ON THE HEATING OF THE SOLAR CORONA AND THE ACCELERATION OF THE LOW-SPEED SOLAR WIND BY ACOUSTIC WAVES GENERATED IN THE CORONA

TAKERU KEN SUZUKI
Division of Theoretical Astrophysics, National Astronomical Observatory, 2-21-1 Osawa, Mitaka, Tokyo 181-8588, Japan; and Department of Astronomy, Faculty of Science, University of Tokyo, 7-3-1 Hongo, Bunkyo-ku, Tokyo 113-0033, Japan; stakeru@th.nao.ac.jp

Received 2002 March 26; accepted 2002 June 16

ABSTRACT

We investigate possibilities of solar coronal heating by acoustic waves generated not in the photosphere but in the corona, focusing on heating in the mid- to low-latitude corona, where the low-speed wind is expected to come from. Acoustic waves of period $\tau \sim 100$ s are triggered by chromospheric reconnection, one model of small-scale magnetic reconnection events recently proposed by Sturrock. These waves, having a finite amplitude, eventually form shocks to shape sawtooth waves (N-waves) and directly heat the surrounding corona by dissipation of their wave energy. Outward propagation of the N-waves is treated based on weak-shock theory, so that the heating rate can be evaluated consistently with physical properties of the background coronal plasma without setting a dissipation length in an ad hoc manner. We construct coronal structures from the upper chromosphere to outside 1 AU for various acoustic wave inputs, with a range of energy flux of $F_{a,0} = (1-20) \times 10^5$ ergs cm$^{-2}$ s$^{-1}$ and a period of $\tau = 60-300$ s. The heating by the N-wave dissipation works effectively in the inner corona, and we find that waves of $F_{a,0} \geq 2 \times 10^5$ ergs cm$^{-2}$ s$^{-1}$ and $\tau \geq 60$ s could maintain the peak coronal temperature, $T_{\text{max}} > 10^6$ K. The model could also reproduce the density profile observed in the streamer region. However, due to its short dissipation length, the location of $T_{\text{max}}$ is closer to the surface than in observation, and the resulting flow velocity of the solar wind is lower than the observed profile of the low-speed wind. Cooperation with other heating and acceleration sources with larger dissipation lengths are inevitably needed to reproduce the real solar corona.

Subject headings: solar wind — Sun: corona — waves

1. INTRODUCTION

The heating of the solar corona is still poorly understood, and it is one of the most challenging but interesting questions to be solved in astrophysics. The origin of the energy that heats the corona is generally believed to lurk in the turbulent convective motions beneath the photosphere. A certain fraction of the kinetic energy of those turbulent motions is carried up to the corona in the shape of nonthermal energy, such as magnetic or wave energy, and thermalization of such energy in the corona results in the heating of the surrounding plasma. Granule motions of surface convection are simply expected to cause wave generation at the photospheric level. The possibility of coronal heating by those waves has been investigated by many researchers (Osterbrock 1961; Ulmschneider 1971; McWhirter, Thome mann, & Wilson 1975). To date, among the various modes of waves, although Alfvén waves have been widely taken up as a convincing candidate for coronal heating and the acceleration of the high-speed wind in the polar coronal holes through the ion-cyclotron damping mechanism (e.g., Cranmer, Field, & Kohl 1999; Hollweg 1999), acoustic waves have not been regarded as a major heating source of the corona because of their dissipative character. Acoustic waves with a finite amplitude inevitably steepen their wave fronts to form shocks when they travel. They are strongly damped through upward propagation in the chromosphere; consequently, only 0.01% of the initial wave energy can reach the corona (Stein & Schwartz 1972, hereafter SS72). The role of acoustic waves in solar coronal heating has been buried in oblivion for a long time.

However, recent observations of dynamical structures of the solar corona (see Aschwanden, Poland, & Rabin 2001 for a recent review) highlight the role of acoustic waves generated not in the photosphere but in the corona. A large number of flares and flarelike events (e.g., Tsuneta et al. 1992; Tsuneta 1996) have been observed by various telescopes up to the present (e.g., Nishio et al. 1997; see Bastian, Benz, & Gary 1998 for a review) in a wide range of energy from $\sim 10^{24}$ to $\sim 10^{32}$ ergs, and it has been statistically argued whether they could supply sufficient energy for heating the global corona. Recent measurements of the EUV frequency distribution ($dN/dE \propto E^{-\alpha}$ of flarelike events give a steeper power-law index ($\alpha = 2.3-2.6$: Krucker & Benz 1998; $\alpha = 2.0-2.6$: Parnell & Jupp 2000) on the lower energy side ($E = 10^{24}-10^{25}$ ergs). This fact makes us infer that small-scale flarelike events, such as nanoflares, might be sufficient (Hudson 1991) for the required energy budget ($\geq 10^{3.5}$ ergs cm$^{-2}$ s$^{-1}$; Withbroe & Noyes 1977). Several models taking into account such small-scale events have been introduced (Sturrock 1999, hereafter S99; Sturrock, Roald, & Wolfson 1999; Roald, Sturrock, & Wolfson 2000, hereafter RSW00; Tarbell, Ryutova, & Covington 1999; Sakai et al. 2000; Furusawa & Sakai 2000). Among these models, S99 proposed a model of chromospheric magnetic reconnection. In this model, a flux tube, newly formed as a result of reconnection events, oscillates vertically to excite acoustic waves in the corona. RSW00 further showed with a simple kinetic model that this process possibly liberates sufficient energy in the corona with the correct relevant parameters. A different mechanism also predicts the generation of acoustic waves in the corona. For example, spicules in the open magnetic field region effectively transport the energy of random motions in the photosphere into the corona and produce longitudinal (acoustic) waves there (Hollweg 1992; Kudoh & Shibata 1999). Thus, acoustic waves are supposed
to be generated constantly in the corona by various dynamical processes. Such acoustic waves could heat the surrounding corona directly, unlike acoustic waves created at the photospheric level. They are expected to effectively heat the inner corona by shock dissipation. Therefore, the density at the coronal base might be large enough to create the dense wind flow observed as the streamer in the mid- to low-latitude region, which is believed to be connected to the low-speed solar wind (Habbal et al. 1997). Acoustic waves might become one of the relevant processes in the region that generates the low-speed wind of which mechanisms have been poorly elucidated.

In the above models that introduce the production of acoustic waves in the corona, the authors concentrated on the amount of kinetic (wave) energy released in the corona. However, heating of coronal gas is accomplished by the thermalization of wave energy through the dissipation of such acoustic waves. Therefore, an appropriate treatment of the propagation and dissipation of these waves is indispensable in understanding the problem of coronal heating. In this paper, we employ a formulation of the weak-shock theory constructed by SS72, originally developed to study acoustic waves generated by convective motions in the photosphere, for propagation and dissipation of the waves produced in the corona. Then, we can explicitly determine the rate of heating by dissipation of the waves for a given wave energy flux (or amplitude) and wave period without tuning other free parameters; we do not have to arbitrarily take the exponential type of heating for mechanical energy input, \( F_m \propto \exp[-(r - R_\odot)/l_m] \), for an assumed constant dissipation length \( l_m \), which is poorly supported by fundamental physical processes, although this conventional shape of the heating law was adopted for lack of an alternative in most previous models (Kopp & Orrall 1976; Withbroe 1988; Sandbæk & Leer 1994) to study the global coronal structure.

The paper is organized as follows: In \( \S \) 2 we present our model. We briefly summarize the generation of acoustic waves in the corona (\( \S \) 2.1) and the weak-shock theory formulated by SS72 (\( \S \) 2.2). Basic equations describing the global coronal structure from 1 \( R_\odot \) to \( \gtrsim 1 \) AU are shown in \( \S \) 2.3, and we present our method for constructing a unique coronal solution with given wave parameters. In \( \S \) 3 we show our model results. First, we discuss the variation of coronal wind structures with respect to input parameters from a theoretical point of view (\( \S \) 3.1). Second, we examine several characteristic properties of the resulting corona (\( \S \) 3.2). Finally, we test whether our model is able to reproduce recent observed results for the low-speed wind in the mid- to low-latitude region (\( \S \) 3.3). In \( \S \) 4 we summarize our results and discuss related topics.

