Induced Emission of Alfvén Waves in Inhomogeneous Streaming Plasma: Implications for Solar Corona Heating and Solar Wind Acceleration

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Induced emission of Alfvén waves in inhomogeneous streaming plasma: implications for solar corona heating and solar wind acceleration

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The results of a self-consistent kinetic model of heating the solar corona and accelerating the fast solar wind are presented for plasma flowing in a nonuniform magnetic field configuration of near-Sun conditions. The model is based on a scale separation between the large transit/inhomogeneity scales and the small dissipation scales. The macro-scale instability of the marginally stable particle distribution function compliments the resonant frequency sweeping dissipation of transient Alfvén waves by their induced emission in inhomogeneous streaming plasma that provides enough energy for keeping the plasma temperature decaying not faster than r^{-1} – in close agreement with *in situ* heliospheric observations.

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The Sun's outer atmosphere, known as the solar corona, is considerably hotter than the Sun's surface, or photosphere. There are numerous speculative discussions and publications that attribute this heating either to waves [1, 2] or to (micro?, nano?) flares [3]. However, these explanations have thus far failed to provide a self-consistent quantitative model of a mechanism that heats the corona to millions of degrees and produces streams of fast solar wind in the heliosphere with the speed of thousands of kilometers per second.

There is an abundance of indirect observational evidence that appears to support wave based heating through resonant ion-cyclotron absorption of high frequency Alfvén waves. This evidence includes observations of preferential heating of heavy ions, obtained both through spectroscopic remote sensing [4] and *in situ* [5], as well as detected *in situ* anisotropic proton distributions at various locations within the solar wind [6]. But in spite of the plethora of available evidence, several attempts to develop a comprehensive quantitative model of coronal heating by resonant frequency sweeping of high frequency Alfvén waves have proved largely unsuccessful.

The original idea that the high frequency Alfvén waves can be used to heat the corona and drive the fast solar wind streams was suggested a while ago [1]. It stemmed from the speculation that the regions of strong magnetic field which define the boundaries of the chromospheric supergranulation network may release the available free energy accumulated in a twisted magnetic field through a series of impulsive reconnection events (or microflares). These events would then give rise to high frequency hydromagnetic waves with periods of much less than a second. As these high-frequency waves oscillate with about the same periods as the helical gyrating motion of ions in the presence of a strong magnetic field ($T \sim 1/\Omega_i \sim Am_pc/ZeB$ for ions with Z/A charge-to-mass ratio), the wave-particle interaction at cyclotron ($\omega_k - kv_\parallel = \pm\Omega_i$) resonance would be able to efficiently convert the available wave energy into heat. The waves with frequencies not in local resonance with any of the coronal ions would propagate away from the transition region through the corona in a rapidly declining magnetic field. In turn, this decline in the magnetic field strength would cause the change of the gyrofrequencies of ions, ultimately bringing those ions into resonance with yet-to-be-absorbed waves of lower frequencies.

This speculative scenario is the origin of the term *frequency (or wave) sweeping*. It was quickly included into a number of coronal heating models, both magnetohydrodynamic (MHD) [2, 7] and fixed shape kinetic [8] ones. But in order to simultaneously attain asymptotic values of the fast solar wind speeds and observed high values of coronal proton temperatures, these models had to either use an elevated level of waves in the high frequency range [2], or rely on collisional isotropization – even during collision-free expansion [7], or claimed that waves were completely unable to heat the corona and accelerate the solar wind [8]. Overall, the general conclusion reached by these models stated that the monotonically decaying spectrum of waves $|b_k|^2$ (up to a limiting flat power law $|k|^\rho$ with a single $\rho = -1$ exponent) generated at the base of the corona seems to be insufficient.

