

SPATIAL STRUCTURE AND COLLISIONLESS ELECTRON HEATING IN BALMER-DOMINATED SHOCKS

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ABSTRACT

Balmer-dominated shocks in supernova remnants (SNRs) produce strong hydrogen lines with a two-component profile composed of a narrow contribution from cold upstream hydrogen atoms and a broad contribution from hydrogen atoms that have undergone charge transfer reactions with hot protons. Observations of emission lines from edge-wise shocks in SNRs can constrain the gas velocity and collisionless electron heating at the shock front. Downstream hydrogen atoms engage in charge transfer, excitation, and ionization reactions, defining an interaction region called the shock transition zone. The properties of hot hydrogen atoms produced by charge transfers (called broad neutrals) are critical for accurately calculating the structure and radiation from the shock transition zone. This paper is the third in a series describing the kinetic, fluid, and emission properties of Balmer-dominated shocks, and it is the first to properly treat the effect of broad neutral kinetics on the shock transition zone structure. We use our models to extract shock parameters from observations of Balmer-dominated SNRs. We find that the inferred shock velocities and electron temperatures are lower than those of previous calculations by $<10\%$ for $v_s < 1500 \text{ km s}^{-1}$ and by $10\%–30\%$ for $v_s > 1500 \text{ km s}^{-1}$. This effect is primarily due to the fact that excitation by proton collisions and charge transfer to excited levels favor the high-speed part of the neutral hydrogen velocity distribution. Our results have a strong dependence on the ratio of the electron to proton temperatures, $\beta \equiv T_e/T_p$, which allows us to construct a relation $\beta(v_s)$ between the temperature ratio and the shock velocity. We compare our calculations to previous results by Ghavamian and coworkers.

Subject headings: shock waves — supernova remnants

Online material: color figures

1. INTRODUCTION

Balmer-dominated shocks in supernova remnants (SNRs) encounter upstream gas containing a substantial fraction of neutral hydrogen atoms. Emission from the shock is characterized by strong Balmer and Lyman lines with two-component profiles, which consist of a narrow contribution from direct excitation of the hydrogen atoms entering the shock and a broad contribution from the excitation of “broad neutrals,” hydrogen atoms that have been produced by charge transfer reactions with protons in the downstream gas (Chevalier & Raymond 1978; Chevalier et al. 1980). Such shocks are seen in many SNRs, including parts of the Cygnus Loop (Ghavamian et al. 2001, hereafter G01), Tycho and Kepler’s remnants (Kirshner et al. 1987, hereafter KWC87; Smith et al. 1991, hereafter S91; G01; Fesen et al. 1989, hereafter F89), RCW 86 (G01; Ghavamian et al. 2007b, hereafter G07b), SN 1006 (KWC87; S91; Ghavamian et al. 2002, hereafter G02), and several remnants located in the LMC (Tuohy et al. 1982, hereafter T82; S91; Ghavamian et al. 2003, hereafter G03; Ghavamian et al. 2007a, hereafter G07a).

Profiles of broad emission lines from edgewise observations of shocks in SNRs can be used to infer shock velocities. These calculations can be compared to analyses of the shock front proper motion to compute distances to the objects (Chevalier et al. 1980; KWC87). Observations that resolve the spatial profile of the combined $H\alpha$ emission can constrain the neutral fraction and density of the upstream gas (Raymond et al. 2007).

In addition to their use in diagnosing parameters of SNRs, Balmer-dominated shocks provide an important probe into the

plasma physics of collisionless, nonrelativistic shocks. Models of the broad component width and integrated broad-to-narrow intensity ratio can be used to derive the ratio of the electron to proton temperatures as a function of the shock velocity. Such a relation can provide insight into the physical mechanisms at work in the collisionless plasma.

This is the third in a series of papers investigating the hydrodynamics, kinetics, and line emission from Balmer-dominated shocks. In Heng & McCray (2007, hereafter Paper I), we calculated velocity distribution functions for the broad neutrals and computed the ratio of the broad-to-narrow line emission as a function of the shock velocity. In Heng et al. (2007, hereafter Paper II) we calculated the density structure of the shock transition zone, where hydrogen atoms passing through the shock front undergo charge transfer reactions, emit radiation, and become fully ionized.

The approximations employed in Paper II limit the validity of the results to shocks entering the upstream gas with velocities of $<3000 \text{ km s}^{-1}$, due to our treatment of the broad neutrals as a single fluid with the same bulk velocity as the ions (the “restricted three-component model” from § 3 of Paper II). In fact, the broad neutrals in the shock transition zone are not a fluid, but have distinct anisotropic distribution functions that depend on the number of charge transfer reactions they engage in (Paper I). In the present paper, we treat charged species in the shocked gas as fluids and describe the hydrogen atoms with appropriate kinetic distribution functions. In addition, we provide an improved calculation of the broad-line velocity profile for $v_s \gtrsim 2000 \text{ km s}^{-1}$. Using this methodology, we calculate the structure of the shock transition zone more accurately than in Paper II, and we characterize hydrogen line emission from shocks with $300 < v_s < 10,000 \text{ km s}^{-1}$.

We find that we can self-consistently determine the shock parameters for most Balmer-dominated SNRs. Our results yield lower values of the inferred shock velocity and the proton-electron

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temperature equilibration than those derived using previous models. We compute the dependence of the electron temperature on the shock velocity and compare it to the results of G07b. We note that our calculations are unable to fit the observations for several SNRs; in these cases, our basic model must be augmented by new physics to account for the data.

In § 2 of this paper, we describe our physical model. In § 3, we display the equations employed to calculate the structure of the shock transition zone and describe our numerical solution. In § 4, we compute the spatial emission profile and hydrogen line spectra from the shock transition zone. In § 5, we analyze our results and describe how observations of Balmer-dominated SNRs should be interpreted in light of our new calculations. Finally, in § 6, we discuss the implications of our results for collisionless electron heating, explore the limitations of our model, and identify areas for future research.

2. PHYSICAL MODEL

We consider a shock with a velocity of $300 < v_s < 10,000 \text{ km s}^{-1}$ traveling through the ISM, which consists of cold, partially neutral hydrogen and neutral helium. The preshock fraction of helium, relative to hydrogen, is denoted f_{He} . If the total upstream density of protons and hydrogen atoms is n_0 , then the fraction of preshock protons is defined to be $f_p \equiv n_p/n_0$.

In the frame of the shock, the cold upstream neutrals and charged particles flow downstream with a uniform velocity of v_s . At the shock, we assume that the protons are heated in a thermal distribution to $T_p \sim 3m_p v_s^2/16k_B \sim 10^7 \text{ K}$ for $v_s \sim 10^8 \text{ cm s}^{-1}$, where m_p is the proton mass, and k_B is Boltzmann's constant. At the high postshock temperatures in Balmer-dominated remnants, the emission from hydrogen excitation is much stronger than that from the forbidden transitions in metals seen in radiative shocks (e.g., Shull & McKee 1979). The length scale for the thermalization of protons is set by the proton cyclotron gyroradius, $l_{\text{gyro}} \sim 10^8 \text{ cm}$ at magnetic field strength $B \sim 10^{-4} \text{ G}$, which is much less than the length scale for the shock transition zone, $l_{\text{zone}} \gtrsim 1/(n_0\sigma) = 10^{14} \text{ cm}$ (set by the cross sections σ for the excitation, charge transfer, and ionization reactions, as well as the total upstream density).

Without any energy transfer between the shocked protons and electrons, the electron temperature is expected to be a factor of m_e/m_p less than that of the protons. Since the collisional timescale for the temperature equilibration between the electrons and protons is in some cases longer than the age of the remnant (e.g., Spitzer 1962), we assume that the electrons remain at a fixed temperature with respect to the protons in the downstream material (note, however, that intraspecies and inter-ion collisional timescales can be much shorter).⁴ However, several authors have argued that electrostatic instabilities at the shock front can increase the electron temperature, with estimates ranging from $T_e \sim 0.1T_p$ to $T_e = T_p$ (see Cargill & Papadopoulos 1988; G07b and references therein). Since the physics of nonrelativistic, collisionless shocks is still poorly understood, we parameterize the ratio of the electron to proton temperatures using the definition

$$\beta \equiv \frac{T_e}{T_p}. \quad (1)$$

One of the goals of our work is to use observations of Balmer-dominated remnants to constrain the value of β in the shock transition zone as a function of the shock velocity.

⁴ Some collisional heating of electrons by protons will occur. We estimate that such an effect can lower the derived value of β (see eq. [1]) by 1%–2%.

In the uniform-velocity upstream gas, we assume that negligible interactions take place between the neutral and charged species. The cold neutral hydrogen atoms passing through the shock are not affected by the discontinuity. In contrast, the bulk velocities of the downstream thermal protons and electrons become approximately $v_s/4$, at which point the ionization and charge transfer reactions between cold atoms and hot protons occur (see § 5.1). However, if a significant amount of the energy dissipated in the shock is used to produce cosmic rays, the ion compression ratio increases, and the ion bulk velocity decreases. Furthermore, the presence of a cosmic-ray precursor can broaden the distribution of the upstream neutrals. We discuss the implications of these physical effects in § 6.

Downstream charge transfer reactions between cold neutrals and hot protons produce a new population of hot atoms, referred to as broad neutrals. Initially, such an interaction will produce a cold proton, which is reenergized through gyromotion around the magnetic field and intraspecies Coulomb collisions with the hot proton population. We assume that such protons rapidly equilibrate back into the thermal pool. However, recent calculations by Raymond et al. (2008) indicate that interactions within the shock front may produce a distinct “pickup” ion population analogous to that in the solar wind.

H α emission is produced by excitation reactions and charge transfers to excited states in the shock transition zone. The narrow and broad components of the line are emitted by the cold hydrogen and the broad neutrals, respectively. The spatial and velocity distributions of the broad neutrals are critical for calculating the correct structure and emission from the transition zone. In Paper II, we assumed that the broad neutrals could be characterized as a single fluid, with a velocity and temperature equal to that of the ions. This treatment is flawed because, unlike the ionic species, the broad neutrals have negligible interactions among themselves. Therefore, each time a neutral engages in a charge exchange, it becomes part of a new distribution with a distinct velocity and temperature (see Paper I) and is not subsumed into the original population. This fact has not, until now, been properly taken into account in the literature.

The detailed kinetic properties of the particle distributions and their reaction rates were calculated in Paper I. We use these methods, with updated values for the atomic cross sections (see Appendix A), in our fluid calculation of the shock transition zone structure. In this picture, the broad neutrals form an infinite set of atomic populations with distinct reaction rates. In practice, as noted in Paper I, the velocity distribution function of atoms experiencing many charge transfers rapidly converges to that of the protons, usually after only three or more interactions. Thus, only three distinct broad components interact with the other species. The bulk velocities and temperatures of the broad components are written as

$$v_k = F_{vk} v_p, \quad (2)$$

$$T_k = F_{Tk} T_p, \quad (3)$$

where v_k and T_k are the velocity and temperature of the broad neutral population after k charge transfers. The coefficients $F_{vk}(v_s, \beta, f_p, f_{\text{He}})$ and $F_{Tk}(v_s, \beta, f_p, f_{\text{He}})$ are calculated using the methods described in Paper I and, in principle, are also functions of v_p and T_p . However, we show in § 5.1 that the proton velocity and temperature do not vary significantly enough in the shock transition zone to affect the values of these coefficients. For values of $k \geq 3$, F_{vk} and F_{Tk} are taken to be unity.

Ionization of helium in the shock transition zone produces singly ionized He⁺ and alpha particles. We assume that the helium ions

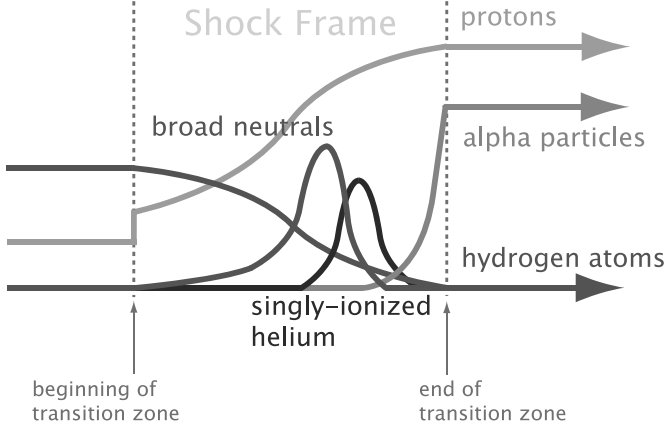


FIG. 1.— Schematic depiction of the density variation in the shock transition zone. The vertical axis represents dimensionless density for the different particle species, while the horizontal axis represents distance. The shock front is marked by the left dotted vertical line. Upstream of the shock, we assume that negligible interactions take place between the particles (although see the discussion in § 2). After the protons and electrons are isotropized and decelerated at the shock front, ionization and charge transfer reactions deplete the cold neutrals and produce hot broad atoms. The transition zone terminates when all the neutral species are ionized, leaving a mix of protons, electrons, and alpha particles. [See the electronic edition of the Journal for a color version of this figure.]

and protons are coupled through rapid inter-ion Coulomb collisions. In this case, all of the ions share the same bulk velocity (v_i) and temperature (T_i). Figure 1 schematically displays the density variation of the different particle species in the shock transition zone.