2. MODEL

2.1. Generation of Acoustic Waves in the Corona

As a new relevant process of coronal heating in the quiet-Sun region, S99 proposed a model of network-field magnetic reconnection at the chromospheric level. The chromosphere, specifically at the minimum-temperature location, is quite a favorable site for reconnection to occur, since the magnetic resistivity is greatest and the width of the current sheet, which is possibly scaled with the pressure scale height, is smallest there. Such reconnection events form a new closed magnetic flux tube, which is initially located far from the equilibrium state and hence will spring upward quickly (consult S99 for a schematic picture). The tube will oscillate about the equilibrium state normal to the solar surface with a period of \( \tau = 2L/v_A \), where \( v_A \) is the Alfvén velocity and \( L \) is the length of the tube. If we use the quoted values in S99 as a typical example of the tube, \( L \approx 1.5 \times 10^8 \) cm, the mean magnetic field \( \sim 100 \) G, and the mean density \( \sim 10^{-12} \) g cm\(^{-3} \), the oscillation periods would be \( \tau \approx 100 \) s. It could keep oscillating several tens of times before it was damped (S99). As a result, these perpendicular oscillations would excite longitudinal waves traveling in the vertical direction. Given that the configurations of the magnetic fields are perpendicular above the flux tube, such waves would propagate upward as acoustic waves in the rarefied atmosphere. They must contribute to the heating of the surrounding corona directly by the dissipation of the waves.

Using a simple kinetic model, RSW00 have estimated the energy liberated by the above mechanism as a function of mean magnetic field strength at the photospheric level. They showed that a mean field of \( \sim 10 \) G could generate the required heating rate (\( \gtrsim 10^{23} \) ergs cm\(^{-2} \) s\(^{-1} \)) to explain the conditions in the quiet-Sun region. A sizable fraction of the released energy is supposed to be transported to the energy of the acoustic waves. In our model, we parameterize the input energy flux of the acoustic waves (not the total energy liberated by the reconnection) as \( F_{\text{w},\theta} \) (ergs cm\(^{-2} \) s\(^{-1} \)). In this paper, to study the physical processes clearly, we focus on the role of such acoustic waves in the coronal heating and construct the coronal structure with the input \( F_{\text{w},\theta} \) and wave period \( \tau \), although some of the released energy would be transformed into other MHD waves depending on complex configurations of the magnetic fields (Tarbell et al. 1999; Sakai et al. 2000).

We would like to remark that other mechanisms also predict the production of longitudinal waves at the coronal height, although we have taken up the chromospheric reconnection model as a typical process in this paper. The generation of spicules in an open magnetic flux tube has been investigated by various authors (Hollweg, Jackson, & Galloway 1982; Hollweg 1992; Kudoh & Shibata 1999). They found that Alfvén waves excited by random motions in the photosphere (Ulrich 1996) effectively transport their energy to the corona. The nonlinear effect of torsional Alfvén waves produces longitudinal waves along the vertical flux tube in the corona. A sizable fraction of the initial energy of the transverse waves in the photosphere is converted to energy of longitudinal waves at the coronal height, and these waves become acoustic waves propagating upward. They could heat the surrounding plasma in the very same way as those triggered by the reconnection events above. Thus, acoustic waves are expected to be universally generated by various mechanisms in the corona far above the photosphere, and therefore, it is quite worth studying their role in coronal heating.

2.2. Dissipation of Acoustic Waves in the Corona

In this section, we describe our method of treating the outward propagation of acoustic waves, after they have been excited in the corona. We first estimate the distance the acoustic waves travel before forming shocks in a plane-parallel geometry. Then, we derive an equation dealing with the variation of N-wave amplitude in spherical geometry.
For simplicity’s sake, in the following discussions, we neglect the effects of magnetic field on the propagation of the waves. This simplification is valid when the circumstantial magnetic configuration is perpendicular above the flux tube generating the waves.

2.2.1. The Distance Acoustic Waves Travel before Forming Shocks

Any acoustic wave having a finite amplitude inevitably changes its shape, makes the wave front steepen, and eventually forms a shock front (e.g., Landau & Lifshitz 1959). The distance the acoustic waves travel before forming the shocks can be estimated for a given wavelength \( \lambda \) and initial amplitude \( \delta v_0 \). Consider acoustic waves propagating in the upward direction \( z \) in an isothermal atmosphere with a density structure of \( \rho = \rho_0 \exp(-z/H_p) \). Assuming waves with initially sinusoidal velocity profiles, \( \delta v = \delta v_0 \sin(2\pi Z/\lambda) \), a wave crest overtakes the preceding trough to form a shock front at

\[
z - z_0 = 2H_p \ln \left[ 1 + \frac{1}{4(\gamma + 1)H_p \delta v_0} \right]
\]

(SS72), where \( z_0 \) is the position at which the waves are created, \( \gamma = 5/3 \) is the ratio of specific heat, and \( c_s \) is the sound velocity.

In our calculations, waves of initial amplitude \( \delta v_0/c_s \) of 0.1–1 are considered (SS 3.1), and therefore, the second term in the logarithm in equation (1) is bounded by an upper limit,

\[
\frac{1}{4(\gamma + 1)H_p \delta v_0} \leq \frac{\lambda}{H_p} = \frac{(c_s^2)}{(\gamma g)} c_s \tau
\]

\[
\approx 0.25 \frac{\tau}{100 \text{s}} \left( \frac{c_s}{2 \times 10^7 \text{cm s}^{-1}} \right)^{-1} ,
\]

where \( g \) is the acceleration of gravity and we have used \( H_p \approx H_p \), which is satisfied in the corona where temperature varies slowly within a scale of \( \lambda (\sim 10^9 \text{cm}) \). Then, equation (1) can be expanded to first order as

\[
z - z_0 \approx \frac{\lambda}{2(\gamma + 1)} c_s \tau.
\]

The factor \( [1/2(\gamma + 1)](c_s/\delta v_0) \) is of the order of unity, indicating that after traveling one wavelength, the waves begin to dissipate energy.

2.2.2. Variation of N-Wave Amplitude

The initially sinusoidal acoustic waves are expected to be transported as N-waves after the formation of the shocks, provided that they are continuously generated from the lower corona. These N-waves propagate outward in the rarefied atmosphere and dissipate their wave energy in heating the corona. For the purpose of giving a reasonable estimate of the heating rate as a function of position, variations of amplitude \( \delta v_w \) of the transported N-waves are treated by weak-shock theory, following the formulation presented by SS72. In the discussions below, the propagation of the wave train is considered to be in a flow tube with a cross section of \( A \) to keep consistent with the model for the global corona presented in §2.3. The quantity \( A \) is a function of \( r \), the distance measured from the center of the Sun, and is modeled in §2.3 taking into account the nonradial expansion of the flow tube. Hereafter, all physical quantities are expressed only as functions of \( r \), unless explicitly declared otherwise.

