THE ROLE OF DIFFUSIVE SHOCK ACCELERATION ON NONEQUILIBRIUM IONIZATION IN SUPERNOVA REMNANTS

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ABSTRACT

We present results of semianalytic calculations which show clear evidence for changes in the nonequilibrium ionization behind a supernova remnant forward shock undergoing efficient diffusive shock acceleration (DSA). The efficient acceleration of particles (i.e., cosmic rays (CRs)) lowers the shock temperature and raises the density of the shocked gas, thus altering the ionization state of the plasma in comparison to the test-particle (TP) approximation where CRs gain an insignificant fraction of the shock energy. The differences between the TP and efficient acceleration cases are substantial and occur for both slow and fast temperature equilibration rates: in cases of higher acceleration efficiency, particular ion states are more populated at lower electron temperatures. We also present results which show that, in the efficient shock acceleration case, higher ionization fractions are reached noticeably closer to the shock front than in the TP case, clearly indicating that DSA may enhance thermal X-ray production. We attribute this to the higher postshock densities which lead to faster electron temperature equilibration and higher ionization rates. These spatial differences should be resolvable with current and future X-ray missions, and can be used as diagnostics in estimating the acceleration efficiency in CR-modified shocks.

Key words: cosmic rays – radiation mechanisms: thermal – shock waves – supernova remnants – X-rays: ISM

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1. INTRODUCTION

In young supernova remnant (SNR) shocks, the acceleration of cosmic rays (CRs) leads to a softening of the equation of state in the shocked plasma. This comes about because the diffusive shock acceleration (DSA) process turns some nonrelativistic particles into relativistic ones and because some of the highest energy relativistic particles escape from the shock. Both of these effects lead to lower postshock plasma temperatures as well as higher postshock densities (e.g., Jones & Ellison 1991; Berezhko & Ellison 1999). The ionization state of shocked gas at a particular time is dependent upon both the gas density and the electron temperature. In light of this, DSA ought to leave its imprint on the ionization structure of the shocked gas. Toward this end, we present what we believe to be the first self-consistent model for SNR evolution which includes the hydrodynamics, the effects of efficient shock acceleration, and a full treatment of the nonequilibrium ionization (NEI) balance at the forward shock (FS).

A number of young SNRs show both nonthermal and thermal emission in the region behind the FS, including SN 1006 (Vink et al. 2003; Bamba et al. 2008), Tycho (Hwang et al. 2002; Cassam-Chenaï et al. 2007), and Kepler (Reynolds et al. 2007). The thermal emission arises when the FS sweeps up the circumstellar medium (CSM) and heats it to X-ray emitting temperatures. As pointed out in Ellison et al. (2007) (hereafter DCE07), the thermal emission is often considerably fainter than the nonthermal emission, but there are certainly examples where the thermal emission is as bright or brighter than any nonthermal emission (Vink et al. 2006). In SNR RX J1713.7–3946, the lack of thermal X-ray emission is an important constraint on the ambient density and significantly impacts models for TeV emission (e.g., Slane et al. 1999; Ellison et al. 2001; Aharonian et al. 2007; Katz & Waxman 2008).

If the DSA process in young SNRs is as efficient as generally believed, with \( \gtrsim 50\%\) of the shock ram kinetic energy going into relativistic particles, nonlinear DSA will influence the SNR hydrodynamics and be important for the NEI calculations (e.g., Decourchelle et al. 2000; Ellison & Cassam-Chenaï 2005). DCE07 took the first steps in self-consistently coupling nonlinear DSA with NEI by tracking the electron temperature \( T_e \) and ionization age (defined as \( n_e t \), where \( n_e \) is the electron density and \( t \) is the time since the material was shocked) as a function of time in hydrodynamic simulations of SNRs where the FS was efficiently producing CRs and, as a result, was substantially modified from test-particle (TP) results. They found that, while both \( T_e \) and \( n_e t \) did differ between the TP and CR-modified cases, in the cases where DSA is highly efficient, the synchrotron emission in the X-ray range is considerably stronger than the thermal X-ray spectrum, and any differences in the thermal X-rays as a result of CR modification are likely to be missed. In this paper, we extend the work of DCE07 by explicitly tracking the NEI state in a CR-modified shock. The lower shock temperature and higher density that result from efficient DSA combine to shorten both the temperature equilibration and ionization equilibrium timescale, and we show that this can have a dramatic effect on the ionization structure between the FS and the (CD). Although we do not calculate the thermal X-ray emission here, the cases we study show that efficient DSA can increase the ionization fraction of important elements and possibly enhance the thermal X-ray emission.

