburg & Zhelezniakov 1958; Nelson & Melrose 1985). Type II solar radio bursts, appearing as narrow frequency bands drifting slowly in the dynamic spectra recorded by radio spectrometers, serve as signatures of shocks propagating outward in the solar corona and interplanetary space.

Particle acceleration by shocks has been intensively studied for decades. The diffusive shock acceleration (DSA) theory is one of the most important mechanisms, in which ions can gain energy from resonant interactions with magnetic turbulence or plasma waves (Axford et al. 1977; Bell 1978; Blandford & Ostriker 1978). The efficient acceleration of electrons, on the other hand, is more difficult since the gyroradii of low-energy electrons are very small compared to those of ions. At quasi-perpendicular shocks, electrons can be accelerated by gradient drift in the magnetic field at the shock along the motional electric field, known as shock drift acceleration (SDA, Armstrong et al. 1985) or fast Fermi acceleration (Wu 1984). However, for a single reflection at a planar shock in the scatter-free limit, the energy gain has been shown to be very limited (e.g., Ball & Melrose 2001). Thus, multiple reflections at the shock are required for efficient acceleration. Some non-planar effects, such as shock ripples and magnetic fluctuations, have been demonstrated to be capable of enhancing electron acceleration in recent numerical simulations (e.g., Burgess 2006; Guo & Giacalone 2010, 2012). For example, Guo & Giacalone (2010) studied the effect of large-scale fluctuations on the acceleration of electrons at perpendicular shocks and found that the large-scale braiding of field lines allows electrons to cross the shock front repeatedly, and small-scale shock ripples also contribute to the acceleration by mirroring and trapping electrons.

In studying electron acceleration at coronal shocks, the effects of the large-scale coronal magnetic field have not received much attention. Closed magnetic structures such as coronal loops, which are ubiquitous in the lower solar corona on various scales, are the fundamental building blocks of the coronal magnetic field at \(<2\text{-}3 R_\odot\). Therefore, if a coronal shock is generated at a sufficiently low height, it can either propagate through closed field lines above the active region or cross nearby closed field regions as it expands laterally. In both cases, an electron trapping geometry can be formed, similar to a collapsing magnetic trap in solar flare models (e.g., Somov & Kosugi 1997; Nishizuka & Shibata 2013). This kind of configuration is usually thought to be efficient for electron acceleration. Thus, a large-scale closed field could potentially be an efficient location for electron acceleration at coronal shocks and relevant solar radiation such as metric type II and type IV radio bursts.

Although it is widely accepted that energetic electrons responsible for type II radio bursts are associated with shock waves, where and how these electrons are accelerated is not yet been completely understood. The observational results for type IIs in previous studies indicate that both the shock nose and the shock flank could be the source region (e.g., Mancuso...
of the closed and open field of the Sun. The black lines represent a streamer-like coronal magnetic field described by an analytical model, and the dashed line shows the boundary of the closed and open field and the current sheet above. A coronal shock with a circular front, centered at the solar surface \( z = 0, r = 1 \ R_e \), expands outward radially with a constant speed of \( U_{sh} \approx 1000 \ \text{km s}^{-1} \), as shown by the thick blue circles. Electrons are injected immediately upstream of the shock, as the shock radius \( R_{sh} \) increases from \( 0.3 \ R_e \) to \( 1.8 \ R_e \). The angle \( \theta_s \) is defined as the angle from one point on the shock front to the streamer axis, with the center being the shock epicenter. A local shock frame \( (x', y', z') \) at a certain point on the shock front is also shown.

Figure 2. Variations of \( \theta_{bn} \) (the angle between the shock normal vector and upstream magnetic field) along the shock front at different instants when the shock reaches various heights from the epicenter \( (R_{sh} = 0.4, 0.5, 0.6, 0.8, 1.0, 1.2, 1.4, 1.6 \ R_e, \) corresponding to \( t = t_1, t_2, \ldots, t_8 \), respectively). The horizontal axis \( \theta_s \) is the angle of one point on the shock front to the streamer axis (see Figure 1). The red asterisks indicate the borders of closed and open field lines at \( t_4 \)–\( t_7 \) (see Figure 1 for \( t_4 \)). The arrows mark the tangent points between the shock flank and the field line at times of \( t_2, t_3, \) and \( t_4 \).