3. SPATIAL STRUCTURE OF THE TRANSITION ZONE

3.1. Basic Equations

We employ a plane-parallel coordinate system in the frame of the shock, in which the shock front is at $z = 0$. The density structure of the transition zone is determined by the conservation of mass flux:

$$\frac{d}{dz}(n_H v_s) = -n_H \sum_{s=e}^{\alpha} n_s \tilde{R}_{iH,s} - n_H n_p \tilde{R}_{T_0,p}, \quad (4)$$

$$\frac{d}{dz}(n_1 v_1) = n_H n_p \tilde{R}_{T_0,p} - n_1 \sum_{s=e}^{\alpha} n_s \tilde{R}_{iB,s} - n_1 n_p \tilde{R}_{T,p}, \quad (5)$$

$$\frac{d}{dz}(n_2 v_2) = n_1 n_p \tilde{R}_{T,p} - n_2 \sum_{s=e}^{\alpha} n_s \tilde{R}_{iB,s} - n_2 n_p \tilde{R}_{T,p}, \quad (6)$$

⋮

$$\frac{d}{dz}(n_k v_k) = n_{k-1} n_p \tilde{R}_{T,p} - n_k \sum_{s=e}^{\alpha} n_s \tilde{R}_{iB,s} - n_k n_p \tilde{R}_{T,p}, \quad (7)$$

⋮

$$\frac{d}{dz}(n_{He} v_s) = -n_{He} \sum_{s=e}^{\alpha} n_s \tilde{R}_{iHe,s}, \quad (8)$$

$$\frac{d}{dz}(n_{He^+} v_i) = n_{He} \sum_{s=e}^{\alpha} n_s \tilde{R}_{iHe,s} - n_{He^+} \sum_{s=e}^{\alpha} n_s \tilde{R}_{iHe^+,s}, \quad (9)$$

$$\frac{d}{dz}(n_{\alpha} v_i) = n_{He^+} \sum_{s=e}^{\alpha} n_s \tilde{R}_{iHe^+,s}, \quad (10)$$

$$\frac{d}{dz}(n_p v_i) = \sum_{s=e}^{\alpha} n_s \left(n_H \tilde{R}_{iH,s} + \sum_{k=1}^N n_k \tilde{R}_{iB,s} \right), \quad (11)$$

where the rate coefficients \tilde{R} for the atomic interactions (in units of $\text{cm}^3 \text{s}^{-1}$) are defined in Appendix A. Charge conservation requires the electron density to obey the relation $n_e = n_p + n_{He^+} + 2n_{\alpha}$. The subscripts iH , iHe , and iHe^+ denote the ionization of hydrogen, helium, and singly ionized helium, respectively, while the index s runs over the charged particles participating in the reactions, $s = \{e, p, He^+, \alpha\}$ (electrons, protons, singly ionized helium, and alpha particles). The subscript T_0 indicates the charge transfer of cold hydrogen atoms with protons. As discussed in § 2, the subscript k denotes the neutral hydrogen population that has engaged in k charge transfers, where the index runs from $k = 1$ to $k = \infty$. While each broad population will have a separate rate coefficient for every atomic interaction, the values are not sensitive to the details of the broad distribution functions. Therefore, we use one rate coefficient to describe the interactions of all the broad populations: the subscripts iB and T denote the ionization and the charge transfer, respectively. We list the references for our interaction cross sections and rate coefficients in Appendix A.

The system of differential equations is completed by expressions for the conservation of momentum and energy flux:

$$\frac{d}{dz} \left[m_p n_H v_s^2 + \sum_{k=1}^{\infty} (P_k + m_p n_k v_k^2) + 4m_p n_{He} v_s^2 + \sum_{s=e}^{\alpha} (P_s + m_s n_s v_i^2) \right] = 0, \quad (12)$$

$$\frac{d}{dz} \left[\frac{1}{2} m_p n_H v_s^3 + \sum_{k=1}^{\infty} \left(P_k + U_k + \frac{1}{2} m_p n_k v_k^2 \right) v_k + 2n_{He} v_s^3 + \sum_{s=e}^{\alpha} \left(P_s + U_s + \frac{1}{2} m_s n_s v_i^2 \right) v_i \right] = 0, \quad (13)$$

where $P = nk_B T$ is the pressure and $U = (3/2)nk_B T$ is the energy density. The ratios P_k/n_k and U_k/n_k are set by the kinetic calculation of the broad neutral distribution functions through equation (3), which assumes that $T_i = 3m_p v_s^2 / 16(1 + \beta)k_B$. Although the ion temperature varies by as much as 10% in the shock transition zone, this produces a negligible effect on the rate coefficients. Thus, the value of T_k is proportional to the local ion temperature. Furthermore, our steady-state fluid calculation neglects the kinetic evolution of the broad neutral distributions, which we expect to be a modest effect.

We define a natural length scale for the transition zone, according to the formula

$$l_z \equiv \frac{v_s}{n_0 \bar{R}}, \quad (14)$$

where \bar{R} is a typical value for the interaction rate coefficients, here taken to be $10^{-8} \text{cm}^3 \text{s}^{-1}$. We also define a set of dimensionless variables with $\chi \equiv z/l_z$, $\eta \equiv n/n_0$, $u \equiv v/v_s$, and $\epsilon \equiv k_B T / (m_p v_s^2)$. As discussed above, broad neutrals that have engaged in three or more charge transfers can be described with the thermal proton distribution function. We therefore define $\eta_N = \sum_{k=3}^{\infty} \eta_k$. Summing over equations (7) with $k \geq 3$ and using these definitions, the formulas for the shock transition zone structure can be written as

$$\frac{d}{d\chi} \eta_H = -\eta_H \sum_{s=e}^{\alpha} \eta_s \mathcal{R}_{iH,s} - \eta_H \eta_p \mathcal{R}_{T_0,p}, \quad (15)$$

$$\frac{d}{d\chi} (\eta_1 u_1) = \eta_H \eta_p \mathcal{R}_{T_0,p} - \eta_1 \sum_{s=e}^{\alpha} \eta_s \mathcal{R}_{iB,s} - \eta_1 \eta_p \mathcal{R}_{T,p}, \quad (16)$$

$$\frac{d}{d\chi}(\eta_2 u_2) = \eta_1 \eta_p \mathcal{R}_{T,p} - \eta_2 \sum_{s=e}^{\alpha} \eta_s \mathcal{R}_{iB,s} - \eta_2 \eta_p \mathcal{R}_{T,p}, \quad (17)$$

$$\frac{d}{d\chi}(\eta_N u_i) = \eta_2 \eta_p \mathcal{R}_{T,p} - \eta_N \sum_{s=e}^{\alpha} \eta_s \mathcal{R}_{iB,s}, \quad (18)$$

$$\frac{d}{d\chi} \eta_{\text{He}} = -\eta_{\text{He}} \sum_{s=e}^{\alpha} \eta_s \mathcal{R}_{i\text{He},s}, \quad (19)$$

$$\frac{d}{d\chi}(\eta_{\text{He}^+} u_i) = \eta_{\text{He}} \sum_{s=e}^{\alpha} \eta_s \mathcal{R}_{i\text{He},s} - \eta_{\text{He}^+} \sum_{s=e}^{\alpha} \eta_s \mathcal{R}_{i\text{He}^+,s}, \quad (20)$$

$$\frac{d}{d\chi}(\eta_{\alpha} u_i) = \eta_{\text{He}^+} \sum_{s=e}^{\alpha} \eta_s \mathcal{R}_{i\text{He}^+,s}, \quad (21)$$

$$\frac{d}{d\chi}(\eta_p u_i) = \sum_{s=e}^{\alpha} \eta_s \left(\eta_{\text{H}} \mathcal{R}_{i\text{H},s} + \sum_{k=1}^N \eta_k \mathcal{R}_{iB,s} \right), \quad (22)$$

$$\frac{d}{d\chi} \left[\eta_{\text{H}} + \sum_{k=1}^N (\eta_k \epsilon_k + \eta_k u_k^2) + 4\eta_{\text{He}} + \sum_{s=e}^{\alpha} \left(\eta_s \epsilon_s + \frac{m_s}{m_p} \eta_s u_s^2 \right) \right] = 0, \quad (23)$$

$$\frac{d}{d\chi} \left[\eta_{\text{H}} + \sum_{k=1}^N (5\eta_k \epsilon_k u_k + \eta_k u_k^3) + 4\eta_{\text{He}} + \sum_{s=e}^{\alpha} \left(5\eta_s \epsilon_s u_s + \frac{m_s}{m_p} \eta_s u_s^3 \right) \right] = 0, \quad (24)$$

where $\mathcal{R} \equiv \tilde{R}/\bar{R}$.

3.2. Solution Method

Equations (15)–(24) are a set of 10 coupled, nonlinear ordinary differential equations, which are solved with the constraint of charge conservation. We compute equations (15)–(22) directly in terms of the number density flux, denoted $y \equiv \eta u$. The ion velocity and temperature, u_i and ϵ_i , determine the broad velocity u_k and the temperature ϵ_k through the relations (2) and (3). Once y and u are known, the dimensionless number density $\eta = y/u$ can be substituted into the right-hand side of the conservation equations.

We solve equations (23) and (24) subject to the boundary conditions $y_{\text{H}}(0) = 1 - f_p$, $y_{\text{He}}(0) = (1 - f_p)f_{\text{He}}$, $y_p(0) = f_p$, and $y_k(0) = y_{\text{He}^+}(0) = y_{\alpha}(0) = 0$. The upstream neutral temperature is taken to be zero, while the protons and electrons are in thermal distributions with temperatures of $T_u \ll T_p$. Integrating these equations, substituting the initial conditions, and using the definition for y yields

$$\left(\sum_{k=1}^N F_{v,k} y_k + \sum_{s=p}^{\alpha} \frac{m_s}{m_p} y_s \right) u_i^2 + [y_{\text{H}} + 4y_{\text{He}} - 1 - (1 + \beta)f_p \epsilon_u - 4(1 - f_p)f_{\text{He}}] u_i + \left(\sum_{k=1}^N \frac{F_{T,k}}{F_{v,k}} y_k + \beta y_e + \sum_{s=p}^{\alpha} y_s \right) \epsilon_i = 0, \quad (25)$$

$$\left(\sum_{k=1}^N F_{v,k}^2 y_k + \sum_{s=p}^{\alpha} \frac{m_s}{m_p} y_s \right) u_i^2 + \left(5 \sum_{k=1}^N F_{T,k} y_k + 5\beta y_e + 5 \sum_{s=p}^{\alpha} y_s \right) \epsilon_i + y_{\text{H}} + 4y_{\text{He}} - 1 - 5(1 + \beta)f_p \epsilon_u - 4(1 - f_p)f_{\text{He}} = 0. \quad (26)$$

Note that the sums over charged species run from protons to alpha particles and do not include electrons, which appear separately in the equations (we neglect terms proportional to m_e/m_p). The system of equations (25) and (26) can be solved simultaneously to yield a quadratic equation in u_i with coefficients that are functions of χ . The solution to the quadratic has two positive roots. Only one of the roots is less than unity; it represents the physical solution for the bulk ion velocity, which must be less than the shock velocity. We use a standard Runge-Kutta method to solve equations (15)–(24), at each integration step using equations (25) and (26) to compute the ion velocity, which determines the particle densities.

Figure 2 shows the density structure of the shock transition zone for $v_s = 1000 \text{ km s}^{-1}$, $f_p = 0.5$, and $\beta = 1$. Using equation (14), the spatial coordinate has been converted from χ to the physical distance z behind the shock front. We use an external density of $n_0 = 1 \text{ cm}^{-3}$ and a helium fraction of $f_{\text{He}} = 0.1$ for all calculations in this paper. The left and right panels of Figure 2 show the dimensionless densities for the neutral and charged species, respectively. The electron and proton densities have been scaled by a factor of $1/18$. In the left panel, the solid curve shows a monotonic decrease in the density of the cold hydrogen atoms, which are removed by both the charge transfer and ionization reactions. Charge transfer produces three populations of broad neutrals. At low velocities, charge transfer dominates over ionization, and many broad neutrals with $k > 3$ are produced. The dash-dotted curve shows the density of neutral helium, which is ionized farther downstream than hydrogen due to its smaller ionization rate coefficients. In the right panel, the solid curve shows singly ionized helium, which is produced downstream from the neutral atoms and is then ionized to yield a monotonically increasing population of alpha particles, shown by the dotted curve. The proton and electron densities are depicted by the short-dashed and long-dashed curves, respectively, which saturate when all of the neutral species have been depleted. The final electron density is slightly higher than that of the protons due to the presence of the alpha particles.

4. LINE EMISSION

Once the density structure of the transition zone is determined, we can compute the hydrogen line emission, including the spatial distribution and line profiles for the broad and narrow components. We neglect collisional deexcitation and assume that every atom is excited from the ground state. $\text{H}\alpha$ photons are produced by transitions from atomic levels $3s$ and $3d$ to $2p$, as well as from $3p$ to $2s$. In the latter case, the atomic physics is complicated by a possible transition from $3p$ directly to $1s$, which results in a $\text{Ly}\beta$ photon. If the medium is optically thin to $\text{Ly}\beta$ photons (Case A conditions), this possibility can be taken into account by proper weighting of the distinct angular momentum states in constructing the excitation and charge transfer cross sections for $\text{H}\alpha$ emission:

$$\sigma_{\text{H}\alpha} = \sigma_{3s} + \sigma_{3d} + B_{3p,2s} \sigma_{3p}, \quad (27)$$

where the factor $B_{3p,2s} \approx 0.12$ is the fraction of transitions from $3p$ to $2s$. If the medium is optically thick to $\text{Ly}\beta$ emission (Case B conditions), reabsorption by ground-state hydrogen effectively traps $\text{Ly}\beta$ photons until they are reemitted as $\text{H}\alpha$ photons. In this case, all of the transitions eventually result in $\text{H}\alpha$ emission, and we set $B_{3p,2s} \approx 1$.

