Heating rate of ions via nonresonant interaction with turbulent Alfvén waves with ionization and recombination

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Ion heating due to turbulent Alfvén waves via nonresonant wave-particle interaction is discussed to complement a preceding work. We claim that newborn ions can get extra energy from turbulent Alfvén waves. Since in plasmas newborn ions are created continuously and intrinsically via ionization and recombination, heating associated with these newborn ions may be significant. The heating rate is proportional to the creation rate of newborn ions and the total wave field energy density without dependence on the profile of the wave spectrum. Based on observed nonthermal velocities, we estimate the energy density of the wave field and the corresponding heating rate in the solar atmosphere. The results show that such a heating process could be significant in the upper chromosphere and transition region of the sun. © *2009 American Institute of Physics*.
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I. INTRODUCTION

Plasma heating is an outstanding topic that occupies a special place in plasma physics and astrophysics. For example, the origin of the hot solar corona is still an unresolved issue for more than half a century since its discovery.[1] Because large amplitude Alfvén waves are constantly observed in the solar wind and because these waves are evidently generated near the sun,[2,3] solar physicists conjecture that Alfvén waves may play important roles for the creation of the solar corona.[4,5] On the other hand, ion heating due to Alfvén waves also attracted much attention in plasma physics. However, most of the existing discussions of the topic are based on the notion that emphasizes resonant wave-particle interactions.[6–13] It is believed that resonance may convert part of the wave field energy to particle thermal energy. Nevertheless, the plasma in the solar corona has very low beta values; thermal protons are not able to resonate with Alfvén waves. In order to resolve this and related issues, scientists proposed a number of scenarios that involve nonlinear processes.[14,15]

Recently, a new process is proposed.[16,17] It is suggested that protons can be heated by turbulent Alfvén waves via nonresonant interactions. Such a process is efficient in low-beta plasma. The process is demonstrated with two totally different theoretical methods. The major finding is that turbulent Alfvén waves can lead to enhanced stochastic ion motion. The key of the heating process is pitch-angle scattering.[16] Here we stress that this "heating" is actually parasitic to the turbulent wave field if there are no newborn ions. In other words, if hypothetically the turbulent Alfvén waves should "diminish," the predicted proton kinetic temperature would return to its original value while Alfvén waves are absent. We name the corresponding temperature as an "apparent temperature" and the heating process as pseudo.[18] However, in the present discussion we suggest that the situation is different if ions are continuously created due to ionization. In natural plasmas or partially ionized gases such as in the earth ionosphere, the solar chromosphere, and transition region, newborn ions are always created continuously and intrinsically due to ionization and recombination, which occur simultaneously and may reach some sort of quasiequilibrium state at certain given temperature.

After a neutral atom is ionized, it may be considered as a newborn ion. In this paper, we show that newborn ions can get more energy from the turbulent Alfvén waves than that gained by the background ions. We suggest that this extra energy does not disappear when the Alfvén waves subside. This heating of ions associated with this process is not reversible. Therefore it may have great significance for the discussion of the heating in the upper chromosphere and transition region.

The structure of the paper is as follows. In Sec. II, we discuss the essential different features between the background ions and newborn ions. Considering the ionization and recombination processes, we discuss the heating rate of ions in low-beta plasmas. This is studied in Sec. III. In Sec. IV and V, we briefly discuss possible applications to the solar atmosphere and summarize the primary results.

II. DIFFERENT HEATING OF NEWBORN IONS AND BACKGROUND IONS

A. General solution of ion motion equations

Without loss of generality we consider turbulent Alfvén waves propagating along the ambient magnetic field, $\mathbf{B}_0=B_0\mathbf{i}_z$. The wave magnetic field vector $\delta\mathbf{B}_w$ and electric field vector $\delta\mathbf{E}_w$ can be expressed as

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$$\delta \mathbf{B}_w = \sum_{+,-} \sum_k B_k^\pm(t)(\cos \phi_k^\pm \mathbf{i}_x \pm \sin \phi_k^\pm \mathbf{i}_y),$$
$$\delta \mathbf{E}_\mathbf{w} = -\frac{v_A}{c}\mathbf{b} \times \delta \mathbf{B}_\mathbf{w}, \quad \mathbf{b} = \frac{\mathbf{B_0}}{B_0}, \tag{1}$$

where $\phi_k^\pm = kv_A t - kz + \varphi_k^\pm$ is the wave phase, $\varphi_k^\pm$ is the phase constant, $\omega = kv_A$ is the wave frequency, v_A is the Alfvén speed, $\mathbf{i}_x$ and $\mathbf{i}_y$ are unit vectors, and $\pm$ represent right and left hand polarizations. Here we consider that the wave amplitude $B_k^\pm(t)$ is a function of time, which is different with Ref. 16. The motion of a ion is described by

$$\frac{d\mathbf{v}}{dt} = \frac{q_i}{m_i c}\mathbf{v} \times (\mathbf{B_0} + \delta \mathbf{B_w}) + \frac{q_i}{m_i}\delta \mathbf{E_w}, \quad \frac{d\mathbf{r}}{dt} = \mathbf{v}, \tag{2}$$

where m_i and q_i are the ion mass and electric charge and c is the speed of light in vacuum. Let $u_\perp = v_x + iv_y$ and $v_\parallel = v_z$, then we have

$$\frac{du_\perp}{dt} = -i\Omega_0 u_\perp + i(v_\parallel - v_A)S_k\Omega_k^\pm, \tag{3}$$

$$\frac{dv_\parallel}{dt} = -\mathrm{Im}(u_\perp S_k \Omega_k^{\pm *}), \quad \frac{dz}{dt} = v_\parallel, \tag{4}$$

where

$$\Omega_0 = \frac{q_i B_0}{m_i c}, \quad S_k = \sum_{+,-}\sum_k, \quad \Omega_k^\pm(t) = \frac{q_i B_k^\pm(t)}{m_i c} e^{\pm i\phi_k^\pm(t)}. \tag{5}$$

