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The Roles of Turbulence on Plasma Heating*

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§1. INTRODUCTION

It is very important for the nuclear fusion research to obtain a high ion temperature plasma in a short time compared with the confinement time. So called "Turbulent Heating" is considered to be one of the strong candidates of the further heating methods, with which we overcome the limitation of the ohmic heating. Then, it is necessary by all means to clarify the physical mechanism of the turbulent heating of plasmas, because it provides scaling laws for heating up to fusion temperature by itself, and also it makes us to predict the effects of turbulence which may be appear when the imput power of other heating methods becomes considerably large.

Many experiments of the turbulent heating by high current discharge along the magnetic field have been studied both in linear and toroidal devices, and it has been shown that the ion temperature of the plasma becomes more than keV, however, the turbulent waves which are responsible for the particle heating have not been clearly identified, and the influence of the turbulence on the other plasma phenomena, such as diffusion across the magnetic field, is not so well investigated yet.

In this lecture we review the characteristic features of the turbulent heating referring to the works in IPP-Japan and discuss the roles of turbulence on plasma heating including the toroidal plasma heating from several points of view.

In §2 we summerize theoretically the relation between the heating rate of plasma particles and the thermalization (randomization) frequency and point out the important role of plasma turbulence in the fast thermalization. The general concept presented there could be applied not only in the high current turbulent heating, but also in the heatings with the use of turbulence due to the other means, for example, high power rf waves. In §3 we examine the experimental results, especially paing attension to the relation between the developement of turbulence and the particle heating, and also to the effect of scattering of particles by the turbulent waves. In §4 for application in the toroidal plasmas we consider the influence of the turbulence, being responsible for the particle heating, on the diffusion across the confinement magnetic field, and point out that the turbulent fields in the fast turbulent heating give only a minor effect to the loss of particles across the magnetic field. The final section is devoted to discussions and conclusions.

§2. RELATION BETWEEN THE HEATING RATE AND THE THERMALIZATION FREQUENCY

At the beginning we survay the general relation between the heating rate of plasma particles and the thermalization (randomization) frequency. As a simple example we consider the case of the high frequency electric conduction. The high frequency resistivity is expressed as

$$\eta(\omega) = \frac{m}{ne^2} (\nu + i\omega) , \qquad (2\text{-}1)$$

where m and e are the mass and charge of the electron, n is the number density , ω is the frequency of the electric field and ν is the classical collision frequency. Using Eq.(2-1) and the expression of the current $j = ne^2E_0 /m(\nu+i\omega)$, we have for the heating rate of the electron as

$$\frac{dT_e}{dt} \cong \frac{e^2E_0{}^2}{m\omega^2} \frac{\nu}{1+(\nu/\omega)^2} , \qquad (2\text{-}2)$$

where E_0 is the amplitude of the electric field. The coefficient of Eq.(2-2), $e^2E_0{}^2/m\omega^2$, corresponds to the mean energy with which an electron is stored during one period of oscillation, and the frequency, which gives the speed of thermalization, corresponds to the factor $\nu_T=\nu(1+\nu^2/\omega^2)^{-1}$. When $\nu<<\omega$, we have $\nu_T\simeq\nu$.

Generally speaking, if we write the energy, with which a particle in the plasma is stored during the characteristic time ω^{-1}, by w_s, the heating rate of the particle of that kind is given by

$$\frac{dT}{dt} \cong w_s \ \nu_T , \qquad (2\text{-}3)$$

where we assumed $\nu_T \ll \omega$ and disregarded the numerical factor of the order of one. The thermalization (randomization) frequenc ν_T agrees with the classical collision frequency ν when the randomization is due to the two particle collisions. However, in some cases ν_T is not always given by ν but is given by an enhanced value. We should find such a case of heating mechanism in the further heating of high temperature plasma. In the following we examine the some cases of heating methods in the above-mentioned viewpoint.