An equation for the variation of wave amplitude normalized by ambient sound velocity, \( \alpha_w \equiv \delta v_w/c_s \), can be found from SS72,

\[
\frac{1}{\alpha_w} \frac{d\alpha_w}{dr} = \frac{1}{2} \left( -\frac{1}{p} \frac{dp}{dr} + \frac{1}{E_\lambda} \frac{dE_\lambda}{dr} - \frac{1}{\lambda} \frac{d\lambda}{dr} \right) ,
\]

where \( p \) is gas pressure and \( E_\lambda = \frac{1}{2} \rho (\delta v_w)^2 \lambda = \frac{1}{2} \rho \alpha_w^2 \lambda \) is the wave energy per wavelength \( \lambda \). The variation of the wave energy can be estimated from the entropy generation by the weak shock (SS72; §5, Mihalas & Mihalas 1984) as

\[
\mathbf{V} \cdot \mathbf{E}_\lambda = \frac{dE_\lambda}{dr} - \frac{E_\lambda}{A} \frac{dA}{dr} = -E_\lambda \frac{2(\gamma + 1)}{\lambda} \alpha_w
\]

\[
\approx -E_\lambda \frac{2(\gamma + 1)}{\lambda} \alpha_w ,
\]

where we have used the relation of \( \lambda = c_s(1 + [(\gamma + 1)/2]\alpha_w) \tau \approx c_s \tau \) by assuming \( \alpha_w < 1 \). In general, the period of waves traveling in different media remains constant, implying \( (1/\lambda)(d\lambda/dr) = (1/c_s)(dc_s/dr) \). Hence, equation (4) is reduced to

\[
\frac{d\alpha_w}{dr} = \frac{\alpha_w}{2} \left( -\frac{1}{p} \frac{dp}{dr} - \frac{2(\gamma + 1)\alpha_w}{c_s \tau} - \frac{1}{\lambda} \frac{d\lambda}{dr} - \frac{1}{c_s \tau} \frac{dc_s}{dr} \right) .
\]

This equation determines the variation of the wave amplitude in the solar corona according to the physical properties of the background coronal plasma. The first term on the right-hand side is positive in the solar atmosphere, the density of which decreases outwardly, and the second term of the entropy generation (heating) is negative. In the 1–2 \( R_\odot \) region, where the dissipation is important in heating the corona, these two terms always dominate the other terms, the third term arising from the geometrical expansion and the fourth from temperature variation. Using the \( \alpha_w \) determined above as well as the background physical quantities \( \rho \) and \( c_s \), the wave energy flux \( F_w \) (ergs cm\(^{-2}\) s\(^{-1}\)) is derived as

\[
F_w = \frac{1}{3} \rho c_s^3 (\alpha_w)^2 \left( 1 + (\frac{\gamma + 1}{2}) \alpha_w \right) ,
\]

if recalling that the wave crest moves at a speed of \( c_s(1 + [(\gamma + 1)/2]\alpha_w) \).

Since the model described above is simple, we instead have several limitations that should be taken into account with great care. The first obvious limitation is that the wave amplitude should satisfy the assumption of weak shock, \( \alpha_w < 1 \). Second, it does not take into account the effects of gravity. Third, it is constructed in a static medium, and therefore we cannot apply it to cases of moving media, such as N-waves in the solar wind where the flow velocity \( v \) exceeds the ambient \( c_s \). Of the three limitations, the third one is most easily overcome, because in all the cases we calculate in this paper, at least 99% of the initial wave energy dissipates within \( r < 1.3 R_\odot \), the region where \( v \ll c_s \) is fulfilled (Figs. 1 and 2). The second limitation also seems to have little effect if the wave period is small enough, since SS72 found that weak-shock theory gives reasonable estimates for waves of \( \tau < \frac{1}{2} (2\pi/\omega_w) \), where \( \omega_{wc} = \gamma g/2c_s \) is the acoustic cutoff frequency in a gravitationally stratified atmosphere, after comparing results of weak-shock theory with those of fully nonlinear calculations. As \( \omega_{wc} \approx \)}
(1/800 s)(c_s/2 \times 10^7 \text{ cm s}^{-1})^{-1} \text{ in the corona, weak-shock theory seems to be applicable to waves of } \tau < 2000 \text{ s. To check whether the first limitation is overcome, we need to calculate the } \alpha_w \text{ variation in the solar corona at first hand. Our results show that any waves we calculate give } \alpha_w < 0.5 \text{ in the entire corona (§ 3.1). Thus, weak-shock theory appears to be applicable to our model and to give a reasonable estimate of the heating rate as a function of } r.

2.3. Basic Equations

We here present basic equations to describe one-component coronal wind structure in a flow tube with a cross section of \( A \) under a steady state condition. Then, the equation of continuity becomes

\[ \rho v A = \text{const} \, . \] (8)

The equation of momentum conservation is

\[ \frac{dv}{dr} = -\frac{GM_\odot}{r^2} - \frac{1}{\rho \, dr} \frac{dp}{\rho c_s^2} \frac{1}{1 + [(\gamma + 1)/2]\alpha_w} \mathbf{V} \cdot \mathbf{F}_w \, . \] (9)

The pressure \( p \) is related to \( \rho \) and the temperature \( T \) by the equation of state for an ideal gas,

\[ p = \rho \frac{k_B}{m_{\text{H}} \mu} T \, , \] (10)

where \( m_{\text{H}} \) is the hydrogen mass, \( k_B \) is the Boltzmann constant, and \( \mu \) is the mean atomic weight of the particles in units of \( m_{\text{H}} \). The third term in equation (9) represents wave momentum deposition, neglecting the effects of reflection and refraction (Rosner & Vaiana 1977). The total energy

\[ \Gamma = 7.8 \times 10^5 \text{ ergs cm}^{-2} \text{s}^{-1} \, . \]
The equation is obtained as
\[
\frac{Dx}{C_20} + \frac{v_1^2 + \frac{1}{v}}{C_{13}} + \frac{v_2}{C_{13}} + \frac{1}{v} = \frac{k_B m_H T}{C_0 G M} + \frac{F_w + F_c}{C_{21}} + q_r = 0,
\]
where we adopt the classical form of the conductive flux for ionized gas,
\[
F_c = -\kappa \frac{dT}{dr} = -\kappa_0 T^{5/2} \frac{dT}{dr},
\]
with $$\kappa_0 = 1.0 \times 10^{-6}$$ in cgs units (Allen 1973), and the radiative cooling term $$q_r$$ (ergs cm \(^{-3}\) s \(^{-1}\)) is derived from the tabulated radiative loss function $$\Lambda$$ (ergs cm \(^{-3}\) s \(^{-1}\)) (Landini & Monsignori-Fossi 1990) for an optically thin plasma as $$q_r = n_e n_p \Lambda$$, where $$n_e$$ and $$n_p$$ are the electron and H\(^+\) density, respectively. They are calculated by solving the ionization of H and He (we ignore heavy elements) under the LTE condition. The term $$\mathbf{v} \cdot \mathbf{F}_w (\leq 0)$$ in equation (11) indicates the heating by the dissipation of the waves, which is explicitly written as
\[
\mathbf{v} \cdot \mathbf{F}_w = \left\{ \begin{array}{ll}
0, & r \leq r_d, \\
\rho v \left[ \frac{1}{3} \alpha_w \left( 1 + \frac{\gamma + 1}{2} \alpha_w \right) \frac{d}{dr} \left( \frac{c_s^2}{v} \right) \right] + \left( \frac{2}{3} \alpha_w + \frac{\gamma + 1}{2} \alpha_w^2 \right) \frac{d}{dr} \left( \frac{\alpha_w}{v} \right), & r > r_d,
\end{array}\right.
\]
where $$r_d$$ is the position at which the waves start to dissipate by shaping N-waves and the height $$h_d = r_d - R_\odot$$ from the photosphere corresponds to the sum of the height of the magnetic flux tube ($$\sim 10^9$$ cm: the generation point of the waves) and the distance the acoustic waves travel before forming the shock fronts ($$\sim \lambda \sim 10^9$$ cm; eq. [3]). We adopt $$h_d = 2 \times 10^9$$ cm as a standard value.

To take into account the nonradial expansion of the flow tube due to configurations of the magnetic field, the cross-
sectional area $A$ is modeled as
\[
A = r^2 f_{\text{max}} e^{(r-\sigma)/\alpha} + f_1 e^{(r-\sigma)/\alpha} + 1,
\]  
(14)
where
\[
f_1 = 1 - (f_{\text{max}} - 1) e^{(1-\tau)/\alpha}
\]
(Kopp & Orrall 1976; Withbroe 1988). The cross section expands from unity to $f_{\text{max}}$ most drastically between $r = r_1 - \sigma$ and $r_1 + \sigma$. Of the three input parameters, $f_{\text{max}}$ is the most important in determining the solar wind structure. In this paper, we consider cases for $f_{\text{max}} = 5$ and $f_{\text{max}} = 1$ (pure radial expansion). As for the other two parameters, we employ $r_1 = 1.25 R$ and $\sigma = 0.1 R$, the same values adopted in a model for a “quiet corona” in Withbroe (1988).