We assume that the turbulent waves change slowly and adiabatically with time in a low-beta plasma such that

$$\left| \frac{1}{B_k^\pm(t)} \frac{dB_k^\pm(t)}{dt} \right| \ll k|v_\parallel - v_A| \approx \omega_k \ll \Omega_0. \tag{6}$$

Note that

$$\frac{d\Omega_k^\pm}{dt} \approx \mp ik(v_\parallel - v_A)\Omega_k^\pm. \tag{7}$$

The solution of Eq. (3) can be generally approximated as

$$u_\perp(t) \approx e^{-i\Omega_0 t}\left(u_\perp(0) + S_k \left. \frac{v_\parallel - v_A}{\Omega_0 \pm k(v_A - v_\parallel)} \Omega_k^\pm e^{i\Omega_0 \tau} \right|_0^t \right). \tag{8}$$

In a low-beta plasma, $(v_A - v_\parallel) \approx v_A$ and $\Omega_0 \pm k(v_A - v_\parallel) \approx \Omega_0$, Eq. (8) can be simplified to

$$u_\perp(t) \approx u_\perp(0)e^{-i\Omega_0 t} - \frac{v_A}{\Omega_0}S_k\Omega_k^\pm(t) + \frac{v_A}{\Omega_0}S_k\Omega_k^\pm(0)e^{-i\Omega_0 t}. \tag{9}$$

The first term in Eq. (9) describes the gyromotion of the ion with respect to the background magnetic field. The second term is the transverse motion due to the electric field of the turbulent Alfvén waves, which is time dependent on the transient wave field strength. The last term may be considered as the modification of the gyromotion due to the presence of the turbulent Alfvén wave magnetic field, which depends only on the initial wave field strength when the ion entered into the system.[19] This solution is quite similar to the solution of Eq. (4) in Ref. 16. The only difference is that $B_k^\pm$ in Eq. (9) can be a function of time, while B_k in Eq. (4) of Ref. 16 is a constant. If at the initial time $B_k^\pm(0) = 0$, the third term in the right side of Eq. (9) will be zero.

B. Kinetic energy of background ions associated with turbulent Alfvén waves

In the following, we consider intrinsic turbulent Alfvén waves entering into homogenous multi-ion plasma. To demonstrate the different responses of background ions and newborn ions to the wave field, we let the wave field strength of each mode k develop from an initially infinite small value to a steady state with finite amplitude. In general, the background ions, which existed before the turbulent Alfvén waves entered the system, see the transient development of the wave field. Thus, $B_k^\pm(0) = 0$ for background ions, the last term in Eq. (9) will be zero for background ions. Solution (9) is reduced to

$$u_\perp(t) \approx u_\perp(0)e^{-i\Omega_0 t} - \frac{v_A}{\Omega_0}S_k\bar{\Omega}_k^\pm(t), \tag{10}$$

where $\bar{\Omega}_k^\pm(t) = q_i \bar{B}_k^\pm e^{\pm ik[v_A t - z(t) + \phi_k^\pm]}/m_i c$ and $\bar{B}_k^\pm$ is the final amplitude for wave mode k. The above result is applicable to a single ion. Let us now apply the above result to low-beta plasma, which consists of an ensemble of ions. If the characteristic spatial scale of our system is much larger than typical Alfvén wavelength, then we may take a similar ensemble average of $|u_\perp|^2$ over the initial position of each particle, as well as over time and gyrophase angle as done in Ref. 16. Then we find that the average kinetic energy of a background ion associated with turbulent Alfvén waves is

$$E_{\perp,\mathrm{background}} \equiv \frac{1}{2}m_i\langle |u_\perp|^2 \rangle$$
$$\approx \frac{1}{2}m_i\left(u_\perp^2(0) + v_A^2 S_k \frac{|\bar{\Omega}_k^\pm(t)|^2}{\Omega_0^2} \right) \tag{11}$$

or

$$E_{\perp,\mathrm{background}} \approx T_0 + \frac{m_i}{m_p}\frac{W_B}{n_p} \equiv T_0\left(1 + \frac{m_i}{m_p}\frac{1}{\beta_{p0}}\frac{\delta B^2}{B_0^2} \right), \tag{12}$$

where T_0 is the ion temperature at initial time without turbulent Alfvén waves, n_p is the number density of proton, $\beta_{p0} = 8\pi n_{p0}T_{p0}/B_0^2$ is the proton beta defined with its temperature T_{p0} without Alfvén turbulence, $\delta B^2 = \Sigma_k \bar{B}_k^{\pm 2}$ represents the final wave field strength, and $W_B = \delta B^2/8\pi$ is the wave field energy density. We implicitly assumed in Eq. (12) that heavy ions represent minority, so that the Alfvén speed is mainly determined by the proton density.