\& Raymond 2004; Cho et al. 2008). By combining radio imaging data from the Nançay Radioheliograph (NRH) and high-quality imaging data of solar eruptions from the Atmospheric Imaging Assembly (AIA)/Solar Dynamics Observatory (SDO), several new advances have recently been made (e.g., Bain et al. 2012; Zimovets et al. 2012; Carley et al. 2013; Feng et al. 2015; Zimovets & Sadykov 2015). Despite all this progress, the exact origin of type II emission and the physics accounting for its fine structure are still unclear, partly due to the limited capability of solar radio imaging observations (e.g., Du et al. 2014, 2015).

Many recent works have suggested that the interaction region between CME shocks and streamers could possibly be an important source of both metric and interplanetary type IIs (e.g., Reiner et al. 2003; Mancuso & Raymond 2004; Cho et al. 2007, 2008; Feng et al. 2012, 2013; Kong et al. 2012, 2015; Shen et al. 2013; Chen et al. 2014; Magdalenić et al. 2014; Eselevich et al. 2015). Possible explanations for these results are twofold. First, the streamer region features relatively slow plasma outflows and Alfvénic speeds, compared to those of its surroundings, mainly due to its much higher plasma density. Therefore, the streamer region is expected to facilitate the formation or enhancement of a solar-eruption-driven shock. Second, the shock geometry is likely to be more quasi-perpendicular in the shock-streamer interaction region when the shock encounters the streamer from the flank. Both factors can make the electron acceleration at a shock more efficient. Lately, Kong et al. (2015) studied the effect a streamer’s closed field has on shock-electron acceleration using a test-particle model consisting of a planar shock and an analytical streamer field. It was found that the closed field of the streamer plays the role of an electron trap through which electrons are sent back to the shock front multiple times and are repetitively accelerated through the SDA mechanism. This is likely a fundamental effect considering the fact that a majority of solar eruptions originate from closed field structures above active regions.

In this paper, we extend our previous model by considering a coronal shock with a curved shock geometry propagating through a large-scale, streamer-like coronal field to further explore its effect on shock-electron acceleration. We describe the setup of our numerical model in Section 2 and the simulation results are given in Section 3. Our conclusions and discussion are presented in the Section 4.

2. NUMERICAL MODEL

In this section, we introduce the setup of the numerical model, which is axisymmetric and consists of a streamer-like coronal magnetic field and an outward-propagating non-planar coronal shock. A cross-section view of the background magnetic field and the shock morphology is shown in Figure 1. Electrons are treated as test particles and are assumed to have a negligible effect on the shock fields and the evolution. The equations of motion of the electrons in prescribed electric and magnetic fields are numerically integrated.

Following Kong et al. (2015), the coronal magnetic field is approximated by an analytical solution of a streamer-like field (Low 1986). It describes an axisymmetric magnetic structure containing both closed magnetic arcades and open field lines with a current sheet in a spherical coordinate \( (r, \theta) \). This coronal field model has been used in previous corona and solar wind modeling (e.g., Chen & Hu 2001; Hu et al. 2003b, 2003a). In this study, the magnetic field strength in the polar region on the solar disk is set to 10 G. The simulation domain is given by \( r = [1.0, 3.0] \ R_e \), which is much larger than that in Kong et al. (2015). The magnetic topology of the region of interest is shown in Figure 1 in Cartesian coordinate \( (x, z) \). The
z-axis represents the rotation axis of the Sun, the x-axis is in the solar equatorial plane parallel to the streamer axis, and the y-axis completes the right-handed triad with the solar center being at the origin. The black lines represent the magnetic field lines and the black dashed line denotes the outermost closed field line and the current sheet above. The height of the streamer cusp is taken to be $2.5 \, R_E$. The $y$ component of the magnetic field $B_y$ is set to be 0.