2.4. Boundary Conditions and Computational Method

Now we would like to explain the practical aspects of our method of constructing wind structure with respect to various input properties of the waves. In order to solve both the heating of the corona (energy transfer) and the formation of the solar wind (momentum transfer) consistently, our calculation is performed in a broad region from the inner boundary of the upper chromosphere where the temperature $T_{\text{ch}} = 10^4$ K, at $r_{\text{ch}} = R_s + h_{\text{ch}}$, which is located at $h_{\text{ch}} = 2 \times 10^8$ cm (2000 km) above the photosphere (Allen 1973), to an arbitrary outer boundary at $r_{\text{out}} = 300 R$.

We set four boundary conditions to construct a unique solution for a given input wave energy flux $F_{\text{w},0}$ as follows:
\[
F_w(r_{\text{ch}}) = F_w(r_s) = F_{\text{w},0},
\]
(15)
\[
T(r_{\text{ch}}) = T_{\text{ch}},
\]
(16)
\[
|F_s(r_{\text{ch}})| \approx |F_{s,\text{max}}|,
\]
(17)
\[
\mathbf{v} \cdot \mathbf{F}_s(r_{\text{out}}) = 0,
\]
(18)
where $F_{s,\text{max}}$ in equation (17) is the maximum value of the downward conductive flux in the inner corona. The first condition denotes that the wave energy flux must agree with the given value when the waves start to dissipate. The second condition is also straightforward: the temperature has to coincide with the fixed value at the inner boundary. The third condition is the requirement that the downward thermal conductive flux should become sufficiently small in the upper chromosphere ($T = 10^4$ K), diminishing from its enormous value at the coronal base ($T \approx 10^6$ K). Practically, we continue calculations iteratively until $F_s(r_{\text{ch}})/F_{s,\text{max}} < 1\%$ is satisfied. The fourth condition corresponds to the ordinary requirement that no heat is conducted inward from infinity (Sandbæk & Leer 1994). Note that thanks to the third condition, the coronal base density, which is poorly determined from observations, does not have to be used as a boundary condition. As a result, the number of free parameters to be set in advance is reduced (Hammer 1982a, 1982b; Withbroe 1988). The density at the coronal base or the transition region (TR) is calculated as an output; a larger input $F_{\text{w},0}$ increases the downward $F_s$ in the lower corona, demanding a larger density in the coronal base and TR to enhance radiative cooling to balance with the increased conductive heating.

For numerical integration of the momentum equation (9) and the energy equation (11), we respectively use $v$ and an isothermal sound velocity $a$, defined as
\[
a^2 = \frac{c^2}{\gamma} = \frac{p}{\rho} = \frac{k_B}{m_\text{H}u} T.
\]
(19)
To carry out the integration, the equations shown in the previous section need to be transformed into useful forms. First, an expression for the velocity gradient can be written from equations (8)–(10):
\[
\frac{dv}{dr} = -\frac{GM_\odot}{r^2} + \frac{(a^2/A)(dA/dr) - da^2/dr}{v - a^2/v}.
\]
(20)
Second, an expression for the gradient of $a^2$ is derived from an integrated form of equation (11),
\[
\frac{da^2}{dr} = \frac{\rho v}{\kappa_0 a^2} \left( \frac{k_B}{m_\text{H}u} \right)^{7/2} \left( \frac{v^2}{2} + \frac{\gamma}{\gamma - 1} a^2 \right)
\]
\[- \frac{GM_\odot}{r} + \frac{F_w}{\rho v} + \int_{r_0}^{r} dr' q_s/v - E_{\text{tot}} \right),
\]
(21)
where $r_0$ is a certain reference point of integration that is set below. The quantity $E_{\text{tot}}$ is mathematically an integral constant, which must be conserved in the entire region of the calculation.

Only transonic solutions are allowed for the flow speed $v$. The integration of the three differential equations (20), (21), and (6) is carried out simultaneously from the sonic point, $r = r_s$, to both outward and inward directions by the fourth-order Runge-Kutta method, also deriving the density from equation (8) at each integration. We have used variable sizes for the grid of the integration, setting a smaller mesh size in regions where physical values change rapidly. For instance, in the TR, we set a mesh as small as $10^3$ cm (0.01 km) for one grid. To start the integration, we have to set nine variables, $v_s$, $(dv/dr)_s$, $a^2_s$, $(da^2/dr)_s$, $\alpha_{w,s}$, $(d\alpha_{w}/dr)_s$, $\rho_s$, $E_{\text{tot}}$, and $r_s$ (the subscript “s” denotes the sonic point). We can determine $v_s$, $(dv/dr)_s$, and $r_s$ for a given $a^2_s$ and $(da^2/dr)_s$, by the condition that both the numerator and the denominator of equation (20) are zero at $r_s$ (Parker 1958). The quantity $(d\alpha_{w}/dr)_s$ is also derived from equation (6) for a given $\alpha_{w,s}$. Setting the reference point $r_0 = r_s$ for the integration of the radiative cooling function, we obtain $E_{\text{tot}}$ as
\[
E_{\text{tot}} = \left[ \frac{1}{2} \frac{v^2}{v} + \frac{\gamma}{\gamma - 1} a^2 - \frac{GM_\odot}{r} + \frac{F_w}{\rho v} \right. \left. - \left( \frac{m_\text{H}u}{k_B} \right)^{7/2} \frac{1}{\rho v} \right] \frac{2}{\kappa_0 a^2} \left( \frac{da^2}{dr} \right),
\]
(22)
where all the variables are evaluated at the sonic point. Now we have four variables, $a^2_s$, $(da^2/dr)_s$, $\alpha_{w,s}$, and $\rho_s$, remaining to be regulated by the four boundary conditions of equations (15)–(18). Concrete procedures for finding a unique solution for a given $F_{\text{w},0}$ are described below.

1. One makes an initial guess for $a^2_s$, $(da^2/dr)_s$, $\alpha_{w,s}$, and $\rho_s$.

2. The integration is performed in the outward direction from $r_s$. Leaving $a^2_s$, $\alpha_{w,s}$, and $\rho_s$ unchanged, $(da^2/dr)_s$ is determined by carrying out the integration iteratively to satisfy the outer boundary condition of equation (18).

3. The integration is performed iteratively in the inward direction, improving $a^2_s$, $\alpha_{w,s}$, and $\rho_s$ for a fixed $(da^2/dr)_s$, until they satisfy the three inner boundary conditions, equa-
tions (15)–(17). Physically, the condition of the wave flux (eq. [15]) regulates \(\alpha_{\text{nu}}\), that of the temperature (eq. [16]) regulates \(\alpha_{\text{tu}}\), and that of the conductive flux (eq. [17]) regulates \(\rho_{\text{e}}\). These relations guide the improvement of the respective initial guesses, although they are not independently approved. Unless the above inner boundary conditions are satisfied simultaneously, one returns to step 2, preparing a new set of \(a^w_{\text{d}}(da^w/dr)\), \(\alpha_{\text{nu}}\), and \(\rho_{\text{e}}\).