This negative result was further supported by the application of radiative transfer ideas from the theory of hot-star wind. The damping rates for more than 2000 species of various minor ions were calculated using the fixed shape of the Maxwellian ion distribution function and using observational spectroscopic constraints for their abundances. Then, the damping rates were summed to form a sort of "Alfvén optical depth" in Sobolev approximation [9]. The analysis culminated in the conclusion that frequency sweeping can not sufficiently heat the corona, as the low abundance minor ions will damp all available waves generated at the base of the corona well before they are able reach resonance with the coronal protons. Hence, an additional source of heating is needed not only for the coronal protons, but for the preferential heating of minor ions, as well. Consequently, the hunt for this "dark energy" content of the solar corona, as well as for the heliospheric/solar wind "dark energy" in general, has continued and intensified.

In this Letter, we report the self-consistent coronal heating and fast solar wind acceleration results from our semi-analytic scale-separation kinetic model – which was originally developed for the nonlinear treatment of cyclotron resonant wave-particle interaction in a streaming plasma imposed in a nonuniform magnetic field [10] and extended recently to coronal plasma expansion [11]. Our model calculates the velocity distribution function (VDF) dynamically from an energy balance [12] between ions and waves.

The main claim of this Letter is that this flexible form of the VDF invalidates all the conclusions about the inability of frequency sweeping to heat the corona and accelerate the fast solar wind. This supposed inability resulted from gross overestimations of Alfvén wave damping by height integration in an MHD framework. But in an inhomogeneous, non-equilibrium plasma of the solar corona, the streaming proton velocity distribution – and especially the velocity distributions of low abundance minor ions – quickly reacts and attunes to the presence of waves by displaying some sort of "wave-induced transparency." Moreover, the kinetic macroscale instability, originally discovered in [10] and investigated in more details in [11, 13], demonstrates that those flexible, wave-mediated kinetic ion distributions exhibit "induced emission" of Alfvén waves; that is, the coronal plasma acts as a source of Alfvén waves, rather than as a sink.

Our analysis is based on a numeric integration of the scale separation equations for the zeroth-order gyrophase averaged VDF $f_0(z, \mu, \Xi)$ for various sorts of ions

$$\frac{\partial f_0}{\partial z} + \frac{\sqrt{2\pi^3\lambda^3}}{L^2}\sum_k \hat{L}\left\{\delta(h_k^{\rightleftarrows\prime})\frac{|b_k^{\rightleftarrows}|^2}{\sqrt{h_k^{\rightleftarrows\prime\prime}}}\frac{\mu}{B_z}\hat{L}f_0\right\} = 0. \quad (1)$$

The summation in Eq. (1) includes the differential operator $\hat{L}$ that describes the particle response to the waves (with the local phase speeds $v_{ph} \equiv v_{ph}(z,k) = \omega_k/k$) in inhomogeneous non-uniform plasma and magnetic fields.

$$\hat{L} = \left(1 - \frac{v_{ph}}{\sqrt{u - 2\mu B_z}}\right)\frac{\partial}{\partial\mu} - \frac{2v_{ph}B_z}{\sqrt{u-2\mu B_z}}\frac{\partial}{\partial\Xi}. \quad (2)$$

Here, u denotes twice the kinetic energy ($u = v^2$), and the Dirac delta function $\delta(h_k^{\rightleftarrows\prime})$ is used to select only the inputs from resonances $h_k^{\rightleftarrows\prime} = 0$. L is the scale of the longest wave.

Eq. (1) was obtained [10, 11] by integrating particle trajectories in the (μ, Ξ, θ) coordinate system, where both the magnetic moment $\mu = v_\perp^2/2B_z$ and the total energy $\Xi = v^2 + \Pi(z)/A$ are conserved quantities for a single particle when waves are absent, and where θ is the gyration angle (used as an averaging parameter in obtaining Eq 1). The additional potential $\Pi(z)$ includes the inputs from all the volume forces acting on the solar wind particles, *i.e.* from gravitational force, as well as from the effect of an ambipolar electric field:

$$\Pi(z) = -\frac{GM_\odot Am_p}{z} + ZeT_e \ln n_e(z), \quad (3)$$

Z and A are the ratios of the ion to proton charge and mass.