In Section 2, we outline the changes to our model first presented in DCE07 and discuss several caveats to our approach. In Section 3, we present our examples and discuss the quantitative and qualitative effects of efficient DSA on the ionization state and SNR structure. We also show how these effects might manifest themselves in current and future X-ray observations. In Section 4, we summarize our results and outline our future enhancements to this model.
2. CR-HYDRO + NEI MODEL

Our spherically symmetric model uses the semianalytic DSA calculation developed by Amato & Blasi (2005) and Blasi et al. (2005) and is similar to that used in DCE07, except that we now calculate the NEI explicitly at every time step using plasma parameters that are continually updated as the SNR evolves. In DCE07, the NEI was calculated at the end of the simulation using average plasma parameters. We refer the reader to DCE07 for all details of the CR-hydro simulation apart for those discussed below detailing our dynamic NEI generalization.

The DSA model used here differs from that described in Ellison et al. (2007) and Ellison & Cassam-Chenaï (2005), and previous papers, in two important ways. First, we replace the “effective gamma,” \( \gamma_{\text{eff}} \), approximation with a more realistic model of the effect escaping particles have on the shock dynamics. We now explicitly remove from the shocked plasma the energy that escaping particles carry away from the FS. The ratio of specific heats of the shocked gas used in the simulations, \( \gamma_s \), is determined directly from the particle distribution function including the correct mix of relativistic and nonrelativistic particles. While the old effective gamma had the range \( 1 < \gamma_{\text{eff}} \leq 5/3 \), the ratio of specific heats \( \gamma_s \) is constrained to lie between 4/3 and 5/3. These changes in the way escaping particles are treated, and \( \gamma_{\text{eff}} \) is calculated, become important for later stages of the SNR evolution, but do not produce significant changes in times as short as 1000 yr. The results reported in Ellison et al. (2007) are not modified significantly by these changes.

The second difference is that instead of specifying a fixed injection parameter, \( \chi_{\text{inj}} \) (this is \( \xi \) in Equation (25) in Blasi et al. 2005), which then determines the acceleration efficiency, we now specify a fixed DSA efficiency, \( \epsilon_{\text{DSA}} \), and then determine \( \chi_{\text{inj}} \) accordingly. This change makes the parameterization of the acceleration efficiency more transparent but does not change the basic approximation that is made.

The semianalytic DSA model we use does not calculate the acceleration efficiency self-consistently based on the Mach number, the available acceleration time, and other relevant shock parameters; rather we parameterize the efficiency by \( \chi_{\text{inj}} \), and the model then determines the shock structure self-consistently. Furthermore, the DSA model assumes that the thermal particles have a Maxwell–Boltzmann distribution with a superthermal tail. The actual shape of the quasi-thermal distribution, and the shape at the point where the superthermal tail joins it, is approximated since the semianalytic calculation only self-consistently describes particles with speeds greater than the shock speed, i.e., \( v_p > v_{sk} \). The differences at low energies between what is assumed in the DSA model and the actual quasi-thermal distribution are expected to be small, but these differences may become more important if the contribution to ionization from superthermal particles is considered. Despite the approximations of the semianalytic calculation at quasi-thermal energies, it is the state of the art since the actual quasi-thermal distribution can only be determined with plasma simulations and these are not yet available for SNR parameters.

The ionization structure of shock-heated gas at a particular distance behind the shock in an SNR is determined by the electron density \( n_e \), the electron temperature \( T_e \), and the ionization and recombination rates for each ion of interest. The structure is determined by solving the collisional ionization equations in a Lagrangian gas element behind the shock:

\[
\frac{1}{n_e} \frac{Df(X_i)}{Dt} = C(X_i^{-1}, T_e) f(X_i^{-1}) + \alpha(X_i, T_e) f(X_i^{+1}) - [C(X_i, T_e) + \alpha(X_i^{-1}, T_e)] f(X_i).
\]

Here, \( f(X_i) \) is the fraction of element \( X \) in ion stage \( X_i \), and \( C(X_i, T_e) \) and \( \alpha(X_i, T_e) \) are the ionization and recombination rates out of and into ion \( X_i \), respectively.