(a) Classical case;

A typical example of the classical case is Ohmic heating. If we write the drift velocity of the electron as v_D, the heating rate of the electron becomes

$$\frac{dT_e}{dt} \cong m v_D{}^2 \nu_{ei} \quad , \tag{2-4}$$

where ν_{ei} is the electron-ion collision frequency. From the same physical consideration we have for the neutral particle injection heating as

$$\frac{dT_i}{dt} \cong \frac{n_b}{n} \varepsilon_b \, \nu_{bi} \tag{2-5}$$

, where ε_b is the energy of the ion-beam resulting from the ionization of the neutral beam and n_b is the beam number density, and ν_{bi} is the collision frequency between the beam ion and the plasma ion. Above examples are the cases in which ν_T is directly equal to the classical collision frequency.

(b) Neo-classical case;

At the plasma heating in the magnetic field of confinement, the thermalization frequency is given by the classical collision frequency multiplied by an enhancement factor. We call these cases as neoclassical. One of the remarkable cases is TTMP (transit time magnetic pumping) which was firstly pointed out by Canobbio[1]. In TTMP the ions trapped in the magnetic mirrors formed with the pumping magnetic field make a substantial contribution to the heating rate of plasma ions. When the collision frequency ν is of the order of the bounce frequency $\omega_B \cong k\ v_T \sqrt{b}$ (k: wave number, v_T: thermal velocity of the ion, $b=\tilde{B}/B_0$, B_0: magnetic field strength of the confinement, $\tilde{B}$: magnetic field strength of the pumping wave), we have $\nu_T \simeq \omega_B$ because of the consideration the probability of the detrapping, and we have

$$\frac{dT_i}{dt} \cong w_s\ \nu_T \cong w_s\ \omega_B = b^2\ kv_T\ T_i\ , \qquad (2\text{-}6)$$

where we take $w_s \simeq \sqrt{b}\ \mu\tilde{B} \simeq b^{3/2}\ T_i$ considering the fact that the fraction of the trapped ions is about $\sqrt{b}$. This agrees with the result obtained by Canobbio. This is an example of the plateau of the heating rate expressed as a function of the classical collision frequency and has the same physical meanings as the plateau of the diffusion coefficient in ordinary space in the case of the toroidal confinement. Another example is the stochastic heating in ECRH in a magnetic mirror device[2]. In this case electrons are accelerated coherently in the local ECR regions produced in the mirror by rf power, however, outside this regions they cannot be effectively accelerated and move almost adiabatically. When the relations between the phases of the electron gyration and electric field oscillation at the entry in the resonance region are random, electrons do a random walk in $v_\perp$-space and the heating results. If we take the resonance time τ_r, the recurrence time to the resonance region τ_t, we have $\nu_T \simeq \tau_t^{-1}$ and

$$\frac{dT}{dt} \simeq m\left(\frac{e}{m} E_0 \tau_r\right)^2 \frac{1}{\tau_t} = \frac{e^2 E_0^2}{m} \frac{\tau_r^2}{\tau_t} . \qquad (2\text{-}7)$$

Considering the condition that the phase of gyration should be randomized in the course of the recurrence to the resonance region, the range of ν, in which ν_T is given by τ_t^{-1}, becomes

$$\frac{v_T}{L} < \nu < \frac{v_T}{\ell} , \qquad (2\text{-}8)$$

where v_T is the thermal velocity and ℓ and L are the length of the resonance region along the magnetic field line and the length of a travel until the recurrence to the resonance region respectively. As a matter of fact, the plateau of ν_T, in which the heating rate does not depend on ν, extends below the lower bound of (2-8), because electrons travel in the inhomogeneous magnetic field and their gyration phase are effectively randomized at the entry of the resonance region due to small angle scattering in the course of bounce motion[3]. This stochastic model of resonance heating has been applied in the ICR in the toroidal device[4].