4. One can finally find a unique solution by iterating procedures 2 and 3.

3. RESULTS

3.1. Theoretical Interpretation of the Resulting Coronal Structures

Figure 1 demonstrates our resulting coronal wind structures employing the same input \(F_{w,0} = 7.8 \times 10^5 \text{ ergs cm}^{-2} \text{ s}^{-1}\). But different sets of wave periods and nonradial expansion factors, \(\tau(s), f_{\text{max}}\) = [300, 5], [60, 5], [300, 1]. Figure 2 compares the wind structures adopting three different \(F_{w,0} = (3.2, 5.9, 10) \times 10^5 \text{ ergs cm}^{-2} \text{ s}^{-1}\) for identical inputs of \(\tau = 120 \text{ s}\) and \(f_{\text{max}} = 5\). The distribution of temperature, flow velocity, and electron density from the inner boundary to 1 AU = (215 \(R_\odot\)) are displayed from top to bottom on the left-hand side of both figures. On the right-hand side, we present the variation of the wave amplitude \(\alpha_{\text{nu}}\), the dissipation length \(l_{\text{d}}\) of the N-waves, which is defined as \(l_{\text{d}} = F_w/\mathbf{V} \cdot \mathbf{F}_w\), and the heating per unit mass \((1/\rho_0)\mathbf{V} \cdot \mathbf{F}_w\), respectively. In Table 1 we tabulate several resulting properties of the corona and solar wind as well as the input parameters.

To begin with, we would like to emphasize that the corona heats up to more than \(10^6 \text{ K}\) in every case, because acoustic waves generated in the corona are able to heat the surrounding gas directly, unlike acoustic waves produced in the photosphere. However, the N-waves are rapidly damped so that the heating occurs only in the inner region, as seen in the bottom right panels of Figures 1 and 2. Then, the location of the maximum temperature \(T_{\text{max}}\) is quite close to the surface. In the region \(\lesssim 1.5 R_\odot\), heat is input only by outward thermal conduction, and the flow is mostly accelerated by thermal pressure. As a result, the speed of the solar wind at 1 AU is \(\lesssim 300 \text{ km s}^{-1}\), which is slightly slower than the actual low-speed wind (300–450 km s\(^{-1}\)).

The top right panel of Figure 2 interestingly illustrates that the distributions of \(\alpha_{\text{nu}}\) are almost identical in spite of very different inputs of \(F_{w,0}\). Particularly, the initial N-wave amplitudes \(\alpha_{\text{nu}}(r_{\text{d}})\) at \(r_{\text{d}}\) are within the range between 0.48 and 0.49. This is because \(F_{w,0} = -\rho_0 a^w_{\text{d}} C_2 c\), (eq. [7]) mostly owes its variation to changes of ambient pressure (see § 3.2.1). Moreover, the top right panel of Figure 1 also indicates that the initial \(\alpha_{\text{nu}}(0.5)\) is almost independent of \(\tau\) and \(f_{\text{max}}\). We have found that \(0.45 < \alpha_{\text{nu}}(r_{\text{d}}) < 0.52\) within the parameter regions of \(F_{w,0} = (1–20) \times 10^5 \text{ ergs cm}^{-2} \text{ s}^{-1}, \tau = 60–300 \text{ s}\), and \(f_{\text{max}} = 1–5\). This proves that \(\alpha_{\text{nu}} < 1\) is fulfilled in the entire region, which justifies the assumption of weak shock. The middle right panels of Figures 1 and 2 show that the dissipation length is a drastically varying function of \(r\). This implies that the assumption of constant dissipation length usually taken in previous models for the global corona (Withbroe 1988; Sandbæk & Leer 1994) is very poor for our N-wave process.

In the following discussions, we examine the dependence of the wind structures on the respective input parameters. First, we examine the dependence on wave periods. As illustrated in Figure 1, N-waves with a smaller \(\tau\) dissipate more quickly, and the heating occurs in a thinner region close to the surface. This simply leads to the deposition of wave energy in a denser region. Since the radiative loss \(q_r\) (ergs \(\text{ cm}^{-2} \text{ s}^{-1}\)) is in proportion to \(\rho^2\) for optically thin plasma, a greater fraction of energy supplied in the denser region goes into radiative escape. Consequently, a smaller amount of energy remains to heat the corona and accelerate the flow. A case adopting a smaller \(\tau = 60 \text{ s}\) gives a lower temperature in the corona and therefore a smaller pressure scale height and a more rapid decrease of density, as shown in Figure 1. A lower temperature also makes the sonic point more distant from the solar surface, and then the mass flux of the solar wind becomes much smaller than that expected from the \(\tau = 300 \text{ s}\) case (Table 1).

Second, we study the effects of areal expansion of the flow tube. Comparing the results of adopting the same \(F_{w,0} = 7.8 \times 10^5 \text{ ergs cm}^{-2} \text{ s}^{-1}\) and \(\tau = 300 \text{ s}\) but a different \(f_{\text{max}} = 1\) and 5 in Figure 1, one can notice a significant change of density structure. The model considering the non-radial expansion gives a more drastic decrease of density as a function of \(r\) in spite of a similar initial density at the inner boundary. The temperature in the inner corona is also lower in that model, since a larger fraction of the input energy is lost adiabatically due to geometrical expansion of the flow tube. On the other hand, the decrease of temperature is slower and the temperature in the outer region (\(\gtrsim 8 R_\odot\)) is higher. This is because the lower density reduces both radiative cooling and adiabatic loss. The higher temperature also leads to larger acceleration of the flow there, giving a larger speed to the solar wind in the outer corona.

Third, we investigate the dependence on input wave energy flux. Models employing a larger input \(F_{w,0}\) give a larger density in the inner corona (Fig. 2), because radiative cooling should be raised to offset the increased heating. More accurately, a larger input \(F_{w,0}\) results in a larger downward conductive flux from the corona to the chromosphere, which needs larger radiative cooling to balance the enhanced conductive heating. A larger \(F_{w,0}\) also leads to a higher temperature in the inner corona. However, a more rapid decrease of temperature occurs because of the enhanced radiative escape. A model employing \(F_{w,0} = 10 \times 10^5 \text{ ergs cm}^{-2} \text{ s}^{-1}\) yields a lower temperature in the

| Input Parameters and Output Wind Properties of Each Model |
|----------------------------------|--------|--------|--------|--------|
| \((F_{w,0}, \tau, f_{\text{max}})\) | \(p_{\text{TR}}\) | \(T_{\text{max}}\) | \(r_{\text{max}}\) | \(\eta_{\text{p.v}}\) | \(v_{\text{I.AU}}\) |
| \((7.8, 60, 5)\) | 0.28 | 1.43 | 1.04 | \(3.8 \times 10^6\) | 222 |
| \((7.8, 300, 5)\) | 0.24 | 1.50 | 1.08 | \(7.5 \times 10^5\) | 251 |
| \((7.8, 300, 1)\) | 0.29 | 1.71 | 1.12 | \(3.8 \times 10^6\) | 222 |
| \((3.2, 120, 5)\) | 0.13 | 1.16 | 1.05 | \(2.5 \times 10^6\) | 230 |
| \((5.9, 120, 5)\) | 0.21 | 1.35 | 1.05 | \(1.3 \times 10^7\) | 240 |
| \((10, 120, 5)\) | 0.32 | 1.52 | 1.05 | \(1.3 \times 10^7\) | 230 |

Note.—The quantity \(F_{w,0}\) is in \(10^5 \text{ ergs cm}^{-2} \text{ s}^{-1}\). \(\tau\) is in seconds, \(p_{\text{TR}}\) is the pressure in dyn cm\(^{-2}\) at the TR where \(T = 10^6 \text{ K}\), \(T_{\text{max}}\) is the peak coronal temperature in \(10^6 \text{ K}\), \(r_{\text{max}}\) is the location of \(T_{\text{max}}\) in \(R_\odot\), \(\eta_{\text{p.v}}\) at 1 AU, and \(v_{\text{I.AU}}\) is the flow velocity in \(\text{km s}^{-1}\) at 1 AU.
region of $r \geq 1.5 R_0$ than that adopting $F_{w,0} = 5.9 \times 10^5$ ergs cm$^{-2}$ s$^{-1}$. Therefore, a larger input of $F_{w,0}$ does not simply cause a larger acceleration of the flow, and the relations of $F_{w,0}$($n_p v_1$)$_{1AU}$ and $F_{w,0}$($v_1$)$_{1AU}$ do not show simple positive correlations (Table 1; see also § 3.2.2).