We calculate the electron temperature by assuming that the electrons are heated by Coulomb collisions with protons and helium (Spitzer 1965). We adopt this simple prescription, which gives a lower limit to the equilibration time, knowing that the heating of electrons may, in fact, be far more complicated. For instance, there is reason to believe that collisionless wave–particle interactions with the magnetic turbulence will be important (e.g., Laming 2001), and recent work interpreting hydrogen line widths suggests that the electron-to-proton temperature ratio behind some SNR blast waves depend mainly on the shock speed, a result implying a heating process substantially different from Coulomb collisions (e.g., Ghavamian et al. 2007; Rakowski et al. 2008). However, there remain large uncertainties in connecting the measured line widths to the electron-to-proton temperature ratio (see Heng & Sunyaev 2008), and until particle-in-cell (PIC) simulations are able to model nonrelativistic, electron–proton shocks with parameters typical of SNRs, the plasma physics of electron heating will remain uncertain (see Vladimirov et al. 2008, for a discussion of the limitations of PIC simulations in this regard). In order to model some of the complexity of electron heating, we scale the Coulomb equilibration time with a parameter, \( f_{\text{eq}} \) defined in Equation (3).

At the start of the simulation, we assume that the unshocked electrons and ions are in equilibrium at a temperature \( T_0 = 10^4 \) K. We also assume that unshocked H and He are both 10% singly ionized and all heavier elements are initially neutral. While we note that this is not the precise equilibrium ionization state for \( 10^4 \) K, we emphasize that none of our results depend in any significant way on the ionization state of the unshocked material as long as it is not fully neutral. In all of the results shown here we fix the helium number density at 10% of the proton number density, \( n_{p,0} \).

At each time step, we track the ionic state \( X_i \) within each spherically symmetric fluid element by solving the time-dependent ionization equations for each abundant element (H, He, C, N, O, Ne, Mg, Si, S, Ar, Ca, Fe, and Ni). We solve the coupled set of equations with atomic data extracted from Raymond & Smith (1977), as first presented in Gaetz et al. (1988) and updated by R. J. Edgar (2008, private communication).

In Figure 1, we show an example of the time evolution of the ionization fraction, \( f(X_i) \), of high-ionization states of oxygen (O\( ^{6+} \), O\( ^{7+} \), and O\( ^{8+} \)) in a mass shell that is crossed by the FS 100 yr after the explosion. For this example, as in all we show in this paper, we have fixed parameters typical of Type Ia supernovae, i.e., the kinetic energy in ejecta from the supernova explosion \( E_{\text{SN}} = 10^{51} \) erg, the mass of the ejecta \( M_{\text{ej}} = 1.4 \ M_\odot \), the density of the ejecta follows an exponential density profile as is generally assumed for Type Ia supernovae (Dwarkadas 2000), and we assume that the supernova explodes in a CSM which is uniform with proton number density \( n_{p,0} \) and magnetic field strength \( B_0 \). In all of the models shown here, we take...
In the following examples we investigate the effect the acceleration efficiency, $\epsilon_{\text{DSA}}$, and the CSM proton density, $n_{p,0}$, has on the NEI state of some selected elements.

3.1. Ionization Versus Position

In Figure 2, we plot the ionization fractions of $\text{O}^{6+}$ and $\text{O}^{7+}$ and $\text{Si}^{12+}$ and $\text{Si}^{13+}$ in the top two panels as a function of position behind the FS. In all panels, TP results ($\epsilon_{\text{DSA}} = 1\%$) are shown with dashed curves and efficient DSA results ($\epsilon_{\text{DSA}} = 75\%$) are shown with solid curves. The electron density and electron and ion temperatures are shown in the bottom two panels. As the top two panels clearly show, higher ionization fractions are attained closer to the shock front in the efficient DSA cases, as compared with the TP cases. For instance, in the efficient case, the fraction of $\text{O}^{7+}$ peaks at a distance $R/R_{\text{FS}} \approx 0.98$ behind the shock, while in the TP case, this fraction peaks at $R/R_{\text{FS}} \approx 0.97$. We attribute the increased ionization fractions closer to the shock as a direct result of higher postshock densities in the efficient DSA case. Note that the curves extend from the FS back to the CD, indicating that the region between the FS and the CD is considerably narrower in the efficient acceleration case. This effect produces important morphological consequences (e.g., Decourchelle et al. 2000; Warren et al. 2005; Cassam-Chenaï et al. 2008).