This scale separation equation is obtained under the assumption of a fast energy exchange between waves and particles in each area of resonance by introducing particle phase $h_k^{\rightleftarrows}(z)$ for outward $\rightarrow$ and inward $\leftarrow$ propagating wave modes

$$h_k^{\rightleftarrows}(z) = -\omega_k\left(t(z) - \int^z dz'/v_{ph}^{\rightleftarrows}(z',k)\right) \pm \theta(z)/\lambda, \quad (4)$$

where parameter $\lambda = \Omega_i R_\odot/U_{sw} \gg 1$ represents the ratio of transit or inhomogeneity times (macrotimes) – defined by $R_\odot/U_{sw}$ – to the typical duration of the microscale processes $\sim 1/\Omega_i$. At standard near-Sun conditions, this ratio of timescales can be as large as 10^8 or even greater. Hence, the integrals along the particles trajectories have been asymptotically expanded, leaving only the leading resonant terms in λ (that is, neglecting the nonresonant wave terms that are $\sqrt{\lambda}$ times smaller).

The simplest evolution of waves is considered in the form of the dispersive wave kinetic equation [14],

$$\frac{\partial\omega_k}{\partial k}\frac{\partial W_k^{\rightleftarrows}}{\partial z} - \frac{\partial\omega_k}{\partial z}\frac{\partial|W_k^{\rightleftarrows}}{\partial k} = \lambda^{3/2}\gamma_k^{\rightleftarrows}W_k^{\rightleftarrows}, \quad (5)$$

where we introduced the wave action $W_k^{\rightleftarrows}$ as a ratio of wave spectral density $|b_k^{\rightleftarrows}|^2$ to wave frequency ω_k.

The wave equation includes only geometrical optics terms and the energy exchange term, where $\gamma_k^{\rightleftarrows}$ is the wave damping/growth rate due to resonant particles.

$$\gamma_k^{\rightleftarrows} = \frac{\sqrt{\pi}}{4}\int \delta(h_k^{\rightleftarrows\prime})\frac{\mu}{\sqrt{u-2\mu B_z}}\frac{\hat{L}f_0}{\sqrt{|h_k^{\rightleftarrows\prime\prime}|}}d\mu d\Xi. \quad (6)$$

The Dirac delta $\delta(h_k^{\rightleftarrows\prime})$ emphasizes that u should be substituted from the resonant condition $h_k^{\rightleftarrows\prime} = 0$ or

$$u = 2\mu B_z + v_{ph}^2\left(1 \pm \frac{ZB_z}{A\omega_k}\right)^2. \quad (7)$$

The wave phase speed v_{ph} is calculated in the cold plasma dispersion limit [11, 15]. Our approach could be improved further by replacing cold plasma dispersion with dispersion based on the flexible VDF from Eq. (1).

This closed set of time stationary Eqs. (1) and (5) has been solved numerically using a scale-separation energy balance approach [12]. This approach assumes that the energy exchange between waves and ions happens at the fastest rate; this can be formally understood as a consequence of the very large parameter $\lambda^{3/2}$ in front of both terms responsible for wave-particle interaction.

After the fast wave-particle relaxation, both the waves and the ions evolve at each integration step following the conservation laws inherently included in both equations – that is, the wave action conservation on the right-hand side of Eq. (5), and the magnetic moment μ and total energy Ξ conservation in the source free form $\partial f_0/\partial z = 0$ of Eq. (1). As a consequence of this energy and magnetic moment conservation, we can apply the Liouville theorem in a scale separated method of integration to find a simple expression for the solar wind speed by differentiating the energy conservation equation $< \Xi >=< (v_z + U_{sw})^2 + 2\mu B_z + \Pi(z) >=$ const. Thus, we obtain:

$$\frac{dU_{sw}}{dz} = \frac{1}{2U_{sw}} \left[-\frac{d\Pi}{dz} - 2\left\langle \mu \right\rangle \frac{dB_z}{dz} - \frac{d\left\langle v_z^2 \right\rangle}{dz} \right]. \qquad (8)$$

Here, $< \cdots >$ denote the moments of the distribution function calculated using the ion VDF $f_0(z, \Xi, \mu)$ obtained at each integration step.