C. Kinetic energy of the newborn ions associated with turbulent Alfvén waves

Without loss of generality one can assume that newborn ions are created after the wave field reaches a "steady" state. Thus, $B_k^\pm(t) = B_k^\pm(0) = \bar{B}_k^\pm$ in Eq. (9) for newborn ions; the solution for newborn ions will be as same as the Eq. (4) in

Ref. 16. In other words, protons are actually considered as newborn ions in Ref. 16. Taking the ensemble average of $|u_\perp|^2$, we find that the average kinetic energy of newborn ions associated with turbulent Alfvén waves is

$$E_{\perp,\text{new}} \equiv \frac{1}{2} m_i \langle |u_\perp|^2 \rangle \approx \frac{1}{2} m_i \left(u_\perp^2(0) + v_A^2 S_k \frac{|\bar{\Omega}_k^\pm(t)|^2}{\Omega_0^2} + v_A^2 S_k \frac{|\bar{\Omega}_k^\pm(0)|^2}{\Omega_0^2} \right). \quad (13)$$

It is interesting to note that the second term on the right hand of Eq. (13) is dependent on the ensemble average of the transient wave field that each newborn ion experiences at time t. On the other hand, the third term depends only the ensemble average of the wave field that each newborn ion felt at the initial time when it entered the system. We will return to this in the next subsection. Because we consider that the turbulent Alfvén waves are already at a steady state when the newborn ions enter the system, we have $S_k |\bar{\Omega}_k^\pm(t)|^2 / \Omega_0^2 = S_k |\bar{\Omega}_k^\pm(0)|^2 / \Omega_0^2 = \delta B^2 / B_0^2$. Equation (13) can be expressed as

$$E_{\perp,\text{new}} \approx T_0 + 2 \frac{m_i}{m_p} \frac{W_B}{n_p} \equiv T_0 \left(1 + 2 \frac{m_i}{m_p} \frac{1}{\beta_{p0}} \frac{\delta B^2}{B_0^2} \right). \quad (14)$$

The factor 2 comes from the fact that the last term in Eq. (9) is not zero for newborn ions. Comparing Eq. (14) with Eq. (12), one can see that the average kinetic energy of newborn ion associated with Alfvén waves is two times of the kinetic energy of background ion, if we assumed the same value of final wave amplitude for the two cases.

D. Test-particle simulation and physical explanation

In order to verify that the analytic theory is indeed justified, we carry out a series of test-particle simulation. Two simulation cases A and B are studied, which are supposed for the background ions and newborn ions, respectively. The wave field increases from an initially infinite small value to a finite amplitude in case A, whereas the wave field has an initially finite amplitude in case B. To see whether the ion kinetic energy could be maintained when Alfvén waves are subsided, the wave field strength is finally withdrew to zero in both cases. The turbulent Alfvén waves are imposed and withdrew slowly and adiabatically in both cases.

The numerical scheme is similar to that described in Ref. 16. Fifty-one Alfvén waves are used in the simulation in which frequencies are uniformly distributed in the range $0.01\Omega_p < \omega < 0.05\Omega_p$. The amplitude of each wave mode is considered to be equal but changes gradually with time such that $\delta B^2(t) = \Sigma_k B_k^2(t) = \epsilon(t) B_0^2$, where

$$\varepsilon(t) = \begin{cases} \varepsilon_0 e^{-(t-t_1)^2/\tau^2}, & t < t_1 \\ \varepsilon_0, & t_1 \le t \le t_2 \\ \varepsilon_0 e^{-(t-t_2)^2/\tau^2}, & t > t_2 \end{cases} \quad (15)$$