In Figure 3, we show the same quantities as in Figure 2, except that $n_{p,0} = 0.1$ cm$^{-3}$. The lower CSM density results in lower shocked densities and in less rapid collisional ionization behind the FS. For the ions we show that higher ionization states (i.e., $\text{O}^{7+}$ and $\text{Si}^{12+}$) are considerably less populated downstream from the FS when $n_{p,0}$ is small. The differences resulting from DSA are less prominent but still evident, e.g., with $n_{p,0} = 0.1$ cm$^{-3}$, $\text{O}^{6+}$ peaks behind the shock at $R/R_{\text{FS}} \approx 0.98$ for the efficient case, and at $\approx 0.95$ in the TP case.

1. We assume that only electrons from the thermal population contribute to the NEI. In nonlinear DSA, the energetic population emerges smoothly from the thermal population (a nice example from a relativistic PIC simulation is given in Spitkovsky 2008) and superthermal particles may contribute to ionization (see Porquet et al. 2001 for a TP calculation involving a Maxwell–Boltzmann distribution with nonthermal tail). As we discussed above, superthermal particles are expected to contribute to the ionization at some level. However, the significance of this nonthermal ionization, in shocks undergoing efficient particle acceleration, has not yet been determined and remains an area of active work. For the purposes of this paper, we assume any nonthermal contribution is small.

2. We only model the interaction region between the FS and the CD where we assume cosmic elemental abundances. One reason for emphasizing the FS is that it is not certain that significant CR production occurs at the reverse shock in SNRs (e.g., Ellison et al. 2005).

3. We only consider young SNRs and do not include the effects of radiative cooling. In the high-density limit, radiative losses could be significant and the cooling timescale could be comparable to other dynamical timescales. We will investigate this effect in a subsequent paper.

3. RESULTS

In the following examples we investigate the effect of the acceleration efficiency, $\epsilon_{\text{DSA}}$, and the CSM proton density, $n_{p,0}$, has on the NEI state of some selected elements.

4. This value for $B_0$ is somewhat higher than the typically assumed 3 $\mu$G and reflects the possibility that magnetic field amplification (MFA) may be taking place. We emphasize, however, that we do not include MFA in the DSA calculation performed here. A large upstream magnetic field, $B_{\infty}$, will reduce the effects of efficient DSA, as described in Berezhko & Ellison (1999).

5. In all results shown, we assume that shocked protons and other ions have the same temperature.
Figure 2. Spatial profiles of H- and He-like oxygen and silicon, electron density, and temperature as a function of distance behind the FS. In the bottom panel, the curves labeled $T_e$ are ion (or proton) temperatures and those $T_i$ are electron temperatures. Here, and in Figures 3–5 that follow, we show values from spherically symmetric shells as a function of $R$ or $\Delta R$, not line-of-sight projections. In all panels, solid curves correspond to models with 75% efficiency, while the dashed lines correspond to TP models. These models are for a CSM proton density of $n_{p,0} = 1 \text{ cm}^{-3}$ and are calculated at $t_{\text{SNR}} = 1000 \text{ yr}$. In the model with 75% efficiency, the FS velocity is $\approx 1800 \text{ km s}^{-1}$, while in the TP model, it is $\approx 2200 \text{ km s}^{-1}$ at $t_{\text{SNR}} = 1000 \text{ yr}$. (A color version of this figure is available in the online journal.)