We consider a shock with a circular front propagating outward with a constant radial expansion speed of $U_{sh} \sim 1000 \, \text{km s}^{-1}$, as shown by the thick blue circles in Figure 1. The center of the shock is located at the solar surface in the equatorial plane ($x = 1 \, R_E, z = 0$). The shock is assumed to form with a shock radius $R_{sh} = 0.3 \, R_E$ and the shock compression ratio $X_{sh}$ is set to be 2.5. Note that in Kong et al. (2015), we took $X_{sh}$ to be 4, which is the upper limit for a very strong shock. In recent white-light and EUV observations, a sheath region with enhanced intensity ahead of the CME is commonly identified as a signature of coronal shocks (e.g., Vourlidas et al. 2003; Ma et al. 2011), and the density compression at the shock was deduced to be no more
than 3 (e.g., Bemporad et al. 2014; Susino et al. 2015). The background plasma in the simulation domain is assumed to be at rest. At each point along the shock front, there exists a local shock frame \((x', z')\). The \(x'\)-axis is parallel to the local shock normal, pointing from the shock center radially to that point, and the \(z'\)-axis is perpendicular to the \(x'\)-axis and lies in the \((x, z)\) plane (see Figure 1). In the local shock frame, where the shock is at \(x' = 0\), the plasma carries the magnetic field flow from \(x' < 0\) (upstream) with a speed of \(U_1 \sim U_{sh}\) to \(x' > 0\) (downstream) with a speed of \(U_2 \sim U_{sh}/\delta_{sh}\). The flow speed close to the shock is given by a hyperbolic tangent function \(U(x') = (U_1 + U_2)/2 - (U_1 - U_2) \tanh(x'/\delta_{sh})/2\), where \(U_1\) and \(U_2\) are the upstream and downstream flow speeds in the shock frame. \(\delta_{sh}\) is the shock thickness and is taken to be 0.01 \(U_1/\Omega_{ci}\), where \(\Omega_{ci}\) is the proton gyrofrequency given by \(B_0 \sim 0.2\) G which approximates the average magnitude of magnetic field in the simulation domain. The upstream magnetic field is described by the analytical solution and we use the ideal MHD shock jump conditions in the local shock frame to determine the downstream magnetic field. The motional electron field is deduced using the ideal MHD approximation \(E = -U \times B/c\).

When the shock starts to propagate outward, from \(R_{sh} = 0.3\) \(R_e\) to 1.8 \(R_e\), electrons with an initial energy of \(E_0 = 300\) eV are continuously injected into the immediate region upstream

Figure 4. Same as Figure 3, but for an electron being trapped and accelerated at the loop top of the closed field. Note that in panel (d) the profile of the electron’s drift along the \(y\) direction (the black line) has been shifted 5 s to the left.
of the shock (at $x' = -10 \frac{U_1}{\Omega_{ci}}$) at a constant rate. Their initial pitch angles are given randomly. For each electron, the equation of motion under the Lorentz force is solved in the lab frame. The electron mass is taken to be $\frac{1}{1836}$ of the proton mass. The numerical technique used to integrate the electron trajectories is the Bulirsch–Stoer method (Press et al. 1986), which has been widely used in calculating particle trajectories (e.g., Giacalone 2005; Guo & Giacalone 2015). The algorithm uses an adjustable time-step method based on the evaluation of the local truncation error. It is highly accurate and has been tested to conserve particle energy to a good degree. In this study, a total of $10^7$ electrons are injected. When an electron moves out of the simulation domain or reaches a distance of $10^4 U_1 / \Omega_{ci}$ downstream of the shock, we stop tracking it and terminate the calculation. An ad hoc pitch-angle scattering is included to mimic the effects of electron interaction with coronal plasma turbulence, kinetic waves on electron and ion scales, and Coulomb collisions (e.g., Marsch 2006). This is done by randomly changing the electron pitch angle every $\tau = 10^4 \Omega_{ci}^{-1}$ (~5 s).

3. SIMULATION RESULTS

The particle acceleration efficiency strongly depends on the angle ($\theta_{Bn}$) between the shock normal vector and the upstream magnetic field. Generally, a quasi-perpendicular shock...
($\theta_{Bn} > 45^\circ$) favors the acceleration of electrons (e.g., Holman \& Pesses 1983; Wu 1984). So, before presenting the main simulation results, we first examine the evolution of the shock geometry as the shock propagates outward. Figure 2 shows the variation of $\theta_{Bn}$ along the shock front at eight different instants with different colors when the shock reaches various heights from the epicenter ($R_{sh} = 0.4, 0.5, 0.6, 0.8, 1.0, 1.2, 1.4, 1.6$ $R_e$, corresponding to $t = t_1, t_2, \ldots t_8$, respectively). The horizontal axis $\theta_s$ is defined as the angle of one point on the shock front to the streamer axis, with the center being the shock epicenter (see Figure 1). We can see that the shock geometry changes substantially along the shock front with its propagation. At earlier times ($t_1, t_2, \text{and } t_3$), the shock only moves through the closed field, while at later times (after $t_4$) the shock crosses both the closed and open fields. The locations of the boundary between the closed and open fields at different instants ($t_4$–$t_7$) are marked by red asterisks in Figure 2. In addition, because of the variation of shock radius and streamer field topology, the tangent points of the shock front and field lines (where $\theta_{Bn} \sim 90^\circ$) change with time. It is evident that at the apex of the closed loops (or the shock nose), the shock is always perpendicular. At $t_2$ and $t_3$, another perpendicular shock region exists at the flank within the closed field; at $t_4$, there is also a perpendicular shock region at the flank but in the open field. The locations of the tangent points are indicated by arrows in Figure 2. We see that the shock geometry here is significantly different from that in Kong et al. (2015) due to the use of a more realistic shock with a circular front and a much larger streamer-like field. Note that the evolution of shock geometry can also be seen from Figure 7 as will be shown later.