for case A, and

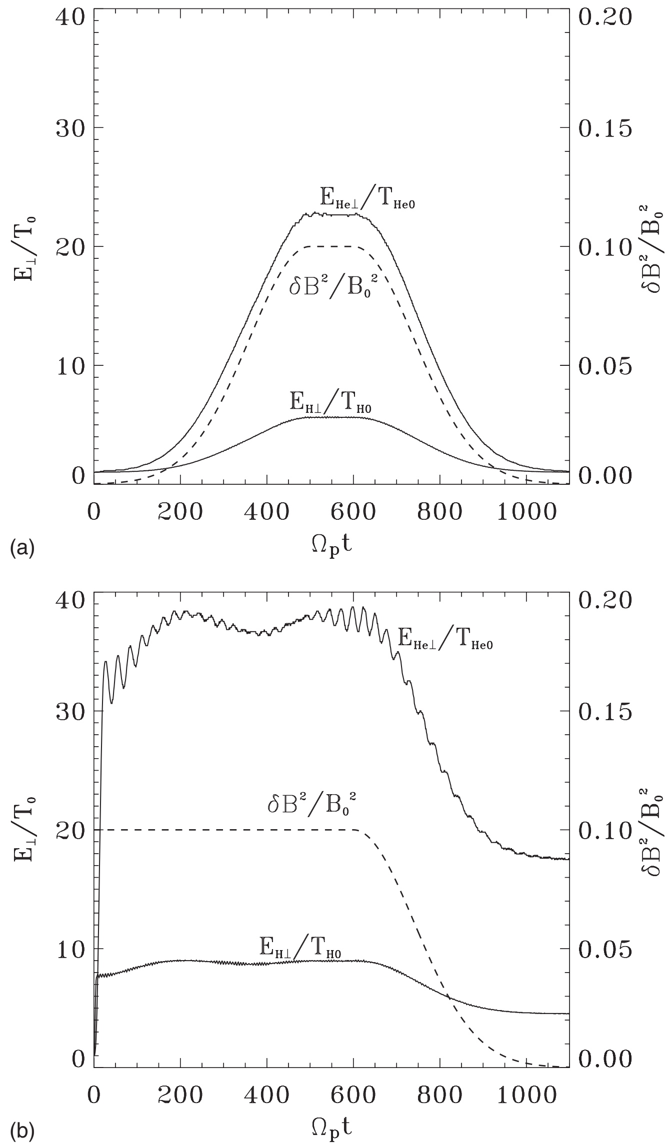

FIG. 1. Average ion kinetic energy and wave field strength acquired with numerical simulation are displayed for simulation case A (background ions, top panel) and case B (newborn ions, bottom panel), respectively. Profiles are plotted vs time.

$$\varepsilon(t) = \begin{cases} \varepsilon_0, & t \le t_2 \\ \varepsilon_0 e^{-(t-t_2)^2/\tau^2}, & t > t_2 \end{cases} \quad (16)$$

for case B. We suppose $t_1 = 500\Omega_p^{-1}$, $t_2 = 600\Omega_p^{-1}$, $\tau = 200\Omega_p^{-1}$, and $\epsilon_0 = 0.1$, where Ω_p is the proton gyrofrequency. The initial velocities of test particles are randomly distributed and possess a Maxwellian distribution with thermal speed of about $0.1 v_A$.

The temporal evolution of the magnitude of the turbulent waves and the heated "temperature" of protons and helium He^{+1} ions for the two cases are shown in Fig. 1. Figures 2 and 3 display the scatter plots of the test particle velocities in the velocity space at different times for cases A and B, respectively.

One can see that the average kinetic energy of newborn ions is about two times of the value of background ions, when the turbulent Alfvén waves are at a steady state. For