We start with standard plasma and magnetic field values at the base of the coronal region. The proton temperature and density at the injection boundary ($z < 2R_\odot$) are equal to $\sim 10^4$ K and $\sim 10^6$ cm^{-3}, respectively. The large scale magnetic fields are approximately $\sim$1-10 G. Under these values, the Alfvén speed is approximately $\sim$1500 km s^{-1}.

We would like to emphasize that in its current form, our model does not include any collision-dependent effects and processes. Strictly speaking, this means that our model cannot be used at the coronal base and should only be applied after the coronal plasma has already reached the collisionless regime – that is, after $z \gtrsim 2R_\odot$. To partially compensate for the absence of collisional isotropization and heat conduction at the initial stage of coronal heating, the initial (pre-accelerated) bulk speed of ions U_{sw0} is chosen to be around $\sim$150 km s^{-1}, roughly on the order of the thermal speed at the collisionless expansion boundary.

The radial dependences of the solar wind velocity U_{sw}, of the Alfvén velocity v_A, and of the proton perpendicular temperature $T_\perp$ are shown in Figure 1. The initial level of Alfvén waves emitted at the base of the coronal region has been chosen to follow a power spectrum $|k|^\rho$ with exponent $\rho = -1$. The amplitudes of the spectral modes have been chosen to match – in the low-frequency limit – the low-frequency waves recently observed [16] within coronal holes, which have amplitudes at the level of $\sim$20-25 km s^{-1}.

Based on our previous estimates, for the particles that are part of resonant wave dissipation in the high-frequency range of the spectrum, the frequency scaling is proportional to $\propto |k|^{-2}$ [10]. Therefore, all typically observed and assumed Kolmogorov or Kraichman-like power spectra – i.e. those with either $\rho = -5/3$ or $\rho = -3/2$ – would produce quantitatively similar results for acceleration and heating. This is true unless the spectral exponent for the injected at the base waves is steeper than -2.

To provide some additional verification for our analysis, in Figure 1 we also include the results for non-self-consistent approach, which is obtained when the level of waves injected at the base is assumed to be large enough that their dissipation is negligible and the waves can be considered to be constantly present at all wave-particle resonances. In this purely wave-dominated expansion, analytical scalings for the radial dependence of the perpendicular proton temperature $T_\perp$ can be sketched easily.