The simulation results demonstrate that injected electrons can be accelerated to energies up to 440 $E_0$. The maximum
energies obtained by the electrons at different instants ($t_1$–$t_8$, corresponding to different shock heights) are 52 $E_0$, 44 $E_0$, 202 $E_0$, 195 $E_0$, 114 $E_0$, 151 $E_0$, 62 $E_0$, and $20 E_0$, respectively. We will show the electron trajectories and energy evolution of three representative cases, with final energies $\geq 20 E_0$, to demonstrate the acceleration mechanism and affecting factors. The electron is injected and accelerated at the shock flank within the closed field in the first case (Figure 3), trapped and accelerated at the loop top of the closed field in the second case (Figure 4), and injected and accelerated at the shock flank in the open field in the third case (Figure 5).

We first analyze the acceleration of an electron at the shock flank within the closed field, which achieves a final energy of $\sim 35 E_0$. In Figure 3, panel (a) shows the electron’s trajectory in the prescribed magnetic field in the lab frame. One can see that the electron basically follows a specific field line, since here the third component of the magnetic field ($B_y$) is set to be 0. It is consistent with previous simulation results that in a magnetic field that has at least one ignorable coordinate, the motion of the charged particle is restricted to the original field line (see, e.g., Guo & Giacalone 2013 and references therein). At the end of this calculation, the electron goes down to the solar surface. In panel (b), the electron’s trajectory is illustrated in three-dimensional (3D) coordinates. The blue arrow points to its injection position. The electron displays a gradient-$B$ drift in the $y$ direction whenever it is reflected at the shock. Also, we can see its back and forth motion along the closed field line. In panel (c), it shows the distance of the electron away from the shock front (i.e., its position $x'$ in the local shock frame) over time. The time starts from the shock formation when $R_{sh} = 0.3 R_\odot$ and $t = t_0$. We can see that the electron is reflected about eight times at the shock front (where $x' = 0$, denoted by the horizontal blue line). An encounter of the electron with the shock may be due to the effect of pitch-angle scattering or to being caught up by the shock. In panel (d), we present the temporal evolution of the electron’s drift in the $y$ direction (the black line) and the electron energy (the red line). It shows that a fast drift of the electron in the $y$ direction is accompanied by a simultaneous sharp increase of its energy. Therefore, the electron gains energy mainly via the SDA mechanism.

Figures 4 and 5 present the simulation results for the other two cases. Compared with the shock-field configuration in Figure 3, which only has one mirroring point at the shock flank, the configuration shown in Figure 4 has two mirroring points because the shock front can intersect the same field line at two different points. With the outward propagation of the shock, the two mirroring points approach each other, giving rise to a collapsing magnetic trap geometry. Since the length of the trap gets shorter and the shock geometry becomes more perpendicular over time, the electron acceleration becomes more efficient. We see that the electron’s energy increases impulsively in the final stage. It takes only $\sim 2$ s for its energy to increase from $\sim 10 E_0$ to $\sim 70 E_0$. The shock-field configuration shown in Figure 5 is similar to that in Figure 3 while the electron is injected into the open field and can only encounter the shock at one end. The electron experiences a generally gradual energy growth, and eventually escapes along the open field line. Due to a nearly perpendicular shock geometry and the effect of pitch-angle scattering, it was accelerated to $\sim 20 E_0$. The two electrons shown in Figures 4 and 5 also gain energy via the SDA mechanism. The difference in the variation of energy as a function of time for these three electrons indicates that the large-scale shock-field configuration can strongly affect the efficiency of acceleration.