Case A and B conditions represent the two extremes of media that are optically thin and thick to $\text{Ly}\beta$ scattering. For the stationary atoms (i.e., cold hydrogen) that produce the narrow line, the

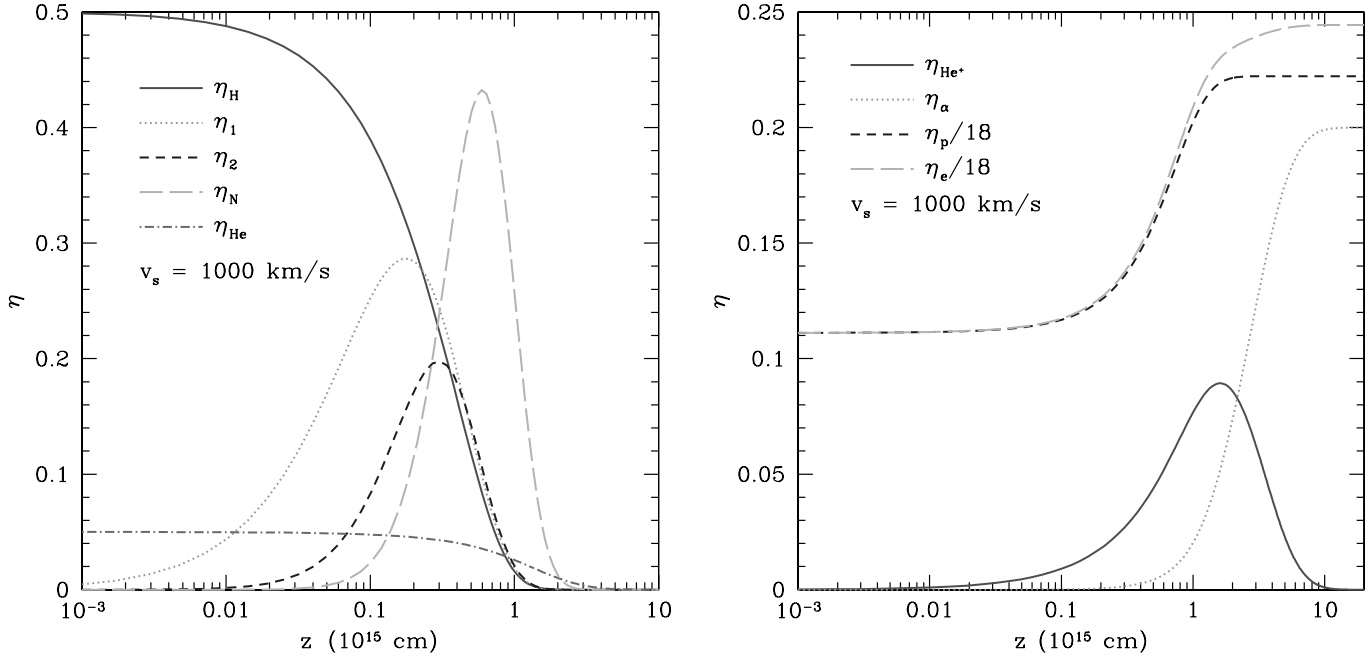


FIG. 2.— Dimensionless density as a function of the physical distance behind the shock front, for the parameters $v_s = 1000 \text{ km s}^{-1}$, $f_p = 0.5$, $\beta = 1$, $n_0 = 1 \text{ cm}^{-3}$, and $f_{\text{He}} = 0.1$. The left panel shows the shock transition zone structure for the neutral species; the solid and dash-dotted curves show the densities of neutral hydrogen and helium, respectively. Charge transfer reactions produce broad neutrals, which have undergone one (*dotted curve*), two (*short-dashed curve*), or three or more (*long-dashed curve*) charge transfer reactions. The right panel shows the transition zone structure for the charged species. Note that the densities of the protons and electrons have been scaled by a factor of 1/18. Singly ionized helium (*solid curve*) is produced downstream from the shock and is ionized to produce alpha particles (*dotted curve*). The proton (*short-dashed curve*) and electron (*long-dashed curve*) densities saturate when the neutral species are fully ionized. [See the electronic edition of the Journal for a color version of this figure.]

optical depth to the scattering of Ly β photons is, at the line center, $\tau_\beta \sim n_{\text{H}} \sigma l_{\text{zone}} \gtrsim 1$, with $n_{\text{H}} \sim 0.5 \text{ cm}^{-3}$, $\sigma \sim 10^{-14} \text{ cm}^2$, and $l_{\text{zone}} \sim 2 \times 10^{14} \text{ cm}$ (see, e.g., Rybicki & Lightman 1979; Cox 2000). The column of the upstream neutrals will also line-scatter Ly β emission, increasing the effective value of τ_β . Thus, in conditions appropriate for many Balmer-dominated SNRs, partial scattering of Lyman photons will occur, producing results intermediate between Cases A and B for the narrow line (G01).

The conversion efficiency of the narrow-component Ly β to H α was first computed by Chevalier et al. (1980). Subsequent Monte Carlo calculations were performed by Laming et al. (1996) and G01. Briefly, models of the neutral hydrogen density and the narrow-component excitation rate are computed as a function of the distance behind the shock. These are used to calculate the profile of excitations to the $3p$ level. Photons are emitted at frequencies distributed according to the preshock temperature in random directions. They are followed as they are absorbed and reemitted as Ly β or H α at different locations until they escape. The conversion fraction depends on the shock speed, the electron-ion temperature ratio, and the preshock ionization fraction, but not on the total upstream density. We assume that the variation with f_p is modest and that the temperature ratio decreases from $\beta \sim 1$ for $v_s \lesssim 500 \text{ km s}^{-1}$ to $\beta < 0.1$ for $v_s > 2000 \text{ km s}^{-1}$, and we use the calculations of Laming et al. (1996), in which 60%–75% of the $3p$ excitations result in Ly β photons. We fit the resulting conversion fraction as a function of the shock velocity, yielding

$$B_{3p,2s} = 12v_{s,5}^{-1} + (0.63 + 3.6 \times 10^{-5}v_{s,5}), \quad (28)$$

where $v_{s,5}$ is the shock velocity in units of 10^5 cm s^{-1} .

4.1. Spatial Emissivity Profile

To determine the emission, we employ the excitation rate coefficients calculated using the methods of Paper I (see also

Appendix A). An emission-line photon is produced when a cold hydrogen atom or a broad neutral is excited by a charged particle or undergoes charge transfer to an excited state. In the former case, the rates must be weighted by the probability of repeated excitation. We use the following equations to calculate the spatial emissivity profiles for the narrow (ξ_n) and broad (ξ_b) components:

$$\xi_n(z) = \frac{n_{\text{H}}}{1 - P_{E_0}} \sum_{s=e}^{\alpha} n_s \tilde{R}_{\Delta n, E_0, s}, \quad (29)$$

$$\xi_b(z) = n_{\text{H}} n_p \tilde{R}_{\Delta n, T_0^*, p} + \sum_{k=1}^N n_k \left(n_p \tilde{R}_{\Delta n, T^*, p} + \frac{1}{1 - P_E} \sum_{s=e}^{\alpha} n_s \tilde{R}_{\Delta n, E, s} \right), \quad (30)$$

where the rate coefficients are labeled by the transition $\Delta n = \text{H}\alpha$, Ly α , and Ly β ; the atomic interaction; and the particle type. The symbols E_0 and T_0^* denote excitation and charge transfer to an excited state for cold hydrogen, while E and T^* denote these interactions for broad neutrals. The probabilities $P_{E_0} = R_{E_0}/(R_{E_0} + R_{I_0} + R_{T_0})$ and $P_E = R_E/(R_E + R_I + R_T)$ are calculated using the total reaction rates per atom or broad neutral for ionization, excitation, and charge transfer. These are calculated using the weighted sum of the rate coefficients; for example, $R_I = \sum_{k=1}^N n_k \sum_{s=e}^{\alpha} n_s \tilde{R}_{iB, s}$.

Figure 3 shows the spatial emissivity profiles for the narrow and broad components as a function of the distance z behind the shock front. The thin vertical lines indicate the centroids $z_{cn, b}$ of the emission components, calculated according to the formula

$$\int_0^{z_{cn, b}} dz \xi_{n, b}(z) - \int_{z_{cn, b}}^{\infty} dz \xi_{n, b}(z) = 0. \quad (31)$$

For the low value of f_p shown in this figure, significant ionization must occur before there are enough charged particles to excite

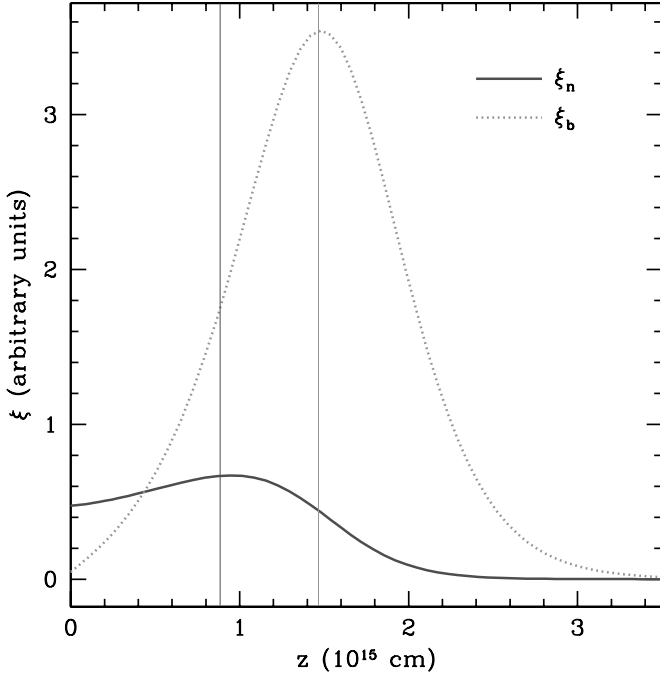


FIG. 3.—Emissivity profiles for narrow and broad emission as a function of the distance behind the shock front. Results are shown for $v_s = 1000 \text{ km s}^{-1}$, $f_p = 0.1$, $\beta = 1$, and $f_{\text{He}} = 0.1$. The centroids of the narrow and broad profiles are indicated by the thin vertical lines. [See the electronic edition of the *Journal for a color version of this figure.*]

the narrow emission and engage in charge transfers to produce broad neutrals. Therefore, the intensity peaks for both components are shifted downstream from the shock front, and the centroid of the broad-line emission is shifted farther downstream than that of the narrow line.

4.2. Line Profiles

The full width at half-maximum (FWHM) of the broad line can be related to the velocity and temperature equilibration of the shock. The line profile is a convolution of the broad neutral and exciting species distribution functions with the cross sections for excitation and charge transfer to an excited state, projected along the line of sight to the observer. In a cylindrical coordinate system (r, θ, z) , where the z -axis is along the shock velocity direction, it is straightforward to project the line profiles for observers oriented both edgewise (along the r -axis) and face-on (along the z -axis) with respect to the shock front. Due to limb brightening, most observations will be selected for shocks viewed edgewise or nearly so.

We calculate the face-on (ϕ_{FO}) and edgewise (ϕ_{EW}) hydrogen line profiles according to the formulas

$$\begin{aligned} \phi_{\text{FO}}(v_z) = & n_{\text{H}} n_p F_{T_p, z}(v_s) \\ & + \int d^3 \mathbf{v}' \left[\sum_{s=e}^{\alpha} n_s f_s(\mathbf{v}') F_{Es, z}(v_z, \mathbf{v}') \right. \\ & \left. + n_p f_p(\mathbf{v}') F_{Tp, z}(v_z, \mathbf{v}') \right], \end{aligned} \quad (32)$$

$$\begin{aligned} \phi_{\text{EW}}(v_r, z) = & n_{\text{H}} n_p F_{T_p, r}(v_s) \\ & + \int d^3 \mathbf{v}' \left[\sum_{s=e}^{\alpha} n_s f_s(\mathbf{v}') F_{Es, r}(v_r, \mathbf{v}') \right. \\ & \left. + n_p f_p(\mathbf{v}') F_{Tp, r}(v_r, \mathbf{v}') \right], \end{aligned} \quad (33)$$

$$F_{Xs, z} = \int_0^{2\pi} d\theta \int_0^{\infty} dv_r v_r \sum_{k=1}^N n_k f_k(\mathbf{v}) \Delta v \sigma_{X, \text{H}\alpha, s}(\Delta v), \quad (34)$$

$$F_{Xs, r} = v_r \int_0^{2\pi} d\theta \int_{-\infty}^{\infty} dv_z \sum_{k=1}^N n_k f_k(\mathbf{v}) \Delta v \sigma_{X, \text{H}\alpha, s}(\Delta v), \quad (35)$$

where $\Delta v \equiv |\mathbf{v} - \mathbf{v}'|$ and the cross sections $\sigma_{X, \text{H}\alpha, s}$ denote the excitation and charge transfer reactions for $\text{H}\alpha$ emission to appropriately weighted angular momentum states, using equations (27) and (28). It should be noted that in equations (32)–(35), the quantities denoted by f are kinetic distribution functions and should not be confused with the preshock ionization fraction f_p . The edgewise profile can be calculated as a function of the position behind the shock. In practice, however, the observed lines are not spatially resolved. We therefore calculate the profiles as a function of v_s and β , spatially averaged over the shock transition zone. Figure 4 shows examples of symmetric, edgewise broad neutral velocity distributions, which are inputs to equation (33). Results are depicted in a reference frame in which the average ion velocity is zero, for $\beta = 0.01, 0.1, 0.5$, and 1 , with $v_s = 1000 \text{ km s}^{-1}$ (left) and 7000 km s^{-1} (right). For all cases, the preshock ionization fraction is set to $f_p = 0.5$.

For $v_s \lesssim 1500 \text{ km s}^{-1}$, excitation by electrons dominates over that by protons (including charge transfer into excited states). Even at moderate values of the electron temperature, $\beta \lesssim 0.1$, the thermal width of the electron distribution function is much larger than that of the broad neutral distribution, and the electrons typically have much higher velocities. Thus, for the majority of the range of integration, $\Delta v \approx v'$, and the electron contribution to equations (32) and (33) is approximately equal to the projected velocity distribution of the broad neutrals, multiplied by the rate coefficient for excitation by electrons. This approximation has been used in previous studies of the line profile (e.g., Chevalier et al. 1980; G01; Paper I). Nevertheless, for shock velocities of $v_s \sim 2000 \text{ km s}^{-1}$, the excitation rates by electrons and protons are comparable, with the proton contribution increasing for faster shocks. The cross sections for excitation by protons increase with the relative speed of the colliding particles. Since the broad neutrals and the ions have comparable speeds, high-speed neutrals are more likely to produce $\text{H}\alpha$ photons. The effect on the line profile is somewhat mitigated by the integration over velocity space, but the observed line width is larger than the velocity width of the broad neutral distribution. This effect is relatively small at low velocity ($< 10\%$) but is significant for high-velocity shocks (see § 6).