(c) Turbulent Case;

In a quiescent plasma the collision frequency of the electron with ions is expressed by the relation,

$$\nu_c^e \simeq \omega_{pe} \frac{1}{n\lambda_D^3} , \qquad (2\text{-}8)$$

without a numerical factor of the order of unity, where λ_D is the electron Debye length. When the plasma becomes turbulent and the fluctuation level becomes larger than that in the thermal equilibrium, the relaxation of the motions of electrons turns out to be subject to the scattering by the turbulent wave fields. Then, the effective collision frequency ν^e_{eff} , which determines the electric resistivity and the heating rate of electrons, is given by the similar equation to Eq(2-8) which has an another factor in place of $(n\lambda_D{}^3)^{-1}$.

Nishikawa proposed a simple model of turbulence giving the anomalous resistivity, in which electorns are scattered by wave packets with effective charge Qe[5]. Let there be n_p wave packets in the unit volume and let ϕ and λ be the typical potential depth of the wave packet and its typical wavelength respectively, we get

$$\nu^e_{eff} \cong \frac{n_p}{n} Q^2 \nu_c{}^e \cong n\lambda^3 (\frac{e\phi}{T_e})^2 \nu_c{}^e \quad . \qquad (2\text{-}9)$$

Using Eq.(2-8) in Eq(2-9), we have

$$\nu^e_{eff} \cong \omega_{pe} \ (\frac{\lambda}{\lambda_D})^3 (\frac{e\phi}{T_e})^2 \ . \qquad (2\text{-}10)$$

Here we try to derive the expression of the effective collision frequency with another approach. By the quasilinear theory of the weak electrostatic wave turbulence[6] the effective collision frequency of the electron ν^e_{eff}, in which the resonant interaction of the electron and the electrostatic waves is taken into account, is given by the following arguments. The diffusion coefficient in velocity space becomes

$$D(v) = \pi \frac{e^2}{m^2} \int d\vec{k}\, E_k^2 \delta(\vec{k}\vec{v} - \omega_k) \quad .$$

Electrons suffer multiple scattering in the turbulent fields, therefore, the effective collision frequency is defined as

$$v^2 \nu^e_{eff} \simeq 2D(v) \ , \tag{2-11}$$

where the inverse of ν^e_{eff} corresponds to the effective deflection time. When we take $E_k = k\phi_k$ and assume the spectrum of turbulent field $E_k{}^2$ be isotropic and $\omega_k/k = v_p < v$, we have

$$\nu^e_{eff} \cong \frac{4\pi^2 e^2}{m^2 v^3} \int_0^\infty dk\ k^3 \phi_k{}^2 \cong \frac{e^2}{m^2 v^3} \cdot \frac{\langle\phi^2\rangle}{\lambda_c} \ , \tag{2-12}$$

where $\langle\phi^2\rangle$ is the ensemble average of fluctuation potential in space and λ_c is the characteristic wave length of turbulent waves. Rewriting Eq.(2-12) we get

$$\nu^e_{eff}(v) \cong \omega_{pe} \left(\frac{\lambda_D}{\lambda_c}\right) \left(\frac{v_{Te}}{v}\right)^3 \ \langle\left(\frac{e\phi}{T_e}\right)^2\rangle \ , \tag{2-13}$$

where v_{Te} is the thermal velocity of the electron. When $\lambda_c \sim \lambda_D$ and $v \sim v_{Te}$ we have

$$\nu^e_{eff} \simeq \omega_{pe} \langle\left(\frac{e\phi}{T_e}\right)^2\rangle \cong \omega_{pe} \frac{W}{nT_e} \ . \tag{2-14}$$

, where W is the energy density of the fluctuation field. In general we may express the effective collision frequency of the electron due to the turbulent wave as

$$\nu^e_{eff} \sim \omega_{pe} \frac{W}{nT_e} \ , \tag{2-15}$$

without a numerical factor depending on the model of turbulence. If we can estimate ν^e_{eff}, we obtain the heating rate of anomalous Ohmic process as

$$\frac{dT_e}{dt} \simeq m\, v_D^{\,2}\, \nu^e_{eff} \quad , \tag{2-16}$$

where $v_D = j/ne$ and $j = (\omega_{pe}^{\,2}/4\pi\nu^e_{eff})E$, and ν^e_{eff} is a function of the accelerating electric field E through the fluctuation energy density W.