### 3.2. Some Characteristic Properties of the Resulting Corona

#### 3.2.1. Peak Temperature and Pressure at the TR

We would like to inspect the relation of the peak coronal temperature $T_{\max}$ and pressure $p_{\max}$ at the TR to the input model parameters. In Figure 3 we present $T_{\max}$ as a function of input $F_{w,0}$ for different $\tau$ and $f_{\max}$. The figure shows that $T_{\max} > 10^6$ K is accomplished for $F_{w,0} \geq 2 \times 10^7$ ergs cm$^{-2}$ s$^{-1}$, even if one chooses waves with a short period ($\tau = 60$ s) and a large expansion factor ($f_{\max} = 5$). The quantity $T_{\max}$ is a monotonically increasing function of $F_{w,0}$ for each set of ($\tau$, $f_{\max}$), and the relations can be fitted by a power law as

$$T_{\max} \propto (F_{w,0})^k, \quad \text{where } k = 0.23-0.26. \quad (23)$$

Figure 4 displays the relation between $F_{w,0}$ and $p_{\max}$. The quantity $p_{\max}$ is evaluated at $T = 10^5$ K, whereas $p_{\max}$ is almost a constant through the TR from the upper chromosphere ($T \approx 10^4$ K; inner boundary) to the coronal base ($T \approx 5 \times 10^5$ K) because of its geometrically thin configuration. The figure indicates that $p_{\max}$ weakly depends on $\tau$ and $f_{\max}$ and has a positive correlation with $F_{w,0}$:

$$p_{\max} \propto (F_{w,0})^l, \quad \text{where } l = 0.75-0.79. \quad (24)$$

A rise of $F_{w,0}$ leads directly to an increase in downward conductive flux, which demands a higher density (or pressure) in the lower corona and the TR to raise radiative cooling to balance the enhanced conductive heating. Alternatively, it could be interpreted that the injection of more energy into the corona can heat more deeply toward the higher density chromosphere by downward thermal conduction.

The combination of the above two relations of equations (23) and (24) roughly gives

$$T_{\max} \propto (p_{\max})^{1/3},$$

for a given $\tau$ and $f_{\max}$, which reminds us of the famous RTV scaling law (Rosner, Tucker, & Vaiana 1978), $T_{\max} \approx 1400(p L_h)^{1/3}$, for closed magnetic loops, where $p_l$ (dyn cm$^{-2}$) is the loop pressure and $L_h$ (cm) is the loop height. Taking $r_{\max}$ instead of $L_h$ for our model, we can actually derive a relation between $T_{\max}$ and $p_{\max}$ from our results as

$$T_{\max} \approx 3000(p_{\max} r_{\max})^{-0.30},$$

showing a form analogous to that of the original RTV law. This is because we consider the same energy balance among thermal conduction, radiative cooling, and heating, with the same boundary condition that the conductive flux should become almost zero at the base (eq. [17]), although the configurations are quite different (a closed loop for the RTV law and an open flow tube for ours). The slight discrepancies of the prefactor and power-law index are caused by the fact that the RTV law was derived on the assumption of spatially uniform heating along the loop, while our heating function is determined by equation (13), which is not uniform at all.

#### 3.2.2. Mass Flux and Coronal Energy Loss

In Figure 5 we show the anticipated proton flux $(n_p v_1)_{1AU}$ at 1 AU as a function of $F_{w,0}$ for different $\tau$ and $f_{\max}$, with
observational constraints, \( (n_r v)_1 \text{AU} = (3.8 \pm 1.5) \times 10^8 \text{ cm}^{-2} \text{ s}^{-1} \) (shaded region), compiled by Withbroe (1988) as the empirical value for the “quiet corona” that is supposed to correspond to the mid- to low-latitude region generating the low-speed wind. Larger \( r \) waves give a greater \( (n_r v)_1 \text{AU} \) owing to the effective transport of dissipated energy to the solar wind flow by avoiding radiative escape. Introduction of the nonradial expansion of the flow tube reduces \( (n_r v)_1 \text{AU} \), because the input energy per unit flow tube normalized at 1 AU decreases for increasing \( f_{\text{max}} \) even though one inputs identical wave energy flux at the inner corona. Therefore, the larger areal expansion straightforwardly reduces the mass flux of the solar wind. As to the dependence on \( F_{w,0} \), \( (n_r v)_1 \text{AU} \) has an upper limit for a given \( r \) and \( f_{\text{max}}. \) (Even models of \( r = 300 \text{ s} \) and \( f_{\text{max}} = 1 \) are supposed to have an upper limit in a \( F_{w,0} > 2 \times 10^6 \text{ ergs cm}^{-2} \text{ s}^{-1} \) area.) Figure 5 shows that only one case employing \((r, f_{\text{max}}) = (300, 1)\) can reproduce the observed proton flux of the slow wind. Unfortunately, more realistic cases considering the nonradial expansion of \( f_{\text{max}} = 5 \) cannot explain the observations, which implies that other mechanisms of heating (and acceleration) necessarily work cooperatively.

We would like to study in detail why \( (n_r v)_1 \text{AU} \) has an upper limit for fixed \( r \) and \( f_{\text{max}}. \) Figure 6 displays the ratios of the three main types of energy loss: downward thermal conduction from the coronal base, radiative escape in the corona, and total amount of energy converted to the flow (mass loss) at \( r_{\text{out}} \), respectively normalized by \( F_{w,0} \) for the cases adopting \( r = 120 \text{ s} \) and \( f_{\text{max}} = 5 \). The values for downward thermal conduction in Figure 6 are taken from the maximum conductive flux \( F_{c,\text{max}} \) in the lower corona (see Fig. 7), and the values for the radiative escape are derived by the integration of the radiative cooling function from points with \( F_{c,\text{max}} \) to \( r_{\text{out}}. \) The region not labeled between “radiative escape” and “flow” denotes the energy carried out of \( r_{\text{out}} \) by thermal conduction. According to Figure 6, the main source of coronal energy loss is downward thermal conduction within our range of \( F_{w,0} \) (strictly speaking, most of the conducted energy finally radiates away), whereas radiative escape comes to play a significant role for larger \( F_{w,0} \), since the density in the corona becomes higher, being subject to equation (24). Although the ratio of energy transferred to the flow is as small as \( \lesssim 3\% \) throughout the range, it has a bimodal tendency. With \( F_{w,0} < 6 \times 10^6 \text{ ergs cm}^{-2} \text{ s}^{-1} \), it increases, which implies that the larger \( F_{w,0} \) is, the more effectively the energy is transferred to the flow. However, with \( F_{w,0} > 6 \times 10^6 \text{ ergs cm}^{-2} \text{ s}^{-1} \), it decreases because of the abrupt dominance of radiative cooling. As a result, the predicted \( (n_r v)_1 \text{AU} \) increases rapidly with increasing \( F_{w,0} \) at first and eventually decreases, as seen in Figure 5.