To emphasize the importance of the different spatial structures of ionization with $\epsilon_{\text{DSA}}$ and $n_{p,0}$, we show, in Figure 4, a close-up view of the shock fronts in Figures 2 and 3. Here, we have plotted the ionization fractions as functions of angular distance behind the shock, assuming a distance of 1 kpc. In the high-density case ($n_{p,0} = 1 \text{ cm}^{-3}$; top panel), the fraction of $\text{O}^6^+$ peaks right behind the shock at $\approx 2^\circ$ downstream, while it peaks $\approx 5^\circ$ behind the shock in the TP case. In the lower density case ($n_{p,0} = 0.1 \text{ cm}^{-3}$; lower panel), $\text{O}^6^+$ peaks $\approx 30^\circ$ behind the shock in the efficient case, but peaks well beyond $50^\circ$ behind the shock in the TP case. Similar results are found for silicon. While these models are not scaled to match any particular Galactic SNR, we believe that the angular separations shown here would easily be resolvable in current and future space-based X-ray observatories even when line-of-sight effects are taken into account. Thus, measuring the relative fraction of H-like, He-like, and even Li-like charge states would provide a useful diagnostic in studies of Galactic SNRs undergoing efficient shock acceleration.

Another interesting feature seen in Figures 2 and 3 is that the electron temperature is almost independent of $\epsilon_{\text{DSA}}$ and only varies by a factor of $\sim 2$ between the $n_{p,0} = 1 \text{ cm}^{-3}$ and $n_{p,0} = 0.1 \text{ cm}^{-3}$ cases. This is in contrast to the ion temperatures, where generally lower ion temperatures occur in the higher density models, due to the lower shock Mach number, and where the large $\epsilon_{\text{DSA}}$ cases have considerably lower ion temperatures than the TP cases. The fact that lower postshock temperatures occur in efficient DSA is well known (e.g., Ellison 2000). The electron temperature is influenced by this and by the higher densities that occur with efficient DSA. The higher postshock densities imply more collisions between electrons...
and ions, and thus more rapid temperature equilibration. The higher electron temperature combined with the higher postshock density leads to more rapid ionization, and thus higher charge states closer to the FS.

3.2. Ionization Versus Equilibration Timescale

As is clear from Figures 1–3, the ionization fraction for high charge state ions can increase with acceleration efficiency. Since the electron temperature is almost independent of \( n_{p,0} \) in these cases, we attribute this effect mainly to the higher postshock densities. However, we have assumed a particular model for temperature equilibration between protons and electrons, namely, that electrons start off cold and equilibration with the hot protons occurs only through Coulomb collisions where the equilibration timescale is given by (Spitzer 1965, Equations (5)–(31))

\[
\tau_{\text{eq}} = \frac{3 m_p m_e k_B^{3/2}}{8 (2\pi)^{1/2} n_p Z^2 Z^2 e^4} \ln \Lambda \left( \frac{T_p}{m_p} + \frac{T_e}{m_e} \right)^{3/2}.
\]

Here, \( m_p \) is the proton mass and \( T_p \) is the shocked proton temperature and definitions of the other terms are given in Spitzer (1965). It is important to note that Equation (2) places strict limits on how low the electron-to-proton temperature ratio can be behind the shock (see Hughes et al. 2000); if other equilibration mechanisms are important, such as plasma wave interactions, equilibration will occur more rapidly. To investigate the effects of more rapid temperature equilibration, we define a parameter, \( 0 \leq f_{\text{eq}} \leq 1 \), and use the equilibration time \( \tau'_{\text{eq}} \) in our calculations where

\[
\tau'_{\text{eq}} = f_{\text{eq}} \tau_{\text{eq}}.
\]

In the results shown in Figures 1–4, we have assumed \( f_{\text{eq}} = 1 \).

In Figure 5, we compare the ionization fraction of \( O^{7+} \) for \( \epsilon_{\text{DSA}} = 1\% \) and \( \epsilon_{\text{DSA}} = 75\% \) calculated with \( f_{\text{eq}} = 1 \) (black curves in all panels) and \( f_{\text{eq}} = 0.1 \) (red curves in all panels). For both values of \( \epsilon_{\text{DSA}} \), \( f(O^{7+}) \) is larger immediately behind the shock for rapid equilibration (\( f_{\text{eq}} = 0.1 \)) but drops below the \( f_{\text{eq}} = 1 \) value further downstream as \( O^{8+} \) becomes populated. The temperature plots in the bottom two panels show that the electrons and protons have come into equilibrium for a range of radii (i.e., \( 0.86 \lesssim R/R_{\text{FS}} \lesssim 0.98 \)) when \( \epsilon_{\text{DSA}} = 75\% \).
and $f_{\text{eq}} = 0.1$, but remain far from equilibrium for $f_{\text{eq}} = 1$ regardless of $\epsilon_{\text{DSA}}$. The equilibration rate changes the ionization structure for this particular ion, producing changes comparable in scale to those produced by efficient DSA.