Next we consider the ion heating in a turbulent plasma, where the effective scattering of ions by turbulent field are important. In ordinary cases we have $v_{Ti} \ll v_p = \omega_k/k$, then ions are not scattered resonantly. In the framework of the quasilinear theory, non-resonant interaction is given by the principal part of the diffusion coefficient and its effect is rather stochastic, namely[7]

$$D(v) \cong \left(\frac{e}{M}\right)^2 \int d\vec{k}\, \frac{\gamma_k}{\omega_k^2}\, k^2 \phi_k^{\,2}, \tag{2-17}$$

where γ_k is the growth(damping) rate of the wave with the wavenumber k and M is the ion mass. If we define the effective collision frequency of the ion by the inverse of deflection time due to the non-resonant interaction with waves, we have

$$\nu^i_{eff} \cong \frac{2D}{v^2} \cong \frac{4\pi e^2}{M^2 v^2} \int_0^\infty dk\, \frac{\gamma_k\, k^4 \phi_k^{\,2}}{\omega_k^{\,2}} \quad . \tag{2-18}$$

Let us introduce the characteristic frequency ω and characteristic wavelength λ, and the correlation time τ through the equation,

$$\int_0^\infty dk\, \frac{\gamma_k\, k^4\, \phi_k^{\,2}}{\omega_k} = \frac{\langle \phi^2 \rangle}{\lambda^2 \omega^2 \tau} \quad . \tag{2-19}$$

Then Eq.(2-18) becomes

$$\nu^i_{eff}(v) \cong \frac{1}{\tau}\left(\frac{c_s^{\,2}}{\omega^2 \lambda^2}\right) \frac{T_e}{Mv^2} \left\langle \left(\frac{e\phi}{T_e}\right)^2 \right\rangle \quad . \tag{2-20}$$

If we take $\lambda \sim \lambda_D$ and $\omega \sim \omega_{pi}$ in the case of ion acoustic wave turbulence and we put $\tau^{-1} = \alpha_c \omega$ where α_c is the ratio of τ^{-1} to ω, we have

$$\nu^i_{eff}(v) \cong \omega \quad \alpha_c \frac{T_e}{Mv^2} <(\frac{e\phi}{T_e})^2> \quad . \qquad (2\text{-}21)$$

Therefore the heating rate of ions is given by

$$\frac{dT_i}{dt} \cong Mv^2 \nu^i_{eff} = \omega \, \alpha_c \, T_e <(\frac{e\phi}{T_e})^2> \; , \qquad (2\text{-}22)$$

and the upper bound of the factor α_c is to be ~ 1[8], however, from the experimental evidence α_c is considered to be more less than unity. If τ^{-1} is given by the mode-mode couplings between turbulent waves and is of the order of the growth (damping) rate due to the non-linear interaction, α_c is estimated to be of the order of W/(nTe)[9], then we have

$$\frac{1}{\tau} \simeq \omega <(\frac{e\phi}{T_e})^2> \cong \omega \frac{W}{nT_e} \quad , \qquad (2\text{-}23)$$

and

$$\nu^i_{eff} \cong \omega \frac{T_e}{Mv^2} (\frac{W}{nT_e})^2 \quad . \qquad (2\text{-}24)$$

The comparative analyses with the experimental results will be given in the following sections.

§3. EXPERIMENTAL RESULTS ON FLUCTUATIONS AND HEATING OF ELECTRONS AND IONS IN TURBULENTLY HEATED PLASMAS

To investigate the effects of fluctuations on the diffusion in the velocity space for the electrons and ions, we describe the experimental results of THE MACH II Project[10], which has been carried out at the Institute of Plasma Physics, Nagoya University.