To examine these differences in terms of energy transfer, we show the variations of the energy flux of four components,

\[
\text{wave: } f_w = \frac{A(r)}{A(r_{\text{ch}})} F_w, \quad (25)
\]

\[
\text{conduction: } f_c = \frac{A(r)}{A(r_{\text{ch}})} F_c, \quad (26)
\]

flow: \[
f_f = \rho \nu \frac{A(r)}{A(r_{\text{ch}})} \left[ \left( \frac{1}{r^2} + \frac{\gamma}{\gamma - 1} a^2 - \frac{GM_\odot}{r} \right) \right] \left( \frac{1}{r^2} + \frac{\gamma}{\gamma - 1} a^2 - \frac{GM_\odot}{r} \right) \right] \right], \quad (27)
\]

\[
\text{radiation: } f_r = \rho \nu \frac{A(r)}{A(r_{\text{ch}})} \int_{r_{\text{ch}}}^{r_{\text{out}}} dr \frac{q_r}{\rho v}, \quad (28)
\]

per flow tube with a cross section of \( A = 1 \text{ cm}^2 \) at the inner boundary. The “flow” term of equation (27) contains three ingredients, kinetic energy of the solar wind, enthalpy, and gravitational energy. Note that except for the wave energy flux, the zero point of the energy is taken here at the inner boundary, to clarify what fraction of the input energy flux is transferred to the other components. In the left panel of Figure 7, the status of the energy transfer in the inner corona is displayed in a linear scale for both the \( x \)- and \( y \)-axis, and in the right panel, that in the broader region is shown in log scale. The ratio of the downward thermal conduction \( f_f/F_{w,0} \) is smaller for larger \( F_{w,0} \), although the absolute value of \( F_c \) is increasing along with \( F_{w,0}. \) At the TR, most of the heat flux by the downward conduction finally escapes as radiation, except for a tiny fraction transferred to the enthalpy to be used to heat the TR. This leads to a smaller ratio of the radiative loss \( F_r/F_{w,0} \) for larger \( F_{w,0} \) in the inner region of less than 1.2 \( R_\odot \). However, the contribution from the radiation continues in a much more distant region in the model employing the largest \( F_{w,0} \), and \( f_r/F_{w,0} \) finally outdoes the other two cases. Consequently, the ratio of the energy transferred to the solar wind becomes smaller than for the model adopting smaller \( F_{w,0} (=5.9 \times 10^8 \text{ ergs cm}^{-2} \text{ s}^{-1}) \), as seen in the right panel. It can be concluded that if \( F_{w,0} \) is larger than a certain threshold, an increase of the input \( F_{w,0} \) does not lead to effective heating of the corona to accelerate the solar wind but results in deposition of the wave energy in the high-density region to be wasted as radiative escape.

3.2.3. Wave Amplitude in the Inner Corona

Ultraviolet and X-ray emission lines of the corona show nonthermal broadenings (Hassler et al. 1990; Erdélyi et al. 1998), which are inferred to originate from wave motions.
We investigate whether the obtained wave amplitudes are consistent with these observations. In Table 2 we show our results for the wave rms velocities in the inner corona, which are calculated as

\[
\langle \delta_w \rangle = \frac{1}{\sqrt{3}} \delta_w
\]

for the N-shaped waves. As the waves propagate upward from the location of \( T = 5 \times 10^5 \) K to that of \( T = 10^6 \) K, the wave amplitude increases as the ambient density decreases. Our models give very similar results for different \( F_{w,0} \) (and \( \tau \) and \( f_{\text{max}} \)), since a change of \( F_{w,0} \) \( (\sim \rho \langle \delta_w \rangle^2 c_s) \) mostly reflects a variation of the density (Fig. 2; § 3.1). The tabulated results are marginally consistent with the nonthermal velocities of 20–40 km s\(^{-1}\) obtained from observations of inner coronal lines in the solar disk and limbs (Erődy et al. 1998), if we assume that the observed nonthermal components totally consist of waves of the acoustic mode. However, this assumption may be too extreme, because other modes of waves actually exist in the real solar corona.

**3.3. Comparison of Coronal Wind Structure with the Observed Low-Speed Wind**

In this section, we study the feasibility of the process of acoustic waves in coronal heating by comparing our results with recent observations. Focusing on coronal heating and wind acceleration in the region where the low-speed wind is formed, we take observational data of the coronal streamer in the mid- to low-latitude region.

### 3.3.1. Density Distribution

In Figure 8 we display our results of electron density from the observation of the streamer (Parenti et al. 2000; Hayes, Vourlidas, & Howard 2001). Our models adopt an identical \( F_{w,0} = 7.8 \times 10^5 \) erg cm\(^{-2}\) s\(^{-1}\) and three sets of \( \langle \tau(s) \rangle, f_{\text{max}} = [300, 5], [60, 5], [300, 1] \) (same as Fig. 1). As for the observation, we show results derived from the line ratio of Si\( \text{IX} \) by the Solar Heliospheric Observatory Coronal Diagnostic Spectrometer (SOHO/CDS) for the region of 1.02–1.19 \( R_\odot \), the ratio of radiative and collisional intensities of the O\( \text{VI} \) line by the Ultraviolet Coronagraph Spectrometer (SOHO/UVCS) for the region of 1.58–1.6 \( R_\odot \).
(Parenti et al. 2000), and the total brightness obtained from the Large Angle and Spectrometric Coronagraph (SOHO/LASCO) for the region $1.5 \ R_\odot \leq r \leq 6 \ R_\odot$ (Hayes et al. 2001). With respect to the CDS and UVCS data, results for both equatorial and midlatitude streamer regions are displayed.

Our results, adopting the same $F_{\text{w},0}$, show almost identical density in the very inner part, being independent of $\tau$ and $f_{\text{max}}$ (Fig. 4). The figure shows that the adopted value $F_{\text{w},0} = 7.8 \times 10^5 \text{ ergs cm}^{-2} \text{ s}^{-1}$ reproduces the CDS observation of the midlatitude streamer well. To fit the data of the equatorial streamer, a slightly smaller input of $F_{\text{w},0}$ is favored. In our model, $\tau$ and $f_{\text{max}}$ control the slope of the density distribution. The quantity $\tau$ influences the temperature in the intermediate region of $2 \ R_\odot \leq r \leq 10 \ R_\odot$ (Fig. 1) and hence regulates the density scale height (or decreasing slope of density) in that region. The data based on SOHO/LASCO indicate that models considering larger $\tau$ ($\approx 300$ s) are more likely. The quantity $f_{\text{max}}$ determines the density decrease in the region of $1-2 \ R_\odot$. The data of CDS ($\leq 1.2 \ R_\odot$) and UVCS ($\sim 1.6 \ R_\odot$) exhibit a drastic decrease of density, which indicates that nonradial expansion is desired. Adjustment of the other parameters of the flow tube geometry ($r_1$ and $\sigma$; eq. [14]) would give a still better fit. Although we do not search further for the best parameter set to fit observation, the figure indicates that our model could reproduce the observed density profile by the choice of the appropriate parameters $F_{\text{w},0} \approx (5-8) \times 10^5 \text{ ergs cm}^{-2} \text{ s}^{-1}$, $\tau \approx 300$ s, and $f_{\text{max}} \approx 5$.

### 3.3.2. Temperature Distribution

Figure 9 compares our results of temperature distribution with observations in the streamer region. Our models employ the same parameter sets as in Figure 8, and the data are the electron temperatures obtained from observations of the line ratio of Fe xvi/Fe x by the CDS and UVCS (Parenti et al. 2000). Although the observed data are electron temperatures, they are supposed to represent the plasma temperature because electron-ion equilibrium is attained in the dense streamer region of $r \leq 2 \ R_\odot$ (Raymond et al. 1998). Therefore, it is reasonable to compare them with the results of our model of a one-fluid coronal plasma. The case with $(F_{\text{w},0}, \tau, f_{\text{max}}) = (7.8 \times 10^5, 300, 5)$ gives a reasonable peak temperature of $\approx 1.5 \times 10^6$ K. However, none of our models can reproduce the observed location of $T_{\text{max}}$. While the location is observationally inferred to be between 1.2 and 1.6 $R_\odot$, all of our models give $r_{T_{\text{max}}} < 1.2 \ R_\odot$. This is because the dissipation length of the N-waves is essentially short, even though one considers long-period waves that are generated in the corona. In summary, acoustic waves excited in the corona can certainly heat the surrounding plasma to $T > 10^6$ K; however, they cannot maintain the high temperature out to a sufficiently distant region by themselves. Therefore, cooperation with other heating sources with larger dissipation lengths are necessary to explain the observed solar corona.