Figure 5 displays the FWHM of the broad, edgewise $\text{H}\alpha$ line profile as a function of v_s , at several values of $\beta = 0.01, 0.1, 0.5$, and 1 . The preshock ionization is set at $f_p = 0.5$. The FWHM increases monotonically with v_s due to the increased temperature of the postshock ion distributions (and hence the broad neutral distributions) behind the faster moving shocks. As β is increased, energy is transferred from the protons to heat the electrons, leading to lower proton temperatures and a smaller FWHM for the broad component.

For fast shocks, in which proton excitation contributes substantially to equations (32) and (33), we expect different transitions (e.g., $\text{H}\alpha$ and $\text{Ly}\alpha$) to have different FWHM relations, since they employ distinct reaction cross sections in the calculation of ϕ . This is a unique prediction made by our calculation, and it may have consequences for studies of $\text{Ly}\alpha$ emission from shocks with $v_s \gtrsim 4000 \text{ km s}^{-1}$, in which the FWHM of

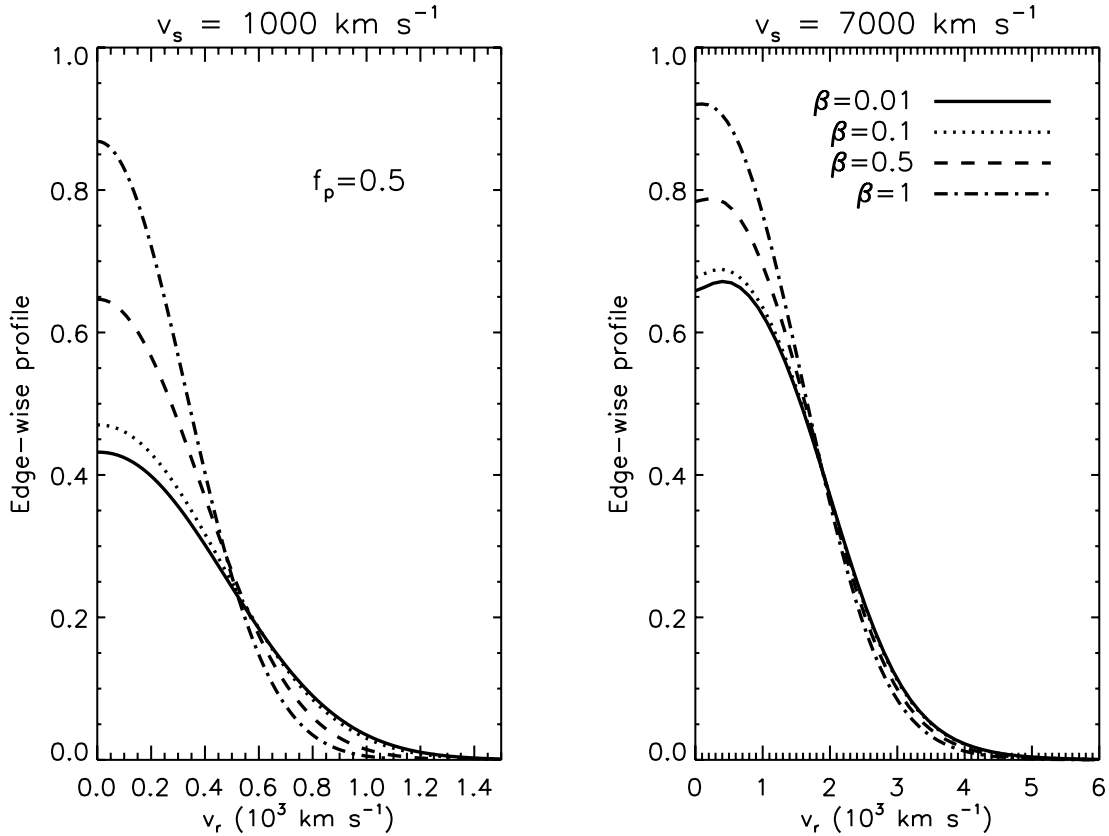


FIG. 4.— Symmetric, edgewise broad neutral velocity distributions as a function of the line-of-sight velocity v_r , in a reference frame in which the proton velocity is zero. Results are shown for a fixed value of $f_p = 0.5$ and several values of $\beta = 0.01, 0.1, 0.5$, and 1 , with $v_s = 1000 \text{ km s}^{-1}$ (left) and 7000 km s^{-1} (right). As the value of β is increased, the temperatures of the proton and broad neutral distribution functions decrease, leading to a smaller FWHM.

the Ly α line increases from 10% to 60% over that of the H α line.

5. DEPENDENCE ON SHOCK PARAMETERS

Here we describe how the structure and emission from the transition zone of Balmer-dominated SNRs depend on the shock velocities (ranging from $v_s = 300$ to $10,000 \text{ km s}^{-1}$), the pre-shock ionization fractions (ranging from $f_p = 0.1$ to 0.9), and the temperature equilibration ratios (ranging from $\beta = 0.1$ to 1), as well as the Case A and Case B conditions for the narrow line.

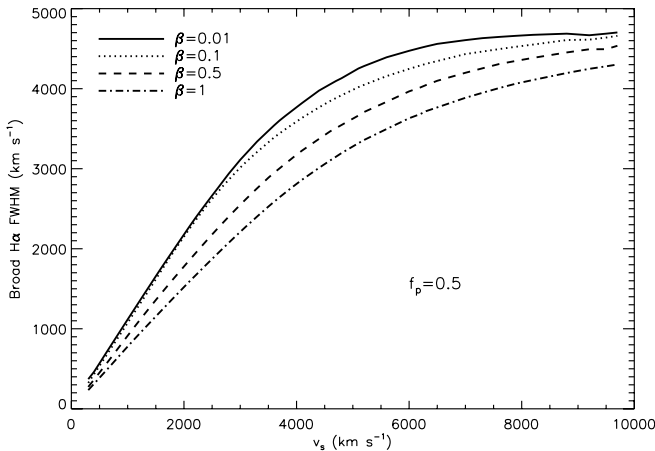


FIG. 5.— Broad, edgewise H α line FWHM as a function of the shock velocity for a fixed value of $f_p = 0.5$ at several values of $\beta = 0.01, 0.1, 0.5$, and 1 for H α transitions. As the value of β is increased, the predicted FWHM decreases (see Fig. 4).

5.1. Spatial Structure of the Shock Transition Zone

The basic features of the shock transition zone density structure were shown in Figure 2. The details of broad neutral production are highly sensitive to the values of v_s, f_p , and β . We can see the effect of increasing the shock velocity by comparing Figure 6 to Figure 2. At high values of v_s , the ionization rate is greater than the charge transfer rate, so relatively few broad neutrals are produced. The densities of the broad neutral populations decrease rapidly with each subsequent charge transfer reaction. The helium ionization rates decrease, causing the He $^+$ and alpha particle production to peak further downstream.

In Paper II, we explored variations in the ion velocity throughout the shock transition zone and found that, for $v_s \gtrsim 300 \text{ km s}^{-1}$, $v_i \approx v_s/4$ with negligible deviation. We confirm this conclusion with our multicomponent models. Figure 7 shows the dimensionless ion velocity and temperature as a function of the distance from the shock front. Results are shown for $v_s = 1000$ and 4000 km s^{-1} , with $f_p = 0.5$ and $\beta = 1$. In the left panel, we display the percent deviation of the ion velocity from $v_s/4$ for the two shock velocities. At $v_s < 1600 \text{ km s}^{-1}$, broad neutrals are produced with a smaller momentum density than that of the original ion population, leading to an increase in the ion velocity by the conservation of momentum (solid curve). For $v_s \gtrsim 1600 \text{ km s}^{-1}$, the opposite occurs, and the ion velocity decreases as the broad neutrals are produced (dotted curve). The deviation of the ion velocity from $v_s/4$ is less than 1% and has a negligible effect on calculations of the reaction rate coefficients. In the right panel, we plot the ion temperature profile. In this case, the broad neutrals are produced with slightly lower temperatures than the ions, leading to a slight heating of the ions. At large distances downstream, the presence

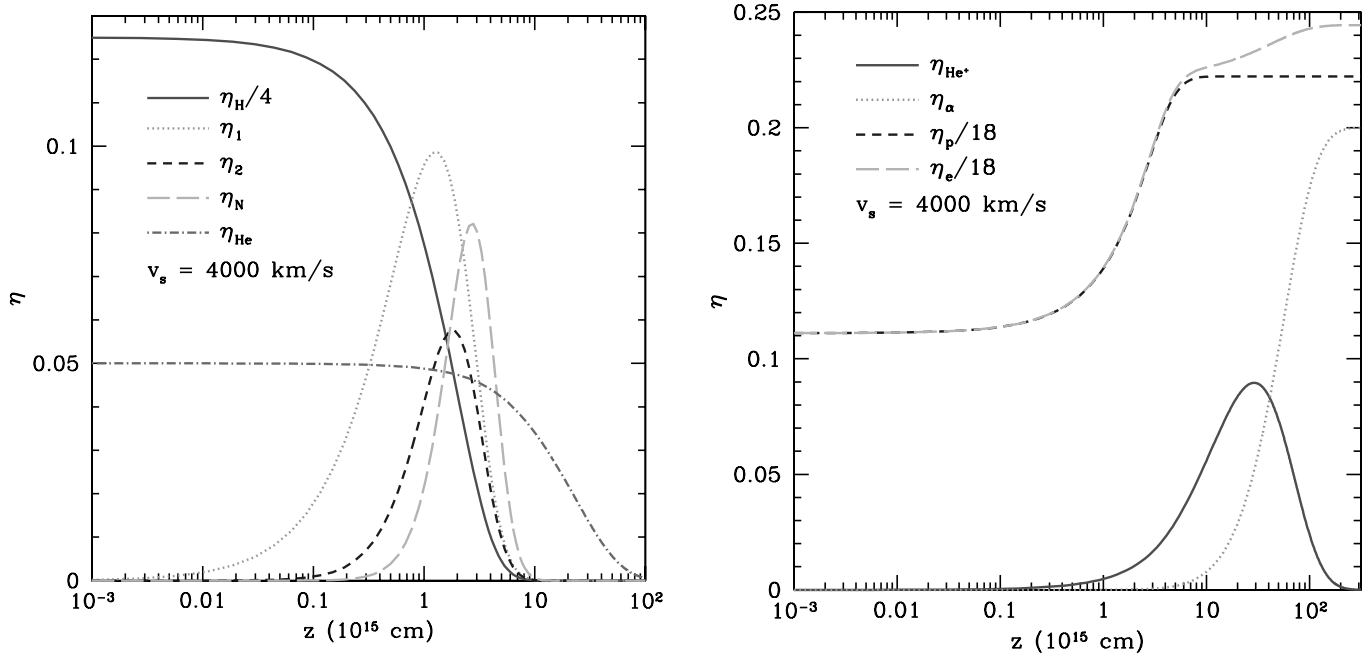


FIG. 6.—Same as Fig. 2, but for $v_s = 4000 \text{ km s}^{-1}$. The density of the cold neutral hydrogen is scaled by 1/4. [See the electronic edition of the Journal for a color version of this figure.]

of alpha particles increases the mean atomic mass in the gas to $1.27m_p$ (under the assumption of a 10% helium abundance), leading to a final ion temperature that is slightly higher than $\epsilon_i \sim 3/[16(1 + \beta)]$. The maximum deviation of ϵ_i from its expected value is of order 10%, which in practice has a negligible effect on the values of the reaction rate coefficients.

5.2. H α Emissivity

In Figure 3, we show a typical emissivity profile at a relatively low shock velocity and ionization fraction. In the left panel of Figure 8, we show the effect of increasing the initial ionization

fraction to $f_p = 0.9$. We see that the neutral density decreases rapidly and the narrow-line emission peaks at the shock front (*left*). Many charge transfer reactions occur close to the shock, pushing the centroid of the broad emission farther upstream compared to the $f_p = 0.1$ case. Comparing the right panel of Figure 8 to Figure 3, we can see the effect of increasing the shock velocity. For the high-velocity case, both the ionization and charge transfer rates are decreased, shifting the centroids of both line components farther downstream in the transition zone.