Plasmas with densities ($10^{12} \sim 10^{14}/cm^3$) and temperature ($2 \sim 10$ eV) are produced by plasma guns and injected into a mirror magnetic field (20 kG). Heating discharge current is fed to the plasma from the capacitors which is charged up to 30 kV, and plasma current goes up to several 10 kA. Fluctuations in the plasma are picked up by probes, and they are classified into two types. Type A fluctuations starts to be observed at the time t = 0.5 μs, when the drift velocity V_D exceeds the ion sound velocity C_s. Spectrum of this type is peaked at about 100 - 300 MHz, and the amplitude is about 10 - 100 Volt. Phase velocity is of the order of C_s. From those fact, we refer the type A fluctuations as ion acoustic turbulent waves. From the direct trace on the oscilloscope, this waves appears as a train of wave packet, and each packet contains several ($5 \sim 10$) waves. This indicates the correlation time τ_c of the turbulent waves is one or two period of the waves. The type B fluctuations appear when the plasma current shows resistive hump. The frequency spectrum peaked in the region of 5 - 10 MHz, which covers ion cyclotron frequency. The amplitude of the wave goes up more than 1 kV, and oscillate only several periods. The electric heating is measured by ruby laser scattering[11], diamagnetic loops and X-ray measurement. The temperature of the electrons reaches up to 10 keV at about the time of resistive hump on the current trace and it saturates. The heating rate is well described by the so called "anomalous ohmic process" $R_A I_H^2$, where R_A is the anomalous plasma resistance.

Ion heating is observed by use of charge exchange fast neutral measurement[10] and Doppler broadening measurement of ion lines[12]. Turbulent heating of hydrogen plasma indicates that the ion energy distribution has a bi-Maxwellian shape, main bulk ions and hot tail. In Fig. 2, typical results on the temporal variation of ion energy distribution of hot tail are shown. From this figure we find that a bump appears at t = 1.45μs (this is the time when the type B fluctuations started to be observed). This bump is thermalized within 0.8 ∿ 1.0 μs.

The mass dependence of the ion heating is measured by turbulent heating of the plasma which is dominantly helium with small amount of impurities. Fig. 3 shows typical results. The ions have two temperature distribution of the energy. The cold component has little dependence on the mass of the ions, however, the temperature of the hot component increases with the mass of the ions, while the number of the hot component decreases.

§4. PARTICLE LOSS ACROSS THE MAGNETIC FIELD

As is seen in the preceding sections, plasma turbulence promotes the thermalization (randomization) of the ordered motion of particles and heats them by itself, however, it enhances the spatial diffusion of particles ineviatbly although such an effect of turbulence should be undesirable. Especially for the heating of the toroidal plasma the diffusion across the magnetic field becomes a source of anxiety.

However, if the heating time required for the plasma temperature to reach the needed value is sufficiently shorter than the diffusion time across the magnetic field due to the same turbulent electric field as what is responsible for the particle heating, the turbulence could not substantially affect the confinement time.

Here let us consider the stochastic particle acceleration with the random electric field in the turbulent plasma. For simplicity, we ignore the effect of the magnetic field on the particle heating. The diffusion coefficient in velocity space is given by

$$D(v) = \frac{<(\Delta v)^2>}{2\ \Delta t} \cong \frac{e^2}{2m^2} \int_{-\infty}^{\infty} dt\ <\vec{E}(0.0)\vec{E}(\vec{r}(t),t)>$$

$$\cong \frac{e^2}{2m^2} <E^2>\ \tau_c \quad , \qquad (4\text{-}1)$$

where τ_c is the autocorrelation time of the random electric field along the particle orbit. Then, for the heating rate we have

$$\frac{dT}{dt} = 2mD \cong \frac{e^2}{m^2} <E^2>\tau_c \quad . \qquad (4\text{-}2)$$