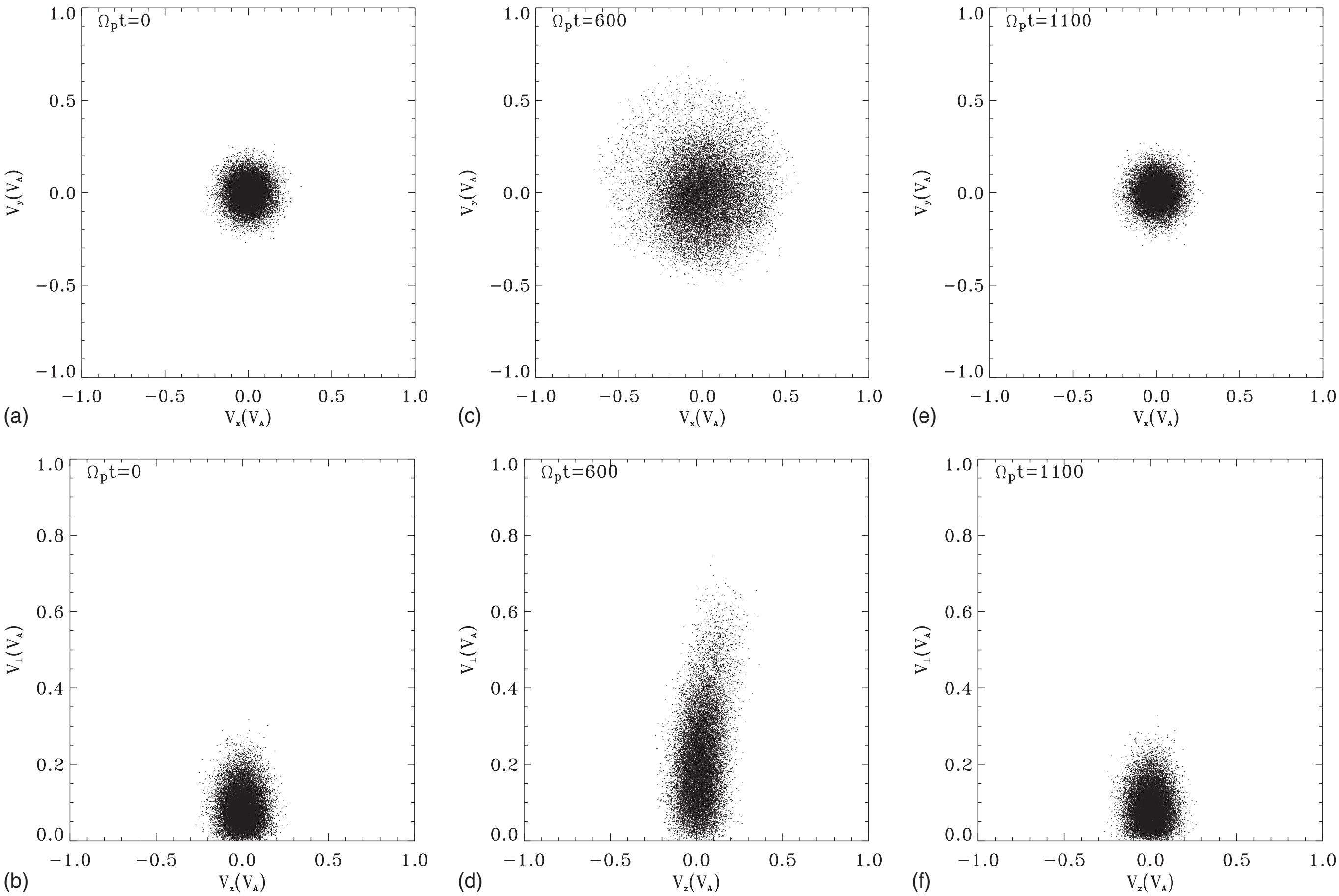

FIG. 2. Velocity scatter plots of protons for simulation case A (background ions) in the $v_x - v_y$ (top panel) and $v_{\perp} - v_z$ spaces (bottom panel) at time $\Omega_p t = 0$, $\Omega_p t = 600$, and $\Omega_p t = 1100$.

background ions, the kinetic energy is parasitic to the turbulent wave field. After the waves disappeared, it would return to its initial value without turbulent Alfvén waves. This can be also understood from Eqs. (9) and (11). The kinetic energy of background ions associated with the waves stems from second term on the right hand of Eq. (9) that is time dependent on the transient wave field strength. It would become zero if the wave field should subside. This heating process can be considered as "pseudoheating" and the corresponding temperature is an "apparent temperature," which is discussed in detail in Ref. 18. According to Eq. (12), the apparent temperature for ion species i is

$$T_i^{\text{apparent}} = \frac{m_i}{m_p} \frac{W_B}{n_p}. \tag{17}$$

On the other hand, the temperature of the newborn ions retains about half of its maximum after the waves disappear.

The kinetic energy of newborn ions associated with waves in Eq. (14) can be divided into two parts that come from the second term and the third term on the right hand of Eqs. (9) and (13), respectively. As discussed in the above subsections, the former is time dependent on the transient wave field strength, which is associated with the induced particle motion that tends to enhance the total wave energy as same as the background ions. It also causes an apparent temperature for newborn ions. Since the third term on the right hand of Eqs. (9) and (13) is independent on the time variation of the wave field strength, one can expect that the latter would not disappear after the wave field subsided. It is due to the "initial pitch angle scattering" when a newborn ion is injected into the finite amplitude Alfvén wave field.

To facilitate this scattering process, let us consider the situation in which intrinsic turbulent Alfvén waves pre-exist and "cold" newborn ions emerge. Here, "cold" means the ion initial thermal speed can be ignored in comparison with the Alfvén speed. This is a reasonable assumption in low-beta plasmas. If we choose to work in the wave frame, the electric field $\delta \mathbf{E}_w$ vanishes and the cold newborn ions are initially streaming toward the stationary waves with velocity $-v_A \mathbf{i}_z$ as shown in Fig. 4. A newborn ion would experience a force $e\mathbf{v}_A \times \delta \mathbf{B}_w(z_0)/c$, where z_0 is the position where the new ion born. This force will scatter part of the initial motion in the parallel direction along the ambient magnetic field to the gyromotion $v_{\perp B}$ in the transverse direction. It is easy to demonstrate that

$$v_{\perp B} \approx v_A \delta B / B_0$$

for $\delta B^2 \ll B_0^2$. Because the wave magnetic field $\delta \mathbf{B_w}$ consists of an infinite number of plane waves whose field directions and phases are randomly distributed, the force field $e\mathbf{v}_A \times \delta \mathbf{B}_w / c$ experienced by different newborn ions is with different strength and in different direction. Moreover, new-