Now we examine the distribution of injection positions for all of the electrons that have been accelerated. In Figure 6, panels (a)–(c) show a color-coded representation of the number of electrons with a final energy of $>5 E_0$, $>10 E_0$, and $>20 E_0$, respectively. One can see from panel (a) that many electrons...
injected in the open field can be accelerated to $>5\ E_0$. In contrast, as shown in panels (b) and (c), for electrons having reached higher energies, almost all of them are injected in the closed field region. This point can also be inferred from Figure 7, as we will show later. This confirms the importance of the large-scale closed field on shock-induced electron acceleration, as pointed out and tested using a planar shock model in Kong et al. (2015).

In the right side of Figure 6, the distributions along the shock front at different instants ($t_1$–$t_8$) are shown as histograms with different colors, binned over 1 (in panels (d) and (e)) or 2 (in panel (f)) degrees. The borders between the closed and open field at $t_4$–$t_7$ (see Figure 2) are marked by triangles in the same color as the corresponding histograms. We see that at earlier times ($t_1$ and $t_2$, i.e., lower shock altitudes), the histogram profiles present a single peak located at the shock flank. Later, at $t_3$ and $t_4$, two peaks are observed, one at the flank and the other around the shock nose with the latter being much higher than the former. At $t_5$ and $t_6$, the peak around the shock nose remains much more prominent than the flank peak, which becomes almost invisible for higher energies (see panels (e) and (f)). From the results, we again see that electrons injected within the closed field can be accelerated efficiently, while only a small number of the electrons injected in the open field at the shock flank can be accelerated to an energy $>20\ E_0$. As explained earlier, electron acceleration within the closed field is strongly enhanced by the trapping effect of the field, while that at the shock flank is mainly due to the quasi-perpendicular shock geometry there.

**Figure 8.** Positions of energetic electrons with different energies ($>5\ E_0$, $>10\ E_0$, and $>20\ E_0$, respectively) as the shock reaches various heights. In panels (a)–(c), the shock fronts at different times ($t_1$–$t_8$) are indicated by blue circles, and electrons are shown by scattering dots with different colors. On the right side of the figure, the distribution along the shock front is shown by histograms with different colors. The borders between the closed and open field at $t_4$–$t_7$ (see Figure 2) are marked by triangles with the same color as the histograms.
To further explore the underlying physics, in Figure 7, we show the injection positions of energetic electrons as red scattering dots superposed with shock fronts and magnetic field lines. Remember that electrons are always injected near the shock front, and so the red dots close to the shock front represent the injection position of electrons at the corresponding time. At earlier times (t1 and t2), energetic electrons are injected near (but not at) the tangent points of the shock and field lines. For example, at t1, injection positions lie in the region \(70^\circ (53^\circ) \gtrsim \theta_s \gtrsim 30^\circ\) where the shock angle \(74^\circ (81^\circ) \lesssim \theta_{Bo} \lesssim 87^\circ\) for electrons \(>10\ E_0\). This configuration looks similar to the foreshock morphology proposed in theoretical models of interplanetary type IIs (e.g., Knock et al. 2001, 2003). At t3, there exist two injection regions of energetic electrons. Aside from the region bound to earlier times (pointed by the green arrow), there appears a new region enclosed by the shock front and the top of closed field lines (pointed by the black arrow) corresponding to the two peaks observed in Figure 6 (the thin red line). At t4, there are also two parts of the injection positions, one part confined by the closed field (the black arrow) and the other part in the open field at the shock flank (the green arrow). In addition, one can see that if electrons are injected very close to the streamer axis, then they cannot be accelerated efficiently. It is found that the shock angle \(\theta_{Bo}\) near the streamer axis (at t3, where \(\theta_s \lesssim 12^\circ\), and at t4, where \(\theta_s \lesssim 5^\circ\)) is \(>86^\circ\). As pointed out in previous studies, particles cannot be reflected by the shock when \(\theta_{Bo}\) is near \(90^\circ\) (e.g., Holman & Pesses 1983; Ball & Melrose 2001).