As discussed by Raymond et al. (2007) and Paper II, the spatial shift between the broad- and narrow-line centroids potentially

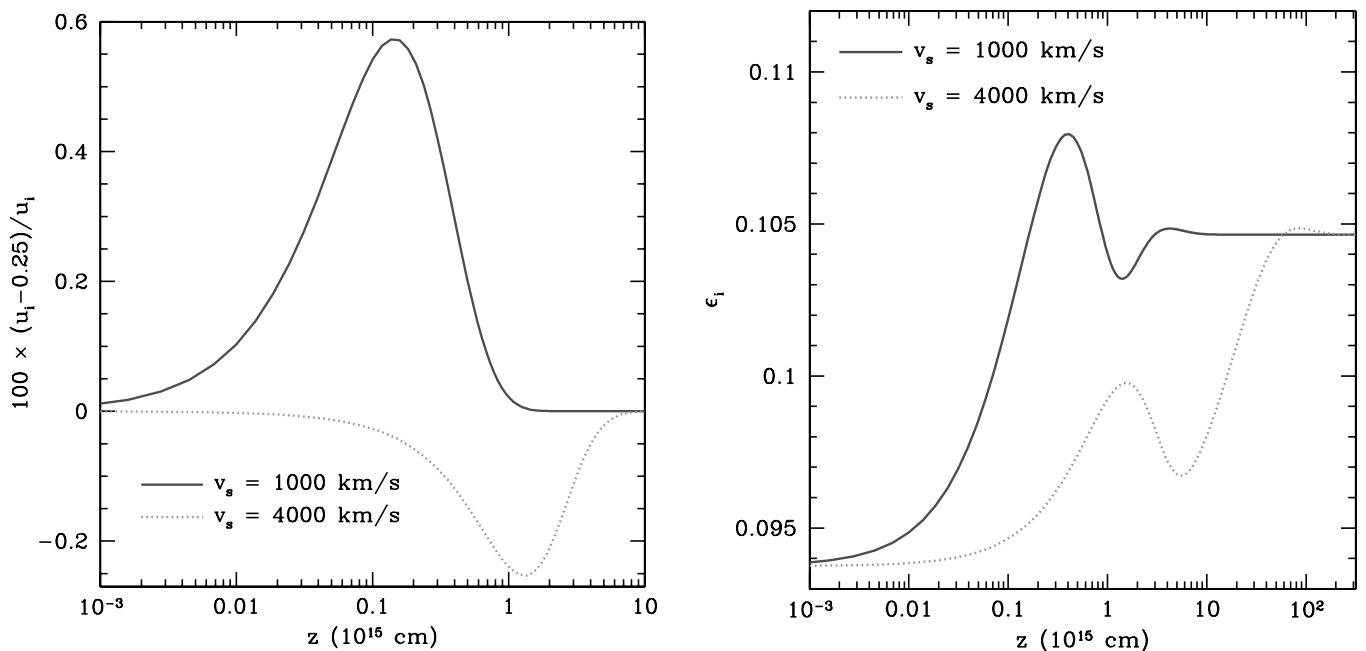


FIG. 7.—Percent deviation of the ion velocity from $v_s/4$ (*left*) and the ion temperature (*right*) as functions of the position behind the shock front for two values of the shock velocity: $v_s = 1000 \text{ km s}^{-1}$ (*solid curves*) and $v_s = 4000 \text{ km s}^{-1}$ (*dotted curves*). Results are shown for $f_p = 0.5$, $\beta = 1$, and $f_{\text{He}} = 0.1$. [See the electronic edition of the Journal for a color version of this figure.]

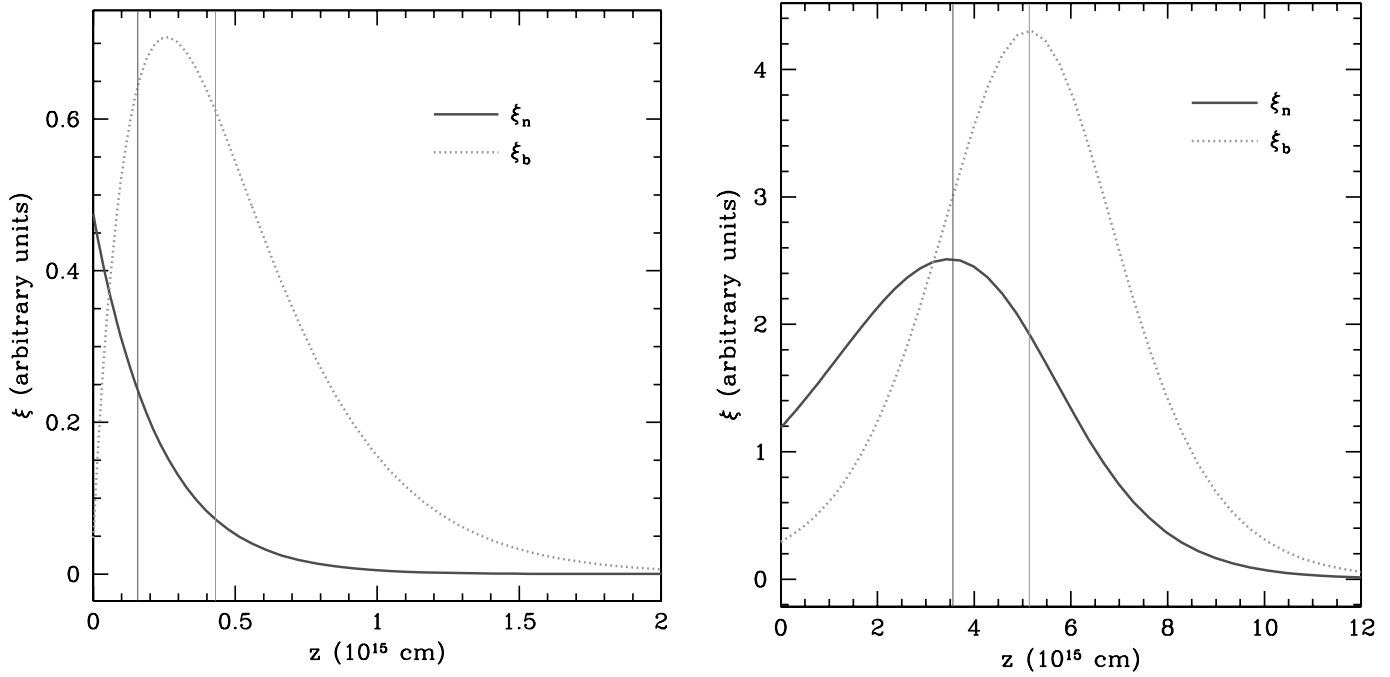


FIG. 8.— Same as Fig. 3, but for $f_p = 0.9$ and $v_s = 1000 \text{ km s}^{-1}$ (left) and $f_p = 0.1$ and $v_s = 4000 \text{ km s}^{-1}$ (right). [See the electronic edition of the Journal for a color version of this figure.]

provides a constraint on the preshock ionization fraction, f_p , and the external density, n_0 . In Figure 9, we plot the spatial shift $z_{\text{sh}} \equiv z_{cb} - z_{cn}$ as a function of the shock velocity v_s , with $n_0 = 1 \text{ cm}^{-3}$. Note that the spatial shift scales as $z \propto 1/n_0$. In the left panel, we show how z_{sh} depends on f_p , holding $\beta = 1$. As the velocity increases, the charge transfer rate decreases relative to the ionization rate, delaying the production of broad neutrals. Consequently, the centroid of ξ_b shifts downstream and the value of z_{sh} in-

creases. For small values of f_p , few protons initially exist to engage in charge exchange. Consequently, z_{cb} shifts downstream from the shock front. In the right panel, we show how the results depend on the temperature equilibration parameter, β , for a fixed value of $f_p = 0.5$. At shock velocities of $v_s \gtrsim 1000 \text{ km s}^{-1}$, increasing β reduces the proton ionization rate for cold hydrogen. This effect causes the peak of broad production to shift downstream from the shock front. At high velocities, increasing β

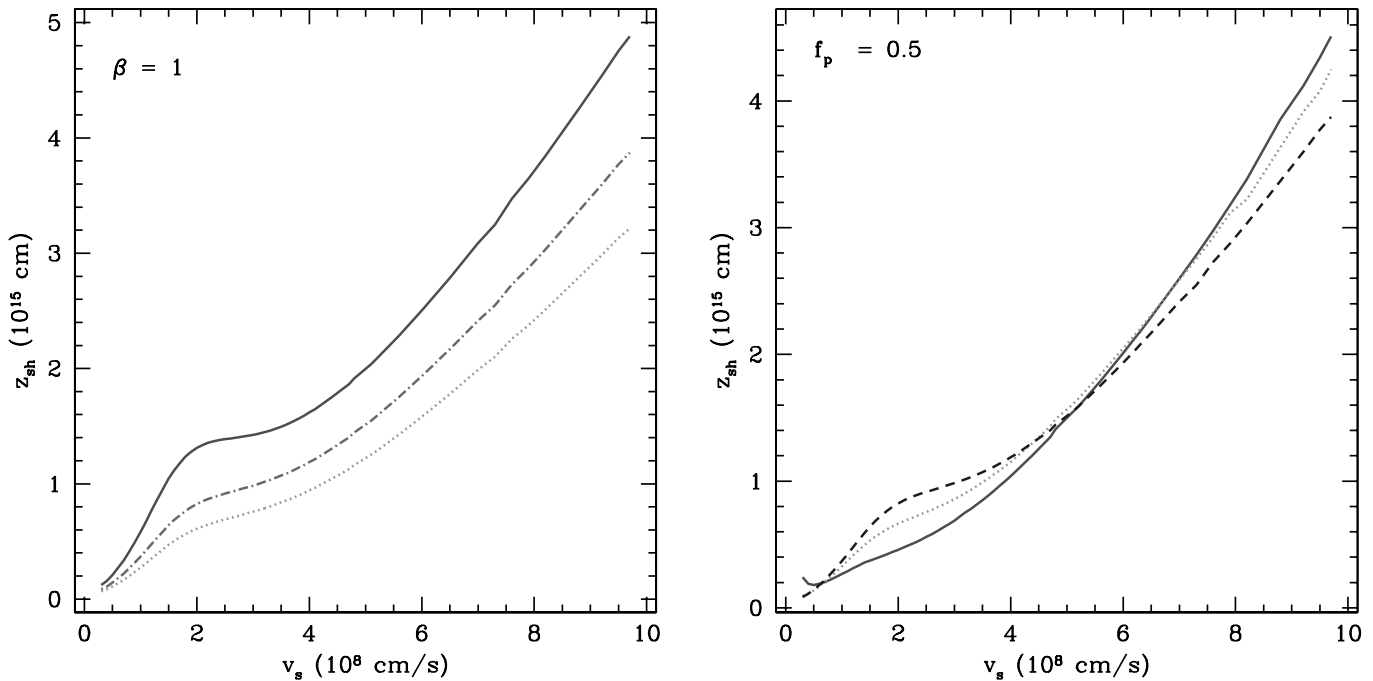


FIG. 9.— Spatial shift, in units of 10^{15} cm , as a function of the shock velocity v_s . The preshock density is set to $n_0 = 1 \text{ cm}^{-3}$; note that z_{sh} scales as $1/n_0$. The left panel shows the variation of the shift with f_p at a fixed value of $\beta = 1$. From top to bottom, the curves show the results for preshock ionization fractions of $f_p = 0.1, 0.5$, and 0.9 . The right panel shows the variation of the shift with β , for a fixed value of $f_p = 0.5$. From bottom to top at $v_s \sim 2000 \text{ km s}^{-1}$, the curves show the results for temperature equilibration ratios of $\beta = 0.1, 0.5$, and 1 . [See the electronic edition of the Journal for a color version of this figure.]

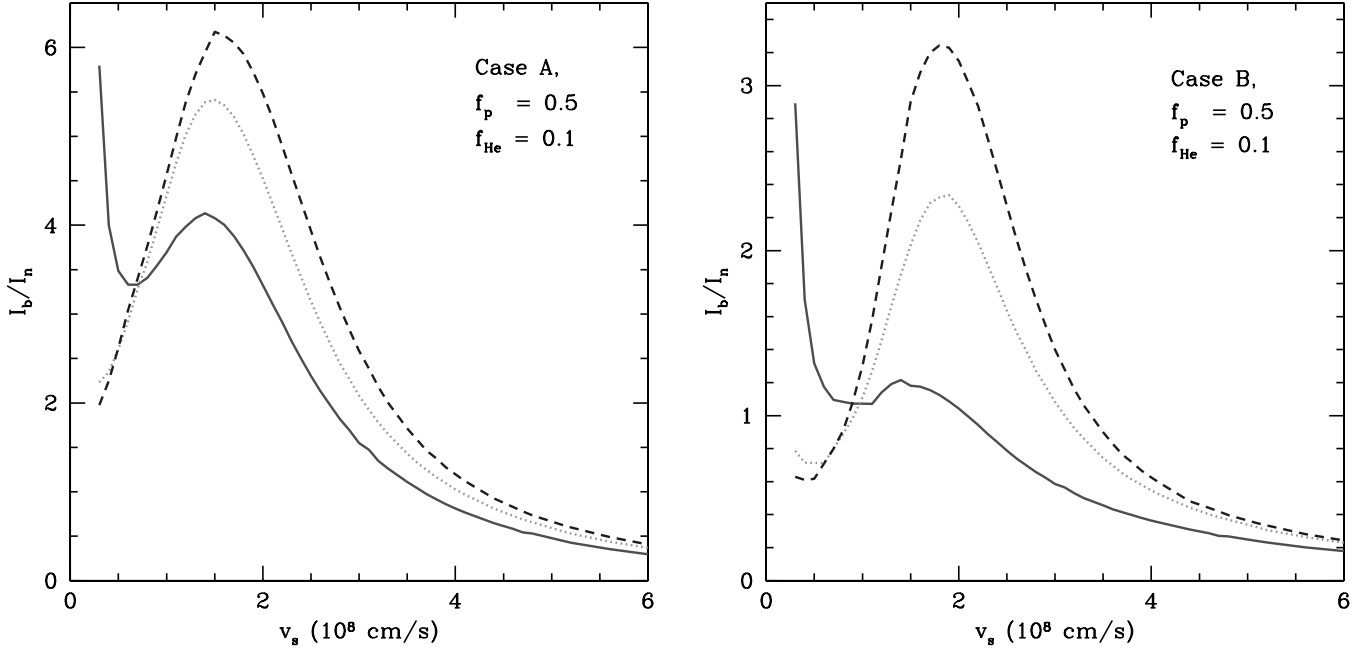


FIG. 10.—Broad-to-narrow intensity ratio I_b/I_n as a function of the shock velocity v_s , with $f_p = 0.5$ and $f_{\text{He}} = 0.1$, for Case A (left) and Case B (right) conditions. From bottom to top, the curves show the results for temperature equilibration ratios of $\beta = 0.1, 0.5$, and 1. [See the electronic edition of the *Journal* for a color version of this figure.]

tends to increase the charge transfer rate for broad neutrals over that for cold hydrogen. This effect causes the peak of broad production to shift closer to the shock front and decreases the value of z_{sh} . For more discussion on the observational significance of the spatial emissivity profile and the shift, see § 6.2.