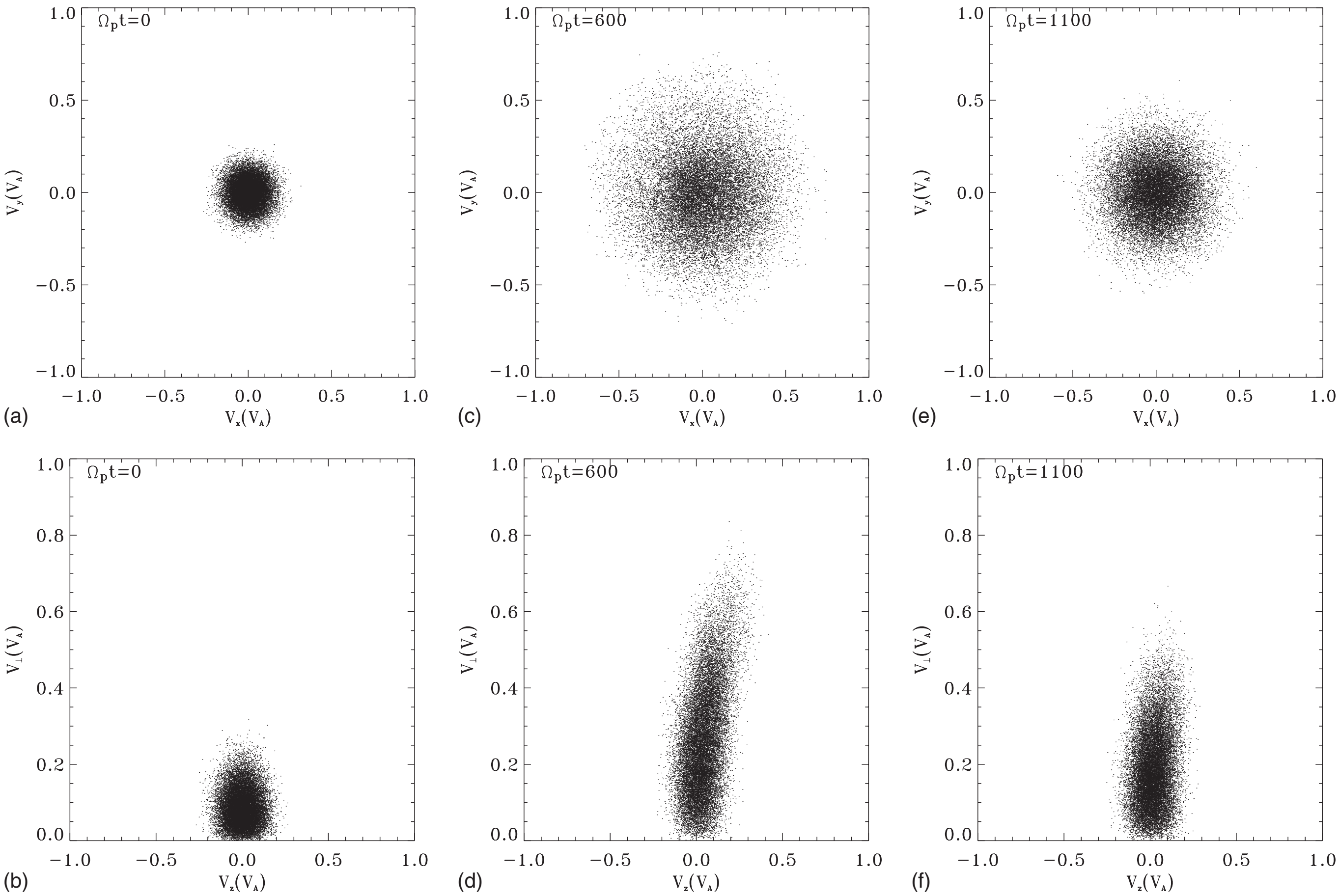

FIG. 3. Velocity scatter plots of protons for simulation case B (newborn ions) in the $v_x - v_y$ (top panel) and $v_\perp - v_z$ spaces (bottom panel) at time $\Omega_p t=0$, $\Omega_p t=600$, and $\Omega_p t=1100$.

born ions are generally created continuously in natural plasma. These give a natural phase mixing in the ion gyromotion leading to heating. Taking ensemble average of the ion gyromotion, the heated perpendicular temperature due to the initial pitch angle scattering can be approximated as

$$T_{\perp,\mathrm{new}}^{\mathrm{real}} = \frac{1}{2} m_i \langle v_{\perp B}^2 \rangle \approx \frac{1}{2} m_i v_A^2 \frac{\delta B^2}{B_0^2} = \frac{m_i}{m_p} \frac{W_B}{n_p}, \qquad (18)$$

which would not disappear after the turbulent Alfvén waves subsided as also demonstrated by the numerical simulation. It is genuine or real heating. It is interesting to note that Li *et al.*[20] also studied the ion pickup by a finite amplitude monochromatic Alfvén wave. Finally, we reiterate that this initial pitch angle scattering heating process does not work for background ions, because they experience an initially infinite small field strength.

III. ION HEATING RATE DUE TO IONIZATION AND RECOMBINATION

In natural plasmas or partially ionized gases, ionization and recombination occur intrinsically and continuously. Neutral atoms or molecules become ionized and ions recombine with electrons at the same time. For a given temperature, ionization and recombination would reach a statistical balance. After a neutral atom is ionized, it may be considered as a newborn ion. However, in principle, neutral atoms (or ions) are ionized (or recombined with electron) in random. After a newborn ion is created, one cannot distinguish it from the background ions. The separation of background ion and newborn ion in Sec. II is arbitrary; it is just for the sake of discussion.

When there are intrinsic turbulent Alfvén waves in homogenous plasma, the discussion presented in the preceding section shows that a newborn ion can get extra energy from the waves during the initial pitch angle scattering. Later on this newborn ion may also become a neutral atom via recombination with an electron. Intuitively it is expected that this newly created neutral atom would preserve the kinetic energy of the ion except some radiative losses. Meanwhile this neutral atom might be ionized again and it would get energy again from the waves. As long as the turbulent Alfvén waves exist enough long time, statistically, every ion would experience this recombination and reionization process at least once. Moreover, the large temperature anisotropy of newborn ions is expected to drive a number of kinetic instabilities due to thermal anisotropy that in turn will transfer energy to other particles.[21] Thus, initial pitch angle scattering of newborn ions would enhance the temperature ensemble averaged for all particles including both ions and neutral atoms. The heating rate is proportional to the creating rate of newborn ion and the extra energy that a newborn ion can obtain during the initial pitch angle scattering, which can be expressed as