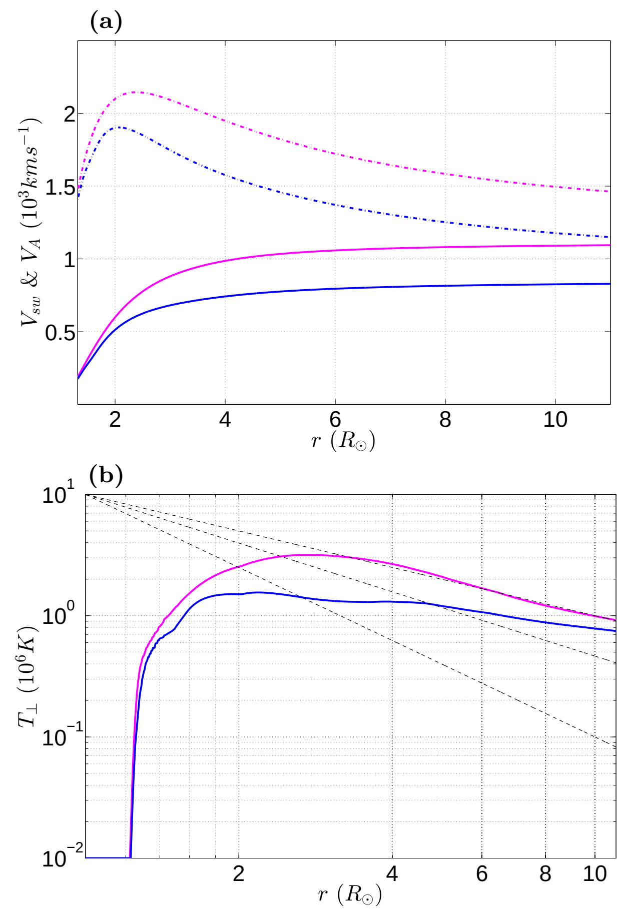

FIG. 1. The radial dependence of **(a)** the proton streaming speed U_{sw} (solid line) and the Alfvén speed v_A (dashed line), and **(b)** the proton perpendicular temperature $T_\perp$ in the coronal hole (bottom blue lines correspond to $\sim$20-25 km s^{-1} injection wave amplitudes, upper magenta lines - infinite wave energy flux). Dashed lines in **(b)** show theoretical scalings for CGL expansion (r^{-2}), for collisional expansion ($r^{-4/3}$), and for kinetic resonance sweeping cooling (r^{-1}).

The infinite supply of outward waves will convert the entire resonant region of the proton VDF to the pitch-angle scattered form. Unlike with dispersion-less waves, the diffusion lines are not simple circles [15]. In leading order, the way v_{ph} changes with distance mimics the way v_A changes; thus, for the proton perpendicular temperature scaling estimates, we can use v_A instead (Figure 2). The initial proton thermal speed v_T at the base of the corona is much smaller than the Alfvén speed ($v_T \ll v_A$); hence, $\alpha \sim \sqrt{2v_T/v_A}$, then $X \sim \alpha v_A$ and $T_\perp \sim X^2 \sim v_A \sim r^{-1}$. This is the same radial dependence of the proton temperature as the result of our numerical solution for the region with infinite waves. For realistic wave amplitudes, this radial dependence is even flatter (Figure 1), so it is in excellent agreement with observations by both Helios (0.3-1AU) and Voyager (1+AU) of $r^{-\epsilon}$, where ϵ is usually between 0.7 and 1.

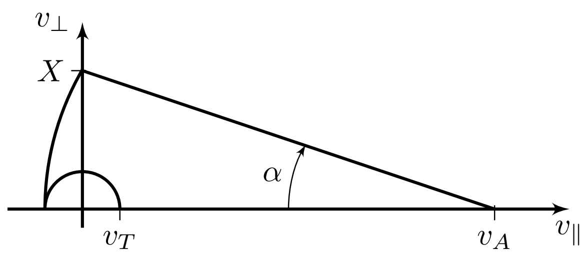

FIG. 2. Sketch of the pitch-angle scattered VDF, illustrating that the proton perpendicular temperature $T_\perp$ for $v_A \gg v_T$ scales as $X^2 \sim v_A \sim r^{-1}$.

We would like to note that just for the simple radial dipole magnetic field $B_z \sim r^{-2}$, the observed maximum of the Alfvén speed is captured remarkably well following the flux conservation in the rapidly accelerated fast solar wind.

Indirect observational evidence of ion heating by frequency sweeping, as well as of the validity of our treatment, can be found in the UVCS/SOHO spectral line profiles of different ions. Figure 3 (originally from [4], Figure 3) shows the spectral line width converted to velocity for the O^{5+} and Mg^{9+} ions. Surprisingly, these results were originally used as an argument against frequency sweeping, stating that since O^{5+} and Mg^{9+} have similar charge to mass ratios, one might expect their velocities to behave similarly if ion-cyclotron resonance is responsible for their broadening – but this is clearly not the case, since the rapid increase of width for these ions begins at different heights.