Another important observational diagnostic for Balmer-dominated SNRs is the ratio of the integrated broad-to-narrow line strengths, defined as

$$I_b/I_n = \frac{\int_0^\infty dz \xi_b(z)}{\int_0^\infty dz \xi_n(z)}. \quad (36)$$

This ratio has a strong dependence on both v_s and β . Figure 10 shows I_b/I_n versus the shock velocity at a fixed value of $f_p = 0.5$, for several values of the temperature equilibration $\beta = 0.1, 0.5$, and 1, using Case A (left) and Case B (right) conditions for the narrow line. For a given β , the variation of the intensity ratio with velocity is the result of competition between charge transfers and ionizations, which contribute equally at $v_s \sim 2000 \text{ km s}^{-1}$. At very low values of v_s , the ionization rate begins to decrease precipitously while the charge transfer rate stays roughly constant, leading to the spike in I_b/I_n . At a fixed shock velocity, increasing β causes a decrease in the proton ionization rate for broad neutrals compared to cold hydrogen, leading to an increase in the intensity ratio. In Case B conditions (right), additional narrow emission due to the absorption of trapped Ly β photons decreases the values of I_b/I_n relative to those of Case A. As noted in § 4, for most Balmer-dominated SNRs, partial line-scattering of Ly β photons yields emission intermediate between Cases A and B.

According to the picture in Paper I, the dependence of I_b/I_n on f_p is expected to be weak, since the number of broad neutrals produced in each population per cold hydrogen atom is fixed. This intuition is confirmed by our multicomponent calculation. When f_p is increased, we expect to see more charge transfers in the shock transition zone. However, this effect is balanced by higher ionization rates for both the broad and cold neutral populations. The net result is a negligible change in the I_b/I_n ratio.

The presence of neutral helium introduces a weak dependence of I_b/I_n on f_p . In Figure 11, we show I_b/I_n as a function of v_s for a fixed value of $\beta = 1$ in Case A (left) and Case B (right) conditions. As f_p decreases, fewer protons engage in charge transfer and ionization reactions close to the shock front, shifting the peak of the broad neutral production downstream. In this case, the broad neutrals persist far enough downstream to interact with the charged helium species produced there, altering the ratio I_b/I_n . This effect is of order $\lesssim 1\%$ for $f_{\text{He}} = 0.1$ and remains small for larger helium fractions. However, it should be noted that when the effects of partial Ly β scattering are included, variations in optical depth with f_p can introduce a dependence of I_b/I_n on the preshock ionization fraction (G01). While we have roughly treated Ly β scattering using the prescription of § 4, a full calculation of line-scattering in the shock transition zone is needed to properly account for these effects and incorporate the dependence of the integrated line ratio on the preshock ionization fraction. This is a source of systematic error in our calculation.

Neglecting the f_p dependence, we can write the broad-to-narrow intensity ratio as

$$I_b/I_n = L_{bn}(v_s, \beta), \quad (37)$$

where I_b/I_n and L_{bn} represent the measured and theoretical values of the intensity ratio, respectively.

5.3. Broad-Line Profiles and Observational Interpretations

Equations (32) and (33) allow us to model the FWHM of the broad line as a function of v_s and β , as shown in Figure 5. In the extreme scenarios of the Case A and Case B conditions, the FWHM has a weak ($\lesssim 1\%$) dependence on the preshock ionization after a spatial average is taken over the shock transition zone. Figure 12 shows examples of broad neutral velocity distributions (which are inputs to eq. [32]) for the case of the face-on orientation of the shock as a function of the line-of-sight velocity v_z , in a frame of reference in which the ion velocity is zero. Results are shown for a fixed value of $f_p = 0.5$ at several values of

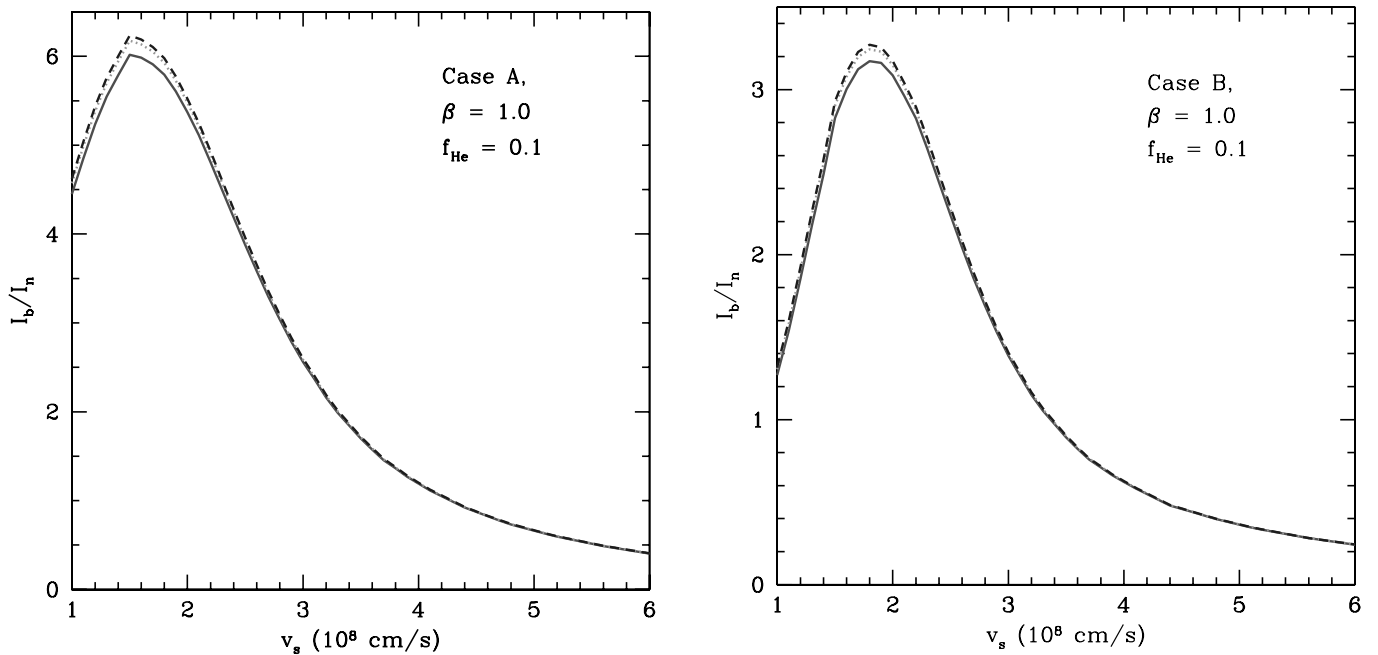


FIG. 11.—Broad-to-narrow intensity ratio I_b/I_n as a function of the shock velocity v_s , with $\beta = 1$ and $f_{\text{He}} = 0.1$, for Case A (left) and Case B (right) conditions. From bottom to top, the curves show results for pre-shock ionization fractions of $f_p = 0.1, 0.5$, and 0.9 . [See the electronic edition of the Journal for a color version of this figure.]

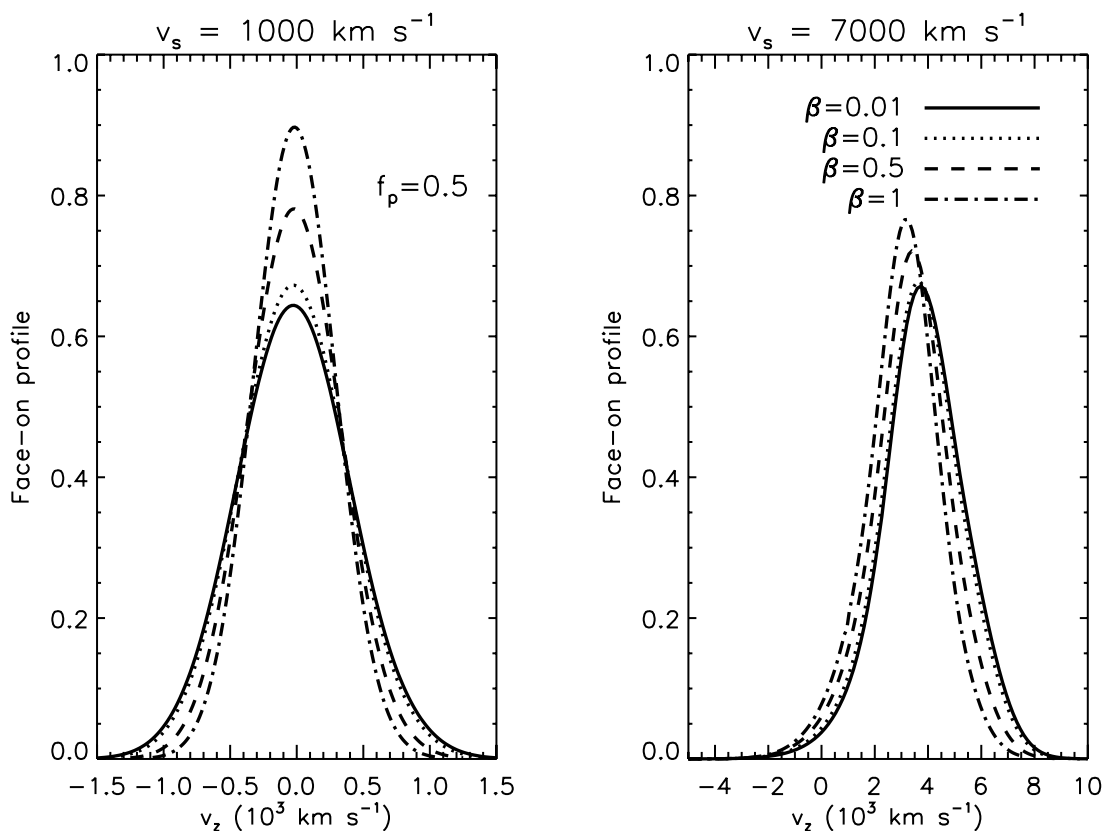


FIG. 12.—Face-on broad neutral velocity distributions as a function of the line-of-sight velocity v_z , in a reference frame in which the proton velocity is zero. Results are shown for a fixed value of $f_p = 0.5$ and several values of $\beta = 0.01, 0.1, 0.5$, and 1 , at shock velocities of $v_s = 1000 \text{ km s}^{-1}$ (left) and $v_s = 7000 \text{ km s}^{-1}$ (right). At low shock velocities, charge transfer is extremely efficient, and the broad neutral distributions are very close to the thermal proton distribution, with the broad neutrals moving slightly slower than the ions (left). At high shock velocities, the broad neutral distribution function is skewed and offset from that of the protons, leading to asymmetric velocity profiles with the broad neutrals moving considerably faster than the ions (right).

TABLE 1
INFERRED SHOCK VELOCITIES AND TEMPERATURE EQUILIBRATIONS FROM SNRS

Object	Reference	H α FWHM (km s $^{-1}$)	H α I_b/I_n	H α v_s (km s $^{-1}$)	β	Ly β FWHM (km s $^{-1}$)	Ly β v_s (km s $^{-1}$)
Cygnus	G01	262 \pm 32	0.59 \pm 0.3
RCW 86	G07b	325 \pm 10	1.06 \pm 0.1	367 $^{+46}_{-26}$	0.577 $^{+0.287}_{-0.136}$
	G01	562 \pm 18	1.18 \pm 0.03	690 $^{+9}_{-131}$	0.826 $^{+0.030}_{-0.510}$
	G07b	640 \pm 35	1.0 \pm 0.2	673 $^{+89}_{-83}$	0.376 $^{+0.317}_{-0.328}$
SNR 0505–67.9	S91	580 \pm 70	≥ 0.7
	G03	785 $^{+95}_{-80}$	0.93 $^{+0.11}_{-0.10}$	808 $^{+104}_{-150}$	0.331 $^{+0.071}_{-0.225}$
	G03	1055 $^{+150}_{-120}$	0.88 $^{+0.11}_{-0.10}$	1024 \pm 107	0.194 $^{+0.103}_{-0.126}$
	G07a	1135 \pm 30	1036 \pm 46 ^a
	G07a	1365 \pm 75	1249 \pm 89 ^a
SNR 0548–70.4	S91	760 \pm 140	1.1 \pm 0.2	814 $^{+153}_{-199}$	0.430 $^{+0.150}_{-0.194}$
SNR 0519–69.0	S91	1300 \pm 200	0.8 \pm 0.2	1178 $^{+185}_{-157}$	0.031 $^{+0.176}_{-0.001}$
	T82	2800 \pm 300	0.4–0.8 ^b
	G07a	3130 \pm 155	2984 $^{+703}_{-185}$
Kepler.....	F89	1750 \pm 200	1.1 \pm 0.25	1589 $^{+191}_{-182}$	0.035 $^{+0.010}_{-0.009}$
Tycho.....	G01	1765 \pm 110	0.67 \pm 0.1	1606 $^{+108}_{-103}$	0.046 $^{+0.007}_{-0.006}$
	KWC87	1800 \pm 100	1.08 \pm 0.16	1634 $^{+95}_{-91}$	0.036 $^{+0.006}_{-0.005}$
	S91	1900 \pm 300	0.77 \pm 0.09	1733 $^{+285}_{-280}$	0.047 $^{+0.018}_{-0.011}$
SN 1006.....	G02	2290 \pm 80	0.84 $^{+0.03}_{-0.01}$	2098 $^{+79}_{-77}$	0.055 $^{+0.002}_{-0.006}$
	S91	2310 \pm 210	0.73 \pm 0.06	2126 $^{+257}_{-205}$	0.069 $^{+0.030}_{-0.018}$
	KWC87	2600 \pm 100	0.77 \pm 0.08	2426 $^{+364}_{-128}$	0.073 $^{+0.125}_{-0.049}$

NOTES.—Case A conditions are assumed for the broad line. The preshock ionization fraction is taken to be $f_p = 0.5$.

^a Values are reported for the range $0.01 \leq \beta \leq 0.1$.

^b We regard this as 0.6 ± 0.2 when solving for β and v_s .

REFERENCES.—(F89) Fesen et al. 1989; (G01) Ghavamian et al. 2001; (G02) Ghavamian et al. 2002; (G03) Ghavamian et al. 2003; (G07a) Ghavamian et al. 2007a; (G07b) Ghavamian et al. 2007b; (KWC87) Kirshner et al. 1987; (S91) Smith et al. 1991; (T82) Tuohy et al. 1982.