#### 3.3.3. Velocity Distribution

In Figure 10 we show the results of the velocity distribution of the solar wind, with observational results in the low-latitude streamer (shaded region). The observational data are from Sheeley et al. (1997), who determined the velocity profile between 2 and $30 \ R_\odot$ from measurements of about 65 moving objects in the streamer belt. They used two different techniques in deriving the results, whereas the shaded area displayed in Figure 10 is based on the straight-line fit method (Fig. 6 [top]; Sheeley et al. 1997; the shaded region is traced from that figure). Figure 10 indicates that the resulting velocity of our models is lower than the observed data in the whole region. The observation exhibits rapid acceleration at $3-5 \ R_\odot$, while our results show gradual acceleration by thermal gas pressure, since the N-waves are damped in less than 1.5 $R_\odot$ and wave pressure cannot contribute to the acceleration of the wind flow, as shown in Figures 1 and 2. Other mechanisms are also required for acceleration of the low-speed wind.
4. SUMMARY AND DISCUSSIONS

We have investigated the process of acoustic waves generated in the corona as a heating source, especially in the mid-to low-latitude corona where low-speed winds come from. We have found that acoustic waves with \( \tau \geq 60 \text{ s} \) and \( F_{w,0} \geq 2 \times 10^5 \text{ ergs cm}^{-2} \text{ s}^{-1} \) could certainly heat up the ambient plasma to \( T \geq 10^6 \text{ K} \) by the dissipation of N-waves, even in a flow tube with an areal expansion of \( f_{max} = 5 \), by balancing with the losses of radiative cooling, downward thermal conduction, and adiabatic loss by the solar wind. Because of its dissipative character, the dissipation of N-waves effectively works in the inner corona and reproduces the density profile observed in the streamer region. However, it cannot contribute to the heating of the outer corona, since most of the wave energy is damped within a region of a few tenths of the solar radius, even if one considers waves with a long period of \( \tau = 300 \text{ s} \). As a result, it is impossible to explain the observed temperature profile and flow velocity of the low-speed wind by only that process. Therefore, other mechanisms with a larger dissipation length should play a cooperative role in the coronal heating and the acceleration of the low-speed wind.

Small-scale reconnection events and spicules also predict the generation of fast shock waves (Lee & Wu 2000; Lee 2001; Hollweg 1982), although we have concentrated on the role of the N-waves (slow shocks along the magnetic field line) excited by those events in this paper. The dissipation length of the fast shocks must be larger than that of N-waves with an identical period, since their phase speed is \( \gtrsim v_A \), which is much larger than that of the N-waves (~1.5 \( v_A \)) for low-\( \beta \) coronal plasma. Hence, cooperation with fast shocks would give a more distant location of \( T_{max} \) to match the observed temperature profile (Fig. 9) and would let the acceleration of the wind flow continue in the outer corona, as suggested by observation of the streamer belt (Fig. 10).

Finally, we had better remark on the issue of anisotropic and selective heating of ions. According to recent observations (Strachan et al. 2002), O vi ions in the streamer have a high perpendicular kinetic temperature, although this is not as extreme as that observed in the high-speed wind. Our results show that the deposition of energy and momentum from the N-waves to the ambient gas is completed in the region of \( r \lesssim 1.5 R_s \), where the particles are thermally well coupled. As a result, any anisotropies obtained by the N-wave heating would be wiped out, and all the ions would have isotropic kinetic temperatures. The process of the N-waves cannot explain the observed anisotropies, which also indicates that other mechanisms causing anisotropic heating have to operate simultaneously.

We thank K. Shibata, K. Ohki, K. Omukai, K. Tomisaka, H. Saio, S. Nitta, S. Inutsuka, T. Kudoh, Y. Yoshii, T. Kajino, and A. Tohsoh, as well as members of DTAP in NAOJ, for many valuable and critical comments and Y. Meican for improvement of presentation in this paper. The author is supported by the JSPS Research Fellowship for Young Scientists, grant 5936.

REFERENCES

Allen, C. W. 1973, Astrophysical Quantities (London: Athlone)
Aschwanden, M. J., Poland, A. I., & Rabin, D. M. 2001, ARA&A, 39, 175
Bastian, T. S., Benz, A. O., & Gary, D. E. 1998, ARA&A, 36, 131
Cranmer, S. R., Field, G. B., & Kohl, J. L. 1999, ApJ, 518, 937
Erdélyi, R., Doyle, J. G., Perez, M. E., & Wilhelm, K. 1998, A&A, 337, 287
Furusawa, K., & Sakai, J. I. 2000, ApJ, 540, 1156
Habbal, S. R., Woo, R., Fineschi, S., O’Neal, R., Kohl, J., Noci, G., & Korendyke, C. 1997, ApJ, 489, L103
Hammer, R. 1982a, ApJ, 259, 767
———. 1982b, ApJ, 259, 777
Hasler, D. M., Rottman, G. J., Shoup, E. C., & Holzer, T. E. 1990, ApJ, 348, L77
Hayes, A. P., Pourlidas, A., & Howard, R. A. 2001, ApJ, 548, 1081
Hollweg, J. V. 1982, ApJ, 254, 806
———. 1999, J. Geophys. Res., 104, 24781
Hollweg, J. V., Jackson, S., & Galloway, D. 1982, Sol. Phys., 75, 35
Hudson, H. S. 1991, Sol. Phys., 133, 357
Kopp, R. A., & Orall, F. Q. 1976, A&A, 53, 363
Krucker, S., & Benz, A. O. 1998, ApJ, 501, L213
Kudoh, T., & Shibata, K. 1999, ApJ, 514, 493
Landau, L. D., & Lifshitz, E. M. 1959, Fluid Mechanics (London: Pergamon)
Landini, M., & Monsignori-Fossi, B. C. 1990, A&AS, 82, 229
Lee, L. C. 2001, Space Sci. Rev., 95, 95
Lee, L. C., & Wu, B. H. 2000, ApJ, 535, 1014
McWhirter, R. W. P., Thonemann, P. C., & Wilson, R. 1975, A&A, 40, 63
Mihalas, D., & Mihalas, B. W. 1984, Foundation of Radiation Hydrodynamics (New York: Oxford Univ. Press)
Nishio, M., Yajii, K., Kosugi, T., Nakajima, H., & Sakurai, T. 1997, ApJ, 489, 976
Osterbrock, D. E. 1961, ApJ, 134, 347
Parenti, S., Bromage, B. J. L., Poletto, G., Noci, G., Raymond, J. C., & Bromage, G. E. 2000, A&A, 363, 800
Parker, E. N. 1958, ApJ, 128, 664
Parnell, C. E., & Jupp, P. E. 2000, ApJ, 529, 554
Raymond, J. C., Suleiman, R., Kohl, J. L., & Noci, G. 1998, Space Sci. Rev., 85, 283
Roald, C. B., Sturrock, P. A., & Wolfson, R. 2000, ApJ, 538, 960 (RSW00)
Rosner, R., Tucker, W. H., & Vainsa, G. S. 1978, ApJ, 220, 643
Rosner, R., & Vainsa, G. S. 1977, ApJ, 216, 141
Sakai, J. I., Kawata, T., Yoshida, K., Furusawa, K., & Cramer, N. F. 2000, ApJ, 537, 1063
Sandberg, O., & Leer, E. 1994, ApJ, 423, 500
Sheeley, N. R., et al. 1977, ApJ, 484, 472
Stein, R. F., & Schwartz, R. A. 1972, ApJ, 177, 807 (SS72)
Strachan, L., Suleiman, R., Panasyuk, A. V., Biesecker, D. A., & Kohl, J. L. 2002, ApJ, 571, 1008
Sturrock, P. A. 1999, ApJ, 521, 451 (S99)
Sturrock, P. A., Roald, C. B., & Wolfson, R. 2000, ApJ, 538, L75
Tarbell, T., Ryutova, M., & Covington, J. 1999, ApJ, 514, L75
Tsuneta, S., et al. 1992, PASJ, 44, L63
Tsuneta, S. 1996, ApJ, 456, 840
Tsuneta, S., et al. 1992, PASJ, 44, L63
Ulmenschneider, P. 1971, A&A, 12, 297
Ulrich, R. K. 1996, ApJ, 465, 436
Withbroe, G. L. 1988, ApJ, 325, 442
Withbroe, G. L., & Noyes, R. W. 1977, ARA&A, 15, 363