To quantify these effects further, we look at a point midway between the CD and the FS, i.e., at $R/R_{\text{FS}} \approx 0.89$ for $\epsilon_{\text{DSA}} = 75\%$ and at $R/R_{\text{FS}} \approx 0.83$ for $\epsilon_{\text{DSA}} = 1\%$ in Figure 5. At these locations, the electron-to-proton temperature ratios are $(T_e/T_p)_{\text{TP}} \approx 0.11$ and $(T_e/T_p)_{\text{NL}} \approx 0.36$, for $f_{\text{eq}} = 1$, and $(T_e/T_p)_{\text{TP}} \approx 0.3$ and $(T_e/T_p)_{\text{NL}} \approx 1$ for $f_{\text{eq}} = 0.1$, i.e., the ratios are about three times larger with rapid equilibration. At these midpoint locations, the ionization fractions of O$^{7+}$ range from $f(O^{7+}) \approx 0.05$ for $f_{\text{eq}} = 1$ and $\epsilon_{\text{DSA}} = 75\%$ to $f(O^{7+}) \approx 0.23$ for $f_{\text{eq}} = 0.1$ and $\epsilon_{\text{DSA}} = 1\%$, i.e., about a factor of 5 span.

The electron temperature ratio for $f_{\text{eq}} = 1$ is $(T_{e,NL}/T_{e,TP})_{f_{\text{eq}}=1} = 1.8 \times 10^5$ K/2.5 $\times 10^5$ K $\approx$ 0.7 and the ratio for $f_{\text{eq}} = 0.1$ is $(T_{e,NL}/T_{e,TP})_{f_{\text{eq}}=0.1} = 3 \times 10^5$ K/6 $\times 10^5$ K $\approx$ 0.5.

For the particular parameters used in this example, the electron temperature stays within a factor of $\sim 2$ for a wide spread in $\epsilon_{\text{DSA}}$ and equilibration time, while $f(\text{O}^{7+})$ varies by a factor of $\sim 5$.

### 3.3. Emission Measure Versus Acceleration Efficiency

As seen in Figure 2 or 3, the plasma density is greatest immediately behind the shock where the electron temperature is lowest. Since the rate for electron temperature equilibration immediately behind the shock where the electron temperature is lowest. Since the rate for electron temperature equilibration immediately behind the shock where the electron temperature is lowest. Since the rate for electron temperature equilibration immediately behind the shock where the electron temperature is lowest.

In Figure 6 we plot the EM for individual ions, $EM = N_X f(X^i, R)n_e n_p dV$, and in Figure 7 we plot ionic DEMs, $\text{DEM} = \sum N_X f(X^i, R)n_e n_p dV/d(\log T_e)$, where $N_X$ is the abundance of element $X$ relative to hydrogen, $f(X^i, R)$ is the ionization fraction for the ion $X^i$ at a distance $R$ behind the shock, and $dV$ is the volume of the shell where EM or DEM is determined. The EM plotted in Figure 6 is a line-of-sight projection normalized to 1 cm$^2$ surface area, and the DEM is obtained by summing over the region between the CD and the FS.

Figure 6 clearly shows that the emission for these ions peaks much closer to the FS and is considerably stronger with efficient DSA than in the TP case. Figure 7 shows that the peak emission for these two ions shifts down in temperature by about a factor of $\sim 2$ ($\sim 1$ keV) when efficient DSA occurs. These two effects are quite significant for individual ions and should be observable. Nevertheless, the emission from a full set of ions needs to be calculated and the results folded through a detectors’ response before the signature of efficient DSA can be quantitatively determined.