Now we examine the positions of energetic electrons as the shock reaches different heights. In Figure 8, panels (a)–(c) show the positions of electrons with different energies, \(>5\ E_0\), \(>10\ E_0\), and \(>20\ E_0\). The shock fronts at different heights are shown by the blue circles and the electrons are represented by the scattering of dots with different colors. For better clarity, only the distribution at eight instants (t1–t8) are presented for each energy level. In the right side of the figure, the distribution of the accelerated electrons along the shock front is shown by histograms binned over 1 (in panels (d) and (e)) or 2 (in panel (f)) degrees. We see that for different energies and different instants, the distributions look quite different. For electrons with lower energies (\(>5\ E_0\)), their positions are more dispersive, appearing both in the open and closed fields; for higher energies (\(>10\ E_0\)), the electrons are more concentrated, mostly within the closed field. Meanwhile, the electron concentration site changes with the outward propagation of the shock, from the shock flank (or the foreshock region) to the region between the shock nose and the top of the closed field. At earlier times (t1 and t2), the energetic electrons mainly concentrate at the foreshock region along the flank. At t3, the energetic electrons start to appear in the tops of the closed loops above the shock nose and remain there later on (before t8). At or after t8, when the shock passes over the streamer cusp, very few electrons are accelerated to high energies, again indicating the importance of the large-scale closed field in shock-induced electron acceleration.

The temporal variation of the distribution of energetic electrons can also be seen from the histograms (right panels of Figure 8). We see that significantly fewer electrons can gain energies \(>20\ E_0\) at t1 and t2 than at later times, suggesting that electron acceleration at the shock flank is not as efficient as that at the top of the closed field, since there is only one mirroring point at the shock flank, while a collapsing
magnetic trap can be formed as the shock sweeps through the tops of the closed field lines. This is consistent with our previous analysis of the three representative electrons. Note that electrons with energies \(> 20 \, E_0 \) (6 keV), as shown in panel (c), are capable of exciting Langmuir waves and radio emission, and so the region with these energetic electrons could potentially be the source of metric radio bursts, such as type II and type IV bursts. The simulation results indicate that the radio-burst source shifts along the shock front over time while moving outward.

The above analysis indicates that the large-scale shock-field configuration plays an important role in the efficiency and location of electron acceleration. To illustrate this more clearly, we present schematics of electron acceleration in different shock-field configurations in Figure 9. The thick blue curve denotes the shock front, the black lines show magnetic field lines, and the red scattering of dots represents energetic electrons. Panel (a) corresponds to the configuration at low shock altitudes (t1 and t2) where the electrons are injected and accelerated at the shock flank, with only one mirroring point at

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**Figure 10.** Energy spectra of electrons at eight instants (t1–t8) corresponding to different shock heights. Panel (a) presents the energy spectra at t1–t4 and panel (b) those at t5–t8. The vertical coordinate represents the number of energetic electrons (\(\Delta N\)) in a certain energy range (\(\Delta(E/E_0)\)). For each instant, all of the energetic electrons in the simulation domain (as shown in Figure 8) are included to determine the spectrum. The starting points of these energy spectra are normalized to the same values in panels (c) and (d).
the shock. As noted above, this configuration is similar to the foreshock morphology proposed in theoretical models of interplanetary type IIs. The configuration around $\Omega = \Omega_1$, when electrons can be accelerated both at the flank and the nose, is shown in panel (b). Panel (c) illustrates the configuration at higher shock altitudes ($\Omega = \Omega_2$). Electrons are trapped by the closed field of the streamer, and energetic electrons are concentrated at the loop tops. Comparing these configurations, we find that the curvature difference between the shock front and closed field lines is important to the acceleration and distribution of energetic electrons. At low shock altitudes, the curvature radius of the shock is smaller than that of the field lines; meanwhile, at higher altitudes, the radius of the shock becomes larger than that of the closed field lines. Panel (d) shows a specific configuration when electrons are accelerated at the shock flank in the open field (around $\Omega_4$) where the quasi-perpendicular shock geometry is crucial.

Figure 10 shows the temporal evolution of the electron energy spectra as the shock propagates outward. Panel (a) presents the energy spectra at $t_1$ and $t_2$ and panel (b) at $t_3$ and $t_4$. The vertical coordinate represents the number of energetic electrons ($\Delta N$) in a certain energy range ($\Delta(E/E_0)$). For each instant, all of the energetic electrons in the simulation domain (as shown in Figure 8) are included to determine the spectrum. To display the temporal evolution more clearly, in panels (c) and (d) we normalize the starting points of these energy spectra to the same value. We see that the spectra at $t_1$ and $t_2$ can be regarded as a double power law, while at later times ($t_3$ and $t_4$) they are more similar to a single power law with the high-energy part of the spectra hardening. The spectral index is about $-3$ to $-4$. Later, from $t_4$ to $t_5$, the energy spectra become softer again, with the spectral index decreasing gradually to about $-6$ at $t_5$. The hardening of the energy spectra, when the main source region of the energetic electrons moves from the shock flank to the shock nose, is consistent with our main result that electron acceleration is more efficient when electrons are trapped between the shock nose and the closed field, in comparison with the shock-flank situation. With the outward propagation of the shock, the electron trapping area within the closed field decreases with time, and the average value of the shock angle $\theta_{\text{sh}}$ becomes smaller (i.e., the shock becomes less perpendicular). These two factors explain the spectral softening later on.