$\beta = 0.01, 0.1, 0.5$, and 1 at shock velocities of $v_s = 1000$ km s $^{-1}$ (left) and $v_s = 7000$ km s $^{-1}$ (right). At low shock velocities, the charge transfer is extremely efficient, yielding broad neutral distribution functions that are nearly identical to those of the thermal protons, with the broad neutrals moving slightly more slowly than the ions. At high shock velocities, the broad neutral distributions are skewed and offset from the proton distribution, leading to the asymmetric profiles depicted in the right panel, with the broad neutrals moving considerably faster than the ions.

We write the FWHM relation as a function of v_s and β :

$$W_{\text{FWHM}} = W(v_s, \beta), \quad (38)$$

where W_{FWHM} and W represent the measured and theoretical values of the line FWHM, respectively. For SNRs with measured values of W_{FWHM} and I_b/I_n , combining equations (37) and (38) self-consistently constrains v_s and β . This is accomplished by inverting $W(v_s, \beta)$ to yield $v_s = W^{-1}(W_{\text{FWHM}}, \beta)$ and calculating the root of the expression

$$L_{bn}(W^{-1}(W_{\text{FWHM}}, \beta), \beta) - I_b/I_n = 0. \quad (39)$$

This procedure produces a pair of values (v_s, β) for which the theoretical calculations L_{bn} and W equal the measured I_b/I_n and W_{FWHM} .

We emphasize that a self-consistent calculation is critical for accurately determining the values of v_s and β . In the previous literature, two bracketing values for the temperature equilibration are sometimes chosen (e.g., $\beta = 0.1$ and 1), and the FWHM relation is employed to give a range of possible values for the shock velocity (e.g., Paper I). However, this procedure has a major flaw. In practice, the self-consistent quantity $L_{bn}(W^{-1}(W_{\text{FWHM}}, \beta), \beta)$ will have a minimum value over the range $\beta \in (m_e/m_p, 1)$. If the value of I_b/I_n for a particular observed shock is less than this

minimum value, then no pair (v_s, β) will yield the observed values I_b/I_n and W_{FWHM} . In such a case, the model breaks down, and quoting a range of possible shock velocities for two bracketing values of β is inappropriate. Additional physics must be invoked in order to account for the observations. Moreover, when no measurement of I_b/I_n exists, the bracketing procedure may or may not yield an accurate estimate of the range of the shock velocities.

We use the methodology described above to self-consistently extract shock parameters from observations of Balmer-dominated SNRs. We summarize our results in Table 1, which lists, from left to right, the object name, the reference, the values of the H α W_{FWHM} and I_b/I_n , calculated values for v_s and β from H α , the Ly β W_{FWHM} , and calculated values for v_s from Ly β . If our models do not yield a fit to the observations, we do not list values for the shock velocity and the temperature equilibration ratio. Our calculations show a characteristic range of β between 0.01 and 0.1 for $v_s > 1000$ km s $^{-1}$. Therefore, in the case of Ly β observations for which no measurement of I_b/I_n exists, we report the derived shock velocities for the range of $\beta = 0.01$ – 0.1 ; further observations are required in order to confirm the accuracy of these estimates. For SNR 0519–69.0, observations exist in both H α and Ly β . To calculate the shock velocities from the Ly β observations, we use the derived value of β from the H α diagnostics.

Our models successfully fit 14 measurements from seven Balmer-dominated SNRs, within the observational uncertainties. In the cases of the Cygnus Loop (G01) and one measurement from SNR 0519–69.0 by T82, the observed ratio I_b/I_n is too low to be accounted for by our calculations. In addition, we are unable to fit the majority of measurements from the object DEM L71/SNR 0505–67.9, which we have omitted from the table. Below, we discuss a possible explanation for these discrepancies in our model.

We note that our inferred shock velocities and temperature equilibration ratios are systematically lower than those in previous

studies and show an increased sensitivity to β compared to those of Papers I and II. The ability to sensitively probe both the shock velocity and the temperature equilibration is of interest for studies of collisionless electron heating in shocks, as described below. Typically, inferred values for v_s are $\sim 10\%$ – 30% smaller than those quoted in Paper I, primarily because of the contribution of broad neutral velocities to the relative speeds in fast neutral-ion interactions. The shock speeds in Tycho’s SNR and SN 1006 are smaller by about 15% and 27%, respectively, than those reported previously. One implication is that the distances to these SNRs derived from the shock speeds and proper motions are correspondingly smaller. The inferred distance to SN 1006 is reduced from 2.18 kpc (Winkler et al. 2003) to 1.6 kpc. The corresponding brightness of the SN is squarely in the middle of the Type I SN distribution at 2.18 kpc, but is 0.7 mag fainter at 1.6 kpc. The smaller distance is more comfortable in comparison with that of the Schweizer-Middleditch star, which lies between 1.05 and 2.1 kpc and whose spectrum shows absorption by SN 1006 ejecta (Burleigh et al. 2000). On the other hand, the ejecta are observed to expand at 7026 km s^{-1} (Hamilton et al. 2007), and the requirement that this material lie within the remnant places a conservative lower limit to the distance of 1.6 kpc, which is just consistent with that derived here.

6. DISCUSSION

6.1. Collisionless Shock Heating

One of the motivations for studying Balmer-dominated SNRs is to probe the physics of collisionless shocks. While various mechanisms for transferring energy from proton to electron populations have been proposed in the literature, there is no consensus on how to predict the value of β for a given set of shock parameters in astrophysical contexts. In a recent paper, G07b attempted to address this problem by deriving a relation for the temperature equilibration ratio versus shock velocity from observations of shocks in SNRs. They reported that their results could be fitted by a curve of $\beta(v_s) \propto v_s^{-2}$, with $\beta = 1$ for $v_s \leq 400 \text{ km s}^{-1}$. Given that the proton temperature scales roughly as $T_p \propto v_s^2$, this conclusion implies that the electrons are heated to a constant temperature, independently of the shock velocity. While G07b discussed a possible mechanism for this dependence, it has not yet been established that theoretical models can produce such an effect.

Using our new model for emission from Balmer-dominated shocks, we have calculated an updated $\beta(v_s)$ relation, which is displayed in Figure 13. The solid curve depicts the proposed v_s^{-2} dependence, with $\beta = 1$ for $v_s < 400 \text{ km s}^{-1}$. For all the points shown in the plot, we have self-consistently fitted measurements of both W_{FWHM} and I_b/I_n . We fix the preshock ionization at $f_p = 0.5$, and we exclude the majority of measurements from DEM L71/SNR 0505–67.9, as well as one measurement each from Cygnus and SNR 0519–69.0 that we are unable to account for with our models. We find from our calculations that the data are not well fitted by a power-law relation $\beta(v_s) \propto v_s^{-\alpha}$; if we set $\alpha = 2$, the fit to the β curve yields a reduced χ^2 of $\chi_r^2 = 62.8/12 = 5.2$.⁵ For shock velocities of $v_s \lesssim 1500 \text{ km s}^{-1}$, our results are fairly close to those of previous models. At higher velocities, $v_s > 2000 \text{ km s}^{-1}$, the deviations are more significant. The minimum value of the temperature equilibration ratio is $\beta \sim 0.03$ at velocities of $v_s \gtrsim 1500 \text{ km s}^{-1}$. This value is greater than the theoretical minimum $\beta = m_e/m_p$ by several orders of magnitude, but it is smaller than the value predicted by some

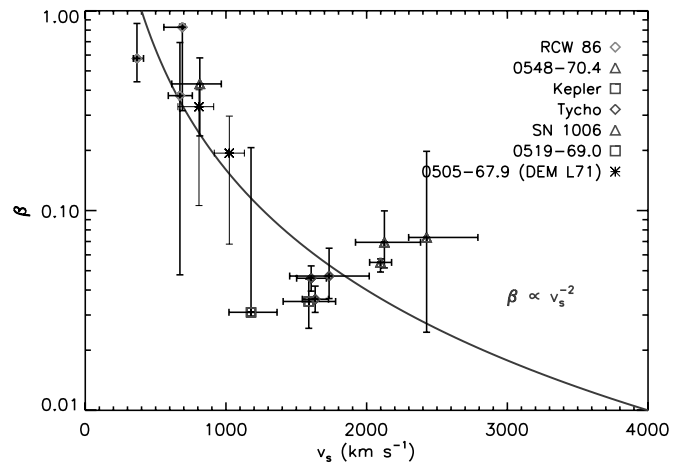


FIG. 13.—Temperature equilibration ratio β as a function of the shock velocity v_s for Balmer-dominated shocks fitted by our calculations. The symbol shape denotes the origin of each data point from the SNRs in our sample. The solid curve shows the dependence $\beta(v_s) \propto v_s^{-2}$ for a heating mechanism that is independent of shock velocity. Fitting this curve to the inferred values from our new model yields a reduced χ^2 of $\chi_r^2 = 62.8/12 = 5.2$. [See the electronic edition of the Journal for a color version of this figure.]

collisionless heating models (e.g., Cargill & Papadopoulos 1988).

6.2. Limitations to Our Model and Future Work

As demonstrated by SNR 0505–67.9, the Cygnus Loop, and SNR 0519–69.0, an additional physical mechanism is needed to account for the observed $H\alpha$ emission seen in some SNRs. One possibility is that a significant amount of the dissipated energy in the shock is transferred to cosmic rays, producing a precursor that can heat and accelerate the upstream gas, altering the shock jump conditions (S91; Hester et al. 1994; Sollerman et al. 2003). In addition, the precursor can “push” on the upstream protons, leading to a velocity differential between the neutrals and charged species (e.g., Berezhko & Ellison 1999). Furthermore, all previous models have assumed that the kinetic distribution functions for protons and electrons in the shock transition zone are Maxwellian. However, recent calculations have shown that the proton distribution can significantly depart from Maxwellian behavior, forming a distinct “pickup” ion population (Raymond et al. 2008). This effect will change the structure of the transition zone, the broad neutral distributions, and the kinetic reaction rates.

The inclusion of a cosmic-ray precursor will affect the predicted broad-to-narrow line strength in several ways. Broadening of the cold neutral distribution function effectively reduces the optical depth to $\text{Ly}\beta$ scattering, which will decrease the conversion efficiency of narrow $\text{Ly}\beta$ to $H\alpha$ and increase L_{bn} . In contrast, additional excitation of cold neutral hydrogen atoms in the precursor will increase narrow-line emission and decrease L_{bn} . While it is not obvious which effect is dominant, evidence for such precursor effects exists in the anomalously large widths of $H\alpha$ lines in most Balmer-dominated SNRs (Sollerman et al. 2003). In addition, $H\alpha$ emission from a spatially resolved precursor in Tycho’s SNR has recently been reported by Lee et al. (2007).

The spatial structure of the shock transition zone provides a way to infer the external density n_0 and the preshock ionization fraction f_p . As noted in § 5.2, the spatial emissivity profile and the centroid shift between the narrow and broad components have a strong dependence on f_p and n_0 . Raymond et al. (2007) were able to spatially resolve the total $H\alpha$ emission from SN 1006 using the Advanced Camera for Surveys (ACS) instrument on the

⁵ Note that such a large value of χ_r^2 should be interpreted in the context of the 10% uncertainties associated with the atomic cross sections used.

Hubble Space Telescope (HST). Raymond et al. (2007) adopted the values of v_s and β from those inferred by G02 using observations of the FWHM and I_b/I_n , and then they attempted to fit for the spatial structure by varying f_p and n_0 . Future instruments similar to the ACS with high spatial resolution (e.g., the WFC3 on board the *HST*, with the narrow $H\alpha$ component isolated by a narrowband filter) will allow for more accurate measurements of the transition zone structure to be made and will provide additional constraints on the shock parameters.

The models presented in our series of papers generically describe the physics of nonradiative shocks interacting with cold, preshock gas. In addition to using them to understand nearby SNRs, we can also calculate hydrogen line emission from SNRs in young, distant galaxies (with redshifts of 3–5), a subject first explored by Heng & Sunyaev (2008). We improve on their estimates of the luminosity ratios of $\text{Ly}\alpha$ and $\text{Ly}\beta$ to $H\alpha$, denoted $\Gamma_{\text{Ly}\alpha/H\alpha}$ and $\Gamma_{\text{Ly}\beta/H\alpha}$, respectively (see Fig. 14; cf. Fig. 1 of Heng & Sunyaev 2008). Heng & Sunyaev (2008) underestimated the broad $\text{Ly}\alpha$ emission at high shock velocities because of their neglect of the multicomponent shock transition zone. Nevertheless, their conclusions remain intact: the sensitivity of $\Gamma_{\text{Ly}\alpha/H\alpha}$ and $\Gamma_{\text{Ly}\beta/H\alpha}$ to β by a factor of ~ 2 over the velocity range $1000 \lesssim v_s \lesssim 4000 \text{ km s}^{-1}$ is a direct and unique way to measure the temperature ratio. A valuable extension to their work will incorporate models of $\text{Ly}\alpha$ and $\text{Ly}\beta$ scattering, taking into account geometric effects.