### 4. DISCUSSION AND CONCLUSIONS

We have presented a calculation of the NEI in a hydrodynamic simulation of SNRs undergoing efficient DSA. While we have only explored a limited range of parameters in this paper, it is clear that the production of CRs by the outer blast wave modifies the SNR evolution and structure enough to produce significant changes in the ionization of the shocked material between the FS and the CD. In particular, higher ionization states are reached at lower electron temperatures (compared with the TP case) because of the increase in postshock density due to the increased shock compression. The calculation of thermal X-ray line emission requires the additional step of coupling the resultant ionization state vectors to a plasma emissivity code, work which is in progress. Nevertheless, our results clearly show that taking DSA into account and dynamically calculating the NEI produces changes in the ionization fractions of important elements that should translate into noticeable changes in the interpretation of X-ray line emission observed from young SNRs.
Our main results are as follows.

1. Compared with the TP case, the increase in ionization that accompanied DSA in our examples suggests that efficient DSA will result in an increase in the overall thermal X-ray emission (see Figure 6). We note that an increase in thermal emission with increasing acceleration efficiency is evident in our earlier results which explored a slightly different parameter space (i.e., Figures 7 and 8 of Ellison et al. 2007). The actual increase may depend importantly on other model parameters, such as the CSM density, and it is important to explore a more expanded parameter space to determine how broadly our results are. This work is in progress. However, regardless of whether or not efficient DSA increases the integrated thermal emission over the TP case, some thermal emission is expected because ionization is not suppressed when efficient DSA occurs. As Figure 1 shows, electrons reach X-ray emitting temperatures well before they come into equilibrium with protons and nearly as rapidly with or without efficient DSA. This occurs even if only Coulomb equilibrium is assumed. This is in contrast to recent claims (e.g., Morlino et al. 2009; Drury et al. 2009) that very weak thermal X-ray emission might result from efficient shock acceleration.

2. Compared with the TP case, ionization occurs more rapidly and, therefore, closer to the FS, with efficient acceleration (see Figures 4 and 6). The differences in spatial structure should be large enough to observe and may be used as a discriminant for the level of CR modification, if a particular ion state is coupled to other known properties, such as the dynamics and ambient conditions.

3. Efficient DSA leads to more efficient Coulomb heating of electrons and faster equilibration with ions, relative to the TP case. This results because the shocked plasma temperature is lower and the shocked density is higher when efficient DSA occurs. We showed, with a simple parameterization of the thermal equilibration time, that the signature of efficient DSA on the ionization state remains apparent for equilibration more rapid than occurs with just Coulomb collisions.

4. Using the DEM, we showed that the maximum emission from a particular ion state occurs at a significantly lower electron temperature with efficient DSA. For the ions shown in Figure 7, the difference in $T_e$ for peak emission is on the order of 1 keV while the maximum DEM remains almost constant. A difference this large will have an important impact on the interpretation of thermal X-ray emission in young SNRs.

Currently, we do not treat radiative or slow shocks, but these regimes are easily explored. For instance in a radiative shock, the cooling time might be comparable to the energy loss time in a CR-modified shock. Increases in the density will enhance the cooling to the point where radiative losses might rival losses from efficient DSA (Wagner et al. 2006). We intend to explore this regime in a forthcoming paper.

While we only considered shocked CSM here, we will consider shocked ejecta in future work. In the ejecta, the electron density can be higher and the temperature may be lower but, more importantly, the abundance structure is far more complicated than for CSM and calculations of X-ray emission are intrinsically more difficult. Furthermore, simple arguments based on the expansion of the ejecta material suggest that the magnetic field may be too low to support DSA by the reverse shock. Nevertheless, there has been speculation that particles are accelerated there (e.g., Gotthelf et al. 2001; Uchiyama & Aharonian 2008; Holder & Vink 2008) and if DSA is efficient at the reverse shock, it will likely alter the ionization balance of the shocked ejecta as much as shown here for the shocked CSM.

Finally, while we have limited our examples here to SNRs expanding into a uniform medium typical of Type Ia supernovae, we emphasize that a wider parameter space should be explored, in terms of both the structure of the ambient medium (i.e., pre-SN winds) and the parameters which determine the CR acceleration efficiency. These cases will be addressed in a follow-up paper.

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