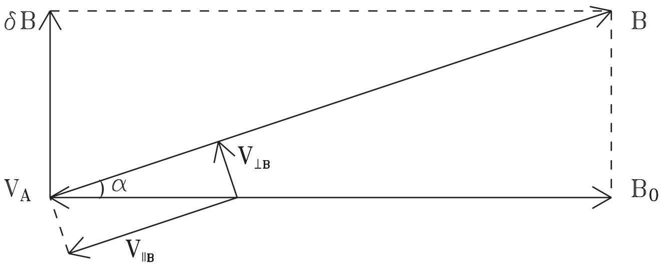

FIG. 4. Scheme plot of the velocity of a "cold" newborn ion and the magnetic field in a frame moving with the Alfvén waves. Here, cold means that the ion initial thermal speed is ignorable in comparison with Alfvén speed in a low-beta plasma.

$$Q_{w,i} = n_e \sigma_i N_i \frac{m_i}{m_p} \frac{W_B}{n_p}, \tag{19}$$

where n_e is number density of electrons, N_i and σ_i are the number density and ionization rate of neutral atom for species i, respectively. Note that the heating rate is proportional to the mass ratio of ion and proton.

When the plasma remains homeostasis, the system may readjust itself to maintain nearly equal rates of ionization and recombination during the heating process. It may be convenient to use the recombination rate of neutral atom instead of the creation rate of newborn ion, which can be expressed as

$$n_e N_i \sigma_i = n_e \alpha n_i = n_e(\alpha_r + n_e \alpha_3) n_i, \tag{20}$$

where n_i is the number density of ion species i, and α_r and α_3 are the radiative and three-body recombination rates. The radiative and three-body recombination rates for proton may be approximated as[22–25]

$$\alpha_{r,p} \approx 2.7 \times 10^{-13} \ T_e^{-1/2} \ \mathrm{cm}^3/\mathrm{s}, \tag{21}$$

$$\alpha_{3,p} \approx 8.75 \times 10^{-27} \ T_e^{-4.5} \ \mathrm{cm}^6/\mathrm{s}, \tag{22}$$

where T_e is the electron temperature in units of electron volt (eV). If we are interested in the region that cover the upper solar chromospheres and the bottom of the corona, where n_e is of the order of 10^{11}–10^8 cm^{-3}, T_e is of the order of 1–100 eV, so the three-body recombination may be ignored in comparison with the radiative recombination. Then the heating rate per unit volume for protons can be simplified to

$$Q_{w,p} \approx 2.7 \times 10^{-13} \ T_e^{-1/2} n_e W_B, \tag{23}$$

where all quantities are in Gaussian cgs units except the temperature expressed in eV.

IV. DISCUSSION AND CONSIDERATION OF SOLAR ATMOSPHERE

The preceding discussion and Eq. (23) imply that the higher the electron density and the lower the electron temperature, the higher the heating rate by turbulent Alfvén waves. This process may have important consequences in the solar chromosphere and transition region where the electrons have high density and low temperature.