The biggest systematic uncertainties in our current model are (1) a proper treatment of Lyman line scattering in the shock transition zone, (2) the effect of a cosmic-ray precursor on L_{bn} , and (3) inclusion of a nonthermal population of protons in calculating the reaction kinetics. The incorporation of these physical effects will be important for future models. Our new results yield substantial differences in the derived shock speed (and hence the inferred distances) for measurements of Balmer-dominated SNRs compared to those of previous models. While our results are not

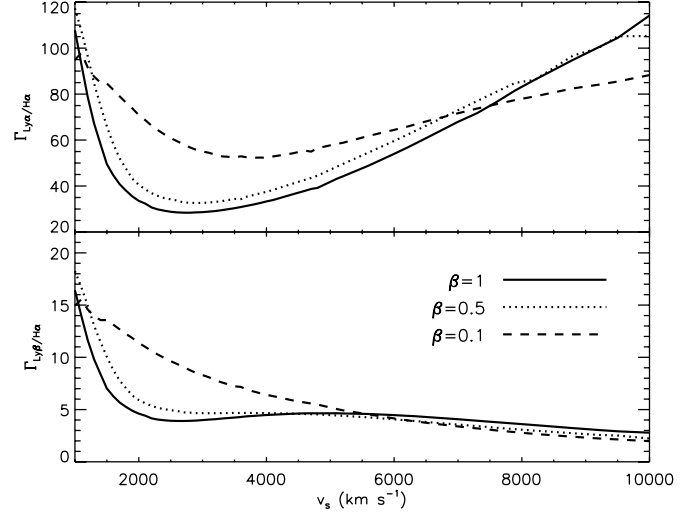


FIG. 14.—Luminosity ratios of broad $\text{Ly}\alpha$ and $\text{Ly}\beta$ to $H\alpha$, denoted $\Gamma_{\text{Ly}\alpha/H\alpha}$ and $\Gamma_{\text{Ly}\beta/H\alpha}$, respectively, as a function of shock velocity (cf. Fig. 2 of Heng & Sunyaev 2007).

consistent with a power-law relation between the shock velocity and the electron temperature, future work is needed to resolve the remaining approximations in our model and to make a more definitive statement about the shock heating of electrons.

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APPENDIX A

SHOCK KINETICS

Here we summarize the important details of our calculations of the rate coefficients, broad neutral velocity distributions, and broad-line profiles. For a more detailed description of our methods, see Paper I. In this work, we treat the nl atomic sublevels separately instead of considering a single, summed n level as in Paper I. We treat charge transfers, excitation, and ionization events between electrons, protons, alpha particles, and hydrogen atoms using the cross sections of Barnett et al. (1990), Belkić et al. (1992), Janev & Smith (1993), Balança et al. (1998), and Harel et al. (1998). Fitting functions for some of these cross sections are provided in Heng & Sunyaev (2008).

We also consider the ionization of helium atoms and singly ionized helium by electrons and protons, using the cross sections of Peart et al. (1969), Angel et al. (1978), Rudd et al. (1983), Shah & Gilbody (1985), Rinn et al. (1986), and Shah et al. (1988). Following Barnett et al. (1990), we fit the cross sections to the function

$$\mathcal{F}(x; \mathbf{A}) = \exp \left[\frac{A_0}{2} C_0 + \sum_{i=1}^8 A_i C_i(x) \right], \quad (\text{A1})$$

where the components of $\mathbf{A} = (A_0, A_2, \dots, A_8)$ are the fitting parameters and the quantities C_i are the Chebyshev polynomials:

$$C_0(x) = 1, \quad (\text{A2})$$

$$C_1(x) = x, \quad (\text{A3})$$

$$C_i(x) = 2xC_{i-1} - C_{i-2}. \quad (\text{A4})$$

We define the fitting variable x as

$$x = \frac{\ln[E^2/(E_{\max}E_{\min})]}{\ln(E_{\max}/E_{\min})}, \quad (\text{A5})$$

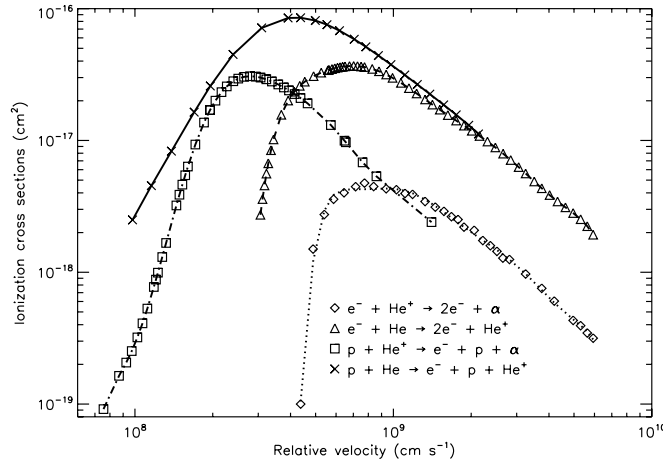


FIG. 15.—Cross sections for the ionization of neutral and singly ionized helium by electrons and protons. For references to the atomic data, see Appendix A.

where E is the relative energy between the interacting particles and E_{\min} and E_{\max} are the minimum and maximum energies for which data are available. We assume a fiducial error of 10% for the data. The cross sections and corresponding fits are displayed in Figure 15, and the fitting parameters are presented in Table 2.

We neglect the charge transfer reactions of neutral and singly ionized helium with protons due to the scarcity of cross sections for these processes in our velocity range of interest. While such interactions may affect helium line emission, we do not expect them to strongly couple to the Balmer radiation. In addition, for ease of computation, we approximate the excitation and ionization of hydrogen by He^+ , using relevant proton rate coefficients.⁶ However, cross sections for this process do exist in the literature and should be used in future calculations (Barnett et al. 1990). Nevertheless, we expect corrections to these cross sections to have a small quantitative effect on our results.

We make several approximations in order to speed up our computations. For temperature ratios of $0.01 \lesssim \beta \lesssim 0.1$, the velocity width of the electron distribution is generally broader than that of the broad neutrals (the characteristic electron velocities are greater than the proton velocities by a factor of 4–14), implying that the rate coefficients for interactions involving electrons are insensitive to changes in the broad neutral velocity distribution. We therefore approximate the electron rate coefficients to be the same for reactions involving cold atoms and broad neutrals. Calculations of the FWHM for line profiles should include excitations by electrons, protons, singly ionized helium, and alpha particles. However, alpha particles contribute $\lesssim 1\%$ due to the relatively smaller density n_α and can typically be omitted.

The reaction rate coefficients used in equations (4)–(11) are defined as

$$\tilde{R}_{X,s} = \int d^3\mathbf{v} \int d^3\mathbf{v}' f_a(\mathbf{v}) f_s(\mathbf{v}') \Delta v \sigma_{X,s}(\Delta v), \tag{A6}$$

where X denotes the interaction (ionization or charge transfer), a denotes the atomic species (broad or cold neutrals), and s denotes the interacting charged particle (electrons, protons, singly ionized helium, or alpha particles). In contrast to equations (33)–(35), the interaction cross sections $\sigma_{X,s}$ used here represent the sum of the reactions to all nl levels for which atomic data are available (cf. eq. [27]). Using the definition of equation (A6), the quantity $n_a n_s \tilde{R}_{X,s}$ gives the number of interactions X between the species a and s per unit volume.

⁶ Our assumption is based on the Weizsacker-Williams approximation, in which scattering dynamics are dominated by the charge of the impacting particle, as opposed to its mass (e.g., Jackson 1998). This approximation is valid when the relative velocity of the collision is greater than that of an electron orbiting the hydrogen atom. At low velocities, $v_s \lesssim 300 \text{ km s}^{-1}$, this assumption breaks down.

TABLE 2
FITS TO HELIUM SPECIES IONIZATION CROSS SECTIONS

Parameter	He + e ⁻	He ⁺ + e ⁻	He + p	He ⁺ + p
A_0	-78.4712	-82.6155	-77.0261	-80.7740
A_1	-0.832236	-0.535745	0.596233	1.50686
A_2	-1.00452	-1.15893	-1.37165	-2.04982
A_3	0.482606	0.644513	0.205854	0.384422
A_4	-0.244927	-0.419765	0.123038	0.353731
A_5	0.121965	0.299795	-0.0671246	-0.268015
A_6	-0.0795100	-0.204216	-0.0135306	-0.0107265
A_7	0.0537985	0.133238	0.0163305	0.0101069
A_8	-0.0521662	-0.0716388	-0.00556328	-0.110770
E_{\min} (eV).....	26.6	54.5	5.0×10^3	2.98×10^3
E_{\max} (eV).....	10^4	10^4	2.38×10^6	1.03×10^6

APPENDIX B

TYPOGRAPHICAL ERRORS IN PAPERS I AND II

We point out several minor typographical errors in Heng & McCray (2007) and Heng et al. (2007). In § 5.1 of Heng & McCray (2007), $\mathcal{R}_{bn}(H\alpha) = I_b(H\alpha)/I_b(H\alpha)$ should be $\mathcal{R}_{bn}(H\alpha) = I_b(H\alpha)/I_n(H\alpha)$. In § 5 of Heng et al. (2007), $R_{H\alpha b_0}$ should be $R_{H\alpha, b_0}$, $R_{H\alpha b}$ should be $R_{H\alpha, b}$, and $R_{H\alpha n}$ should be $R_{H\alpha, n}$. In § 6.2, paragraph 2 of Heng et al. (2007), the sentence that begins “Before CKR80 and HM07, . . .” should be changed to “Before (CKR80 and HM07), . . .”

REFERENCES

- Angel, G. C., Dunn, K. F., Sewell, E. C., & Gilbody, H. B. 1978, *J. Phys. B*, 11, L49
- Balança, C., Lin, C. D., & Feautrier, N. 1998, *J. Phys. B*, 31, 2321
- Barnett, C. F., Hunter, H. T., Fitzpatrick, M. I., Alvarez, I., Cisneros, C., & Phaneuf, R. A. 1990, NASA STI/Recon Technical Report N, 91, 13238
- Belkić, D., Gayet, R., & Salin, A. 1992, *At. Data Nucl. Data Tables*, 51, 59
- Berezhko, E. G., & Ellison, D. C. 1999, *ApJ*, 526, 385
- Burleigh, M. R., Heber, U., O’Donoghue, D., & Barstow, M. A. 2000, *A&A*, 356, 585
- Cargill, P. J., & Papadopoulos, K. 1988, *ApJ*, 329, L29
- Chevalier, R. A., Kirshner, R. P., & Raymond, J. C. 1980, *ApJ*, 235, 186
- Chevalier, R. A., & Raymond, J. C. 1978, *ApJ*, 225, L27
- Cox, A., ed. 2000, *Allen’s Astrophysical Quantities* (4th ed.; New York: AIP)
- Fesen, R. A., Becker, R. H., Blair, W. P., & Long, K. S. 1989, *ApJ*, 338, L13 (F89)
- Ghavamian, P., Blair, W. P., Sankrit, R., Raymond, J. C., & Hughes, J. P. 2007a, *ApJ*, 664, 304 (G07a)
- Ghavamian, P., Laming, J. M., & Rakowski, C. E. 2007b, *ApJ*, 654, L69 (G07b)
- Ghavamian, P., Rakowski, C. E., Hughes, J. P., & Williams, T. B. 2003, *ApJ*, 590, 833 (G03)
- Ghavamian, P., Raymond, J., Smith, R. C., & Hartigan, P. 2001, *ApJ*, 547, 995 (G01)
- Ghavamian, P., Winkler, P. F., Raymond, J. C., & Long, K. S. 2002, *ApJ*, 572, 888 (G02)
- Hamilton, A. J. S., Fesen, R. A., & Blair, W. P. 2007, *MNRAS*, 381, 771
- Harel, C., Jouin, H., & Pons, B. 1998, *At. Data Nucl. Data Tables*, 68, 279
- Heng, K., & McCray, R. 2007, *ApJ*, 654, 923 (Paper I)
- Heng, K., & Sunyaev, R. A. 2008, *A&A*, 481, 117
- Heng, K., van Adelsberg, M., McCray, R., & Raymond, J. C. 2007, *ApJ*, 668, 275 (Paper II)
- Hester, J. J., Raymond, J. C., & Blair, W. P. 1994, *ApJ*, 420, 721
- Jackson, J. D. 1998, *Classical Electrodynamics* (3rd ed.; New York: Wiley)
- Janev, R. K., & Smith, J. J. 1993, *Cross Sections for Collision Processes of Hydrogen Atoms with Electrons, Protons and Multiply Charged Ions* (Vienna: Int. Atomic Energy Agency)
- Kirshner, R., Winkler, P. F., & Chevalier, R. A. 1987, *ApJ*, 315, L135 (KWC87)
- Laming, J. M., Raymond, J. C., McLaughlin, B. M., & Blair, W. P. 1996, *ApJ*, 472, 267
- Lee, J.-J., Koo, B.-C., Raymond, J., Ghavamian, P., Pyo, T.-S., Tajitsu, A., & Hayashi, M. 2007, *ApJ*, 659, L133
- Peart, B., Walton, D. S., & Dolder, K. T. 1969, *J. Phys. B*, 2, 1347
- Raymond, J. C., Isenberg, P. A., & Laming, J. M. 2008, *ApJ*, 682, 408
- Raymond, J. C., Korreck, K. E., Sedlacek, Q. C., Blair, W. P., Ghavamian, P., & Sankrit, R. 2007, *ApJ*, 659, 1257
- Rinn, K., Melchert, F., Rink, K., & Salzborn, E. 1986, *J. Phys. B*, 19, 3717
- Rudd, M. E., Goffe, T. V., Dubois, R. D., Toburen, L. H., & Ratcliffe, C. A. 1983, *Phys. Rev. A*, 28, 3244
- Rybicki, G. B., & Lightman, A. P. 1979, *Radiative Processes in Astrophysics* (New York: Wiley)
- Shah, M. B., Elliott, D. S., McCallion, P., & Gilbody, H. B. 1988, *J. Phys. B*, 21, 2751
- Shah, M. B., & Gilbody, H. B. 1985, *J. Phys. B*, 18, 899
- Shull, J. M., & McKee, C. F. 1979, *ApJ*, 227, 131
- Smith, R. C., Kirshner, R. P., Blair, W. P., & Winkler, P. F. 1991, *ApJ*, 375, 652 (S91)
- Sollerman, J., Ghavamian, P., Lundqvist, P., & Smith, R. C. 2003, *A&A*, 407, 249
- Spitzer, L. 1962, *Physics of Fully Ionized Gases* (2nd rev. ed.; New York: Wiley)
- Tuohy, I. R., Dopita, M. A., Mathewson, D. S., Long, K. S., & Helfand, D. J. 1982, *ApJ*, 261, 473 (T82)
- Winkler, P. F., Gupta, G., & Long, K. S. 2003, *ApJ*, 585, 324