4. CONCLUSIONS AND DISCUSSION

In this paper, we perform a test-particle simulation to investigate the effect of the large-scale coronal field on electron acceleration at an outward-propagating coronal shock with a circular shock. The coronal magnetic field is approximated by an analytical solution of a streamer-like field. Due to the variation of the shock radius and streamer field topology, the shock-field configuration changes significantly with the outward propagation of the shock. We highlight the importance of the large-scale closed field to shock-induced electron acceleration, a result which is consistent with our previous study.

We find that the large-scale shock-field configuration plays an important role in the efficiency and location of electron acceleration. An important factor is the relative curvature of the shock and the magnetic field line across which the shock is sweeping. At low shock altitudes, when the shock curvature is larger than that of the magnetic field lines, the electrons are mainly accelerated at the shock flanks. The configuration is similar to the foreshock morphology proposed in theoretical models of interplanetary type IIs. At higher altitudes, when the shock curvature is smaller, the electrons are mainly accelerated and concentrated around the top of the closed field lines above the shock nose. The result in this configuration is consistent with our previous study using a planar shock with zero curvature. Our calculation reveals the shift of the efficient electron acceleration region during the shock propagation. It is found that some electrons injected into the open field at the shock flank can be accelerated to high energies as well, mainly due to the nearly perpendicular shock geometry there, but not as efficiently as those trapped in the loop top of the closed field. In addition, we find that the energy spectra of electrons are power-law-like, first hardening then softening, with the spectral index varying in a range of $-3$ to $-6$.

Energetic electrons accelerated at coronal shocks are responsible for some solar radio bursts such as type II and type IV bursts. To date, their origin remains unresolved (e.g., Zimovets et al. 2012; Carley et al. 2013; Tun & Vourlidas 2013). Theoretical studies are important to figure out where and how relevant electrons become accelerated. In this work, we highlight the possible role of the large-scale coronal field, and the closed field in particular, in shock-induced electron acceleration. It is likely a fundamental effect, considering the fact that closed magnetic structures such as coronal loops are ubiquitously present in the lower solar corona on various scales, and a majority of solar eruptions originate from closed structures above active regions. Note that the electrons are accelerated mainly through the SDA mechanism, and so the closed field should be regarded as a complementary enhancing factor (keeping electrons upstream of the shock) to shock-electron acceleration mechanisms. Other factors affecting electron acceleration, e.g., magnetic fluctuations as proposed in previous numerical simulations (e.g., Burgess 2006; Guo & Giacalone 2010, 2012), may play a role as well.

Our simulation results demonstrate that electron acceleration is closely related to the local shock-field geometry. It is possible that there is more than one region along the shock front that can produce energetic electrons and then excite radio bursts. Also, as the shock propagates outward, the main source of energetic electrons may shift along the shock front. This is important for studies using type II spectral drift to infer the propagation speed of the shock or type II source (e.g., Reiner et al. 2003; Mancuso & Raymond 2004; Ma et al. 2011; Bain et al. 2012; Kong et al. 2012; Vasanth et al. 2014). A general assumption of these studies is that the radio source propagates outward radially or in a direction along the density gradient. This may not be the case according to our calculations.

For simplicity, we have used an analytical solution to represent the coronal streamer magnetic configuration and the MHD shock jump conditions to determine the shock parameters, without considering the dynamical coupling process between the shock and the coronal plasma. Further studies should be conducted to investigate electron acceleration in a self-consistently solved coronal shock environment.

This work was supported by grants NSBJSF 2012CB825601, NNSFC 11503014, 41274175, 41331068, U1431103, and Natural Science Foundation of Shandong Province (ZR2014DQ001 and ZR2013DQ004). G.L.’s work at UAHuntsville was supported by NSF grants ATM-0847719 and AGS1135432. The numerical simulations reported here...
were carried out on the supercomputer of Shandong University, Weihai.

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