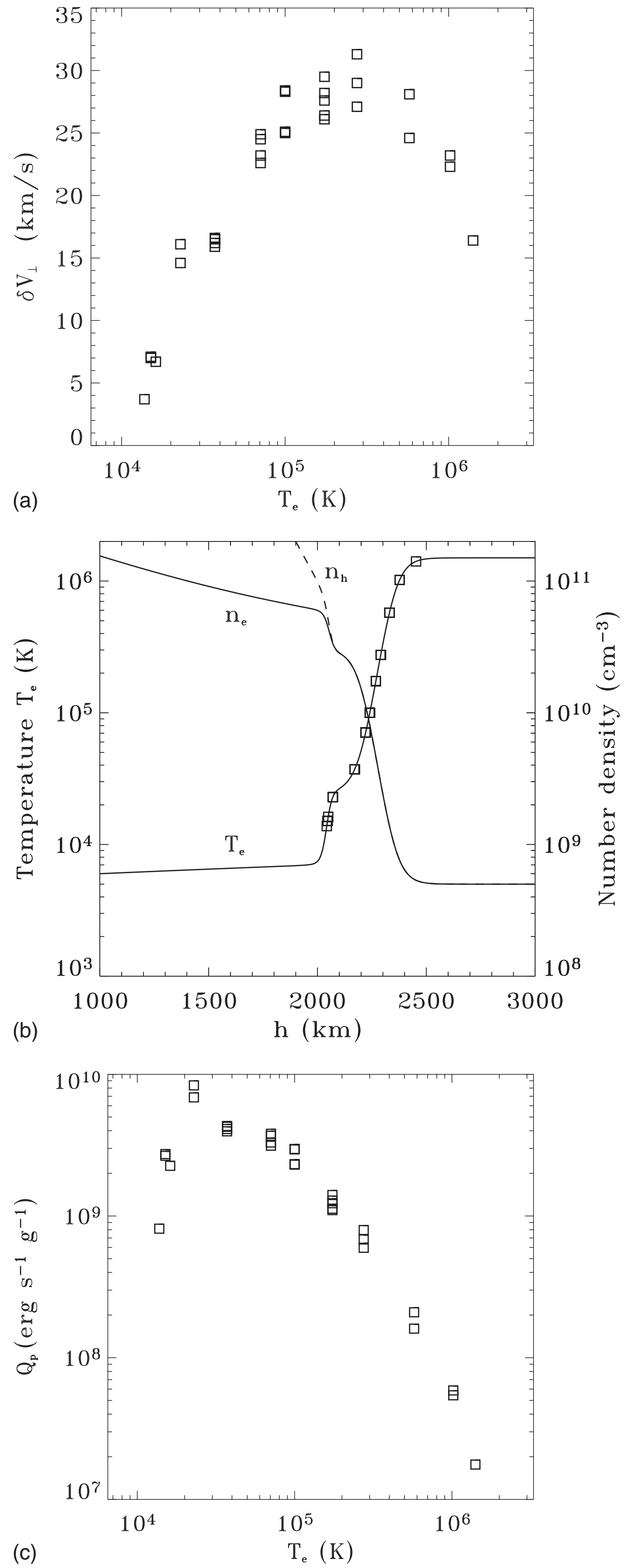

FIG. 5. Top panel: nonthermal velocities vs line formation temperature for the quiet Sun measured using SUMER on board the SOHO spacecraft (replotted from Fig. 5 in Ref. 26). Middle panel: altitude dependence of electron temperature and number density based on the VAL solar atmosphere model, where the square boxes indicate the height of each emission line corresponding to its formation temperature. Bottom panel: heating rate per unit mass in units of erg s^{-1} g^{-1} based on Eq. (23).

Chae *et al.*[26] calculate "nonthermal" velocities in the quiet Sun at temperatures between 10^4 and 2×10^6 K based on the observed linewidths with extra-ultraviolet and far-ultraviolet data taken with Solar Ultraviolet Measurements of Emitted Radiation[27] (SUMER) on board the Solar and Heliospheric Observatory (SOHO) spacecraft. Their results are replotted in Fig. 5 (top panel) for our purpose. Although these observations were mainly made on the solar surface, no significant center-to-limb variation was found from the measured nonthermal velocities. (We admit that the physical meaning of the observed nonthermal motion is still an open issue in solar physics.) In the present discussion we assume that near the limb where the line of sight is nearly perpendicular to the ambient magnetic field, the observed nonthermal velocities may be due to transversal random motion of ions associated with turbulent Alfveń waves. The corresponding energy density of the wave field can be estimated from the measured nonthermal velocities. On the other hand, the electron number density in the source region of each emission line may be estimated on the basis of the VAL solar atmosphere model.[28] The altitude dependence of electron number density and temperature is presented in Fig. 5 (middle panel), where the height of each emission line has been estimated by matching the altitude-dependence of temperature in the VAL model with the assumed formation temperatures of the ions that correspond to each emission line.[26] Then, using Eq. (23), the heating rate of protons due to ionization and recombination effect can be calculated in the solar transition region and the lower corona. The heating rate per unit mass is shown in Fig. 5 (bottom panel).

Recently, Cranmer *et al.*[29] presents a heating and acceleration model extending from the solar chromosphere to 1 a.u. Comparing our heating rate with the heating rate shown in their Fig. 7, one can see that two results are comparable. However, in the lower corona, our heating rate is much lower than the rate estimated in their model. Thus, we consider that ion heating by Alfvén waves due to ionization and recombination effect may be significant in the upper chromosphere and transition region, whereas it may be unimportant in the corona.

An interesting question is: could the pickup ion suffer collisions in one gyro-period? The traditional proton-proton collision frequency can be expressed as $\nu_{pp}=4.8\times 10^{-8}\ n_p\Lambda T_p^{-3/2}\ \mathrm{s}^{-1}$, where n_p and T_p are the proton number density (in units of cm^{-3}) and temperature (in units of eV) and Λ is Coulomb logarithm. For typical values in the solar upper chromosphere and transition region, let $n_p=10^9\ \mathrm{cm}^{-3}$, $T_p=10$ eV, and $\Lambda=20$, one gets $\nu_{pp}=30\ \mathrm{s}^{-1}$. Choosing the ambient magnetic field strength at 20 G, one can obtain the proton gyrofrequency at about $3\times 10^4\ \mathrm{s}^{-1}$. It is reasonable to consider that the pickup ions suffer no collision in one gyro-period in the upper chromosphere.

Finally, it is implicitly assumed in the above discussion that the turbulent Alfvén wave frequencies are higher than the neutral-ion collision frequency (few tens of Hertz in the solar upper chromosphere and transition region), so the motion of neutrals, or newborn ions before they are ionized, has nothing to do with the turbulent Alfvén waves. This turbulent Alfvén waves may be excited by energetic streaming ions produced by transient events such as explosive events and nanoflares observed throughout the upper chromosphere and lower corona regions.[30,31]

V. SUMMARY AND CONCLUSIONS

In this paper we discuss effects of ionization and recombination on the heating of ions by turbulent Alfvén waves via nonresonant interaction. These effects may be important in the upper chromosphere and transition region of the sun. In the following the main conclusions are summarized.

(1) Newborn ions can get more energy from turbulent Alfvén waves than the background ions. Due to intrinsic ionization and recombination processes, newborn ions are created continuously and consequently the resultant heating of ions is irreversible. The heating rate is proportional to the ion mass, the energy density of wave field, and the creation rate of newborn ions, without dependence on the special profile of the wave spectrum.

(2) Making use of the measurements of line broadening due to nonthermal velocity acquired with the SUMER/SOHO program we estimate the heating rate in the solar atmosphere based on our proposed scenario. The results seem to be encouraging and show that ion heating by Alfvén waves due to ionization and recombination effects may be significant in the upper chromosphere and transition region of the sun.

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