As a matter of fact, these plots prove exactly the opposite. If the heating (i.e. the rapid increase of width) of O^{5+} ions starts at $r_{O^{5+}} \sim 1.5R_\odot$, then for Mg^{9+} ions, the same frequencies will be in resonance at $r_{Mg^{9+}} \sim r_{O^{5+}}\sqrt{(Z/A)_{Mg}/(Z/A)_O} \sim 1.64R_\odot$. (This does not take into account a change in Alfvén speed with distance; taking it into account increases this value to something around $1.77R_\odot$.) And this is exactly the behavior of Mg^{9+} ions from the second panel. Thus, heating derived from the observed line width of both O^{5+} and Mg^{9+} ions shows features consistent with height-dependent frequency sweeping.

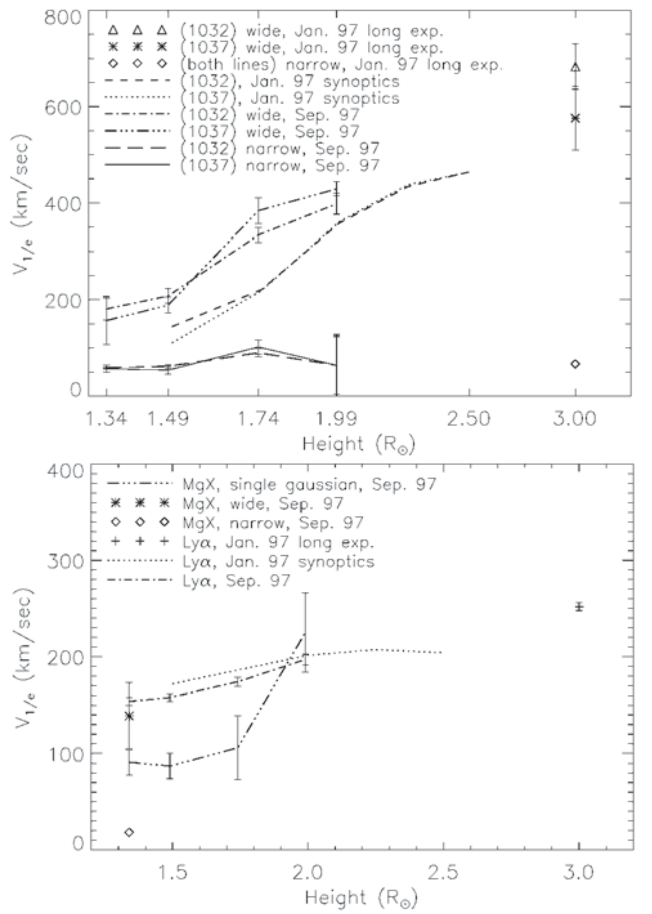

FIG. 3. The line width components expressed as velocities for O^{5+} ions (top) and Mg^{9+} ions (bottom) as a function of height (adopted from from [4], Figure 3).

In conclusion, we investigated a kinetic mechanism of heating of the solar corona and acceleration of the solar wind by frequency sweeping of Alfvén waves injected at the base of the corona. The mechanism operates in an inhomogeneous, non-equilibrium plasma of the solar corona and shows that ion VDFs quickly react to transient Alfvén waves and develop first the “wave-induced transparency” of the corona medium, and then produce more Alfvén waves by “induced emission” – creating a sort of Alfvénic laser by enabling the highly non-equilibrium streaming coronal plasma to act as a source of Alfvén waves, rather than as a passive sink.

The presented results are based on our use of a highly simplified model and at this stage can provide no more than a qualitative understanding of the physics of the solar wind. Quantitative understanding will require a far more complete model. Nevertheless, our claims that kinetic effects are essential to that understanding are well supported by these results.

We would like to emphasize again that Alfvén wave-mediated coronal plasma expansion can be adequately described only using fully kinetic treatment without any implied assumptions for the fixed shape (either Maxwellian, bi-Maxwellian, or fixed kinetic shell) of the ion velocity distribution. Even indirect assumptions based on the fixed shape considerations, such as speculations about the marginal stability of shell-like ion distributions, should be used with extreme care, as they grossly overestimate the resonant cyclotron damping of transient Alfvén waves and, as a result, destroy the heating of the corona and the acceleration of the solar wind.

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