The heating time required to reach the temperature T is given by

$$\tau_H = \frac{mT}{e^2 \langle E^2 \rangle \tau_c} . \tag{4-3}$$

On the other hand we consider the particle diffusion originated by the drift across the magnetic field due to the same turbulent electric field. Consider the cylindrical plasma with constant magnetic field B along the axis and let us take the component of the fluctuation field $E_\perp$ along the direction being perpendicular both to the magnetic field and to the radius of the cylinder. The spatial diffusion coefficient of the guiding center of the particle becomes

$$D_\perp = \frac{\langle (\Delta r)^2 \rangle}{2\Delta t} \cong \frac{c^2}{2B^2} \int_{-\infty}^{\infty} dt \ \langle E_\perp(0.0) E_\perp(\vec{r}(t),t) \rangle$$

$$\cong \frac{c^2}{2B^2} \langle E_\perp{}^2 \rangle \tau_c , \tag{4-4}$$

where we adopted the same autocorrelation time as in Eq.(4-1). The diffusion time for the cylindrical plasma with the radius a is given by

$$\tau_D \cong \frac{a^2}{2D_\perp} \cong \frac{B^2 a^2}{c^2 \langle E_\perp^2 \rangle \tau_c} . \tag{4-5}$$

From Eq.(4-3) and (4-5) we have the ratio of τ_H to τ_D as

$$\frac{\tau_H}{\tau_D} = \frac{D_\perp}{D} \cdot \frac{T}{ma^2} \lesssim \frac{T}{ma^2 \Omega^2} = \left(\frac{\rho}{a}\right)^2 \tag{4-6}$$

where we used $\langle E_\perp{}^2 \rangle \lesssim \langle E^2 \rangle$ and $\rho = (T/m)^{1/2}/\Omega$ is the Larmor radius of the particle with the thermal speed of the needed temperature. Under the condition of the toroidal confinement we should have $\rho << a$, then, we have $\tau_H << \tau_D$. Therefore, if we do the ion heating with use of the dissipation of the turbulent field, we have only a short heating time as compared with the diffusion time due to the same field. For example for $T_i = 10$ keV and $B = 50$ kG we have $\rho \cong 0.3$ cm, then, for a=10 cm we have $\tau_H \cong 10^{-3}\tau_D$.

When the diffusion due to the enhanced fluctuation is not described by the type of the equation as Eq.(4-4), that is, the classical type being proportional to B^{-2}, but is described by the Bohm diffusion being proportional to B^{-1}, the diffusion time becomes drastically short. And in the latter case the turbulence lead to rapid loss of particles across the magnetic field. However, it should be noted that the Bohm type diffusion is considered to take place due to the rather coherent wave excited near the surface of a plasma such as a drift wave[13], therefore, if we could make the fluctuation field turbulent, that is, we could make the correlation time more shorter, the dissipation of the field energy would be enhanced and the diffusion across the magnetic field would be suppressed. Then, if we have an ideal turbulent field with a short correlation time in the plasma, we should not be depressed with the loss across the magnetic field due to the turbulent waves.

§5. DISCUSSIONS AND CONCLUSIONS

It should be enphasized that the enhanced fluctuation in a turbulent plasma gives its field energy to the plasma particles while it can play the role of the fast thermalization of the ordered motions of particles that are produced in the plasma by some accelerating processes.

In the THE MACH II experiments the type A fluctuation starting at the early stage of the current discharge changes the plasma state turbulent and makes the transport coefficients anomalous. The type B fluctuation appearing in the stage of the resistive hump has a large amplitude, so that a fair amount of ions can be trapped in the potentials. If the observed bump in the energy distribution of ions is supposed to be formed by a large scale acceleration due to this potential, the relaxation in a short time of the order of 1 μsec should be caused with the anomalous process described in §2-(c). Suppose that the thermalization frequency of trapped ions in the waves of the type B fluctuation is given by the effective collision frequency due to scattering by the type A fluctuation. As the correlation time is of the order of the oscillation period, then we have $\alpha_c \simeq (2\pi)^{-1}$ in Eq.(2-21). If we take $Mv^2 \simeq T_e$ from the observations and use the value $\langle(e\phi/T_e)^2\rangle \simeq 3 \times 10^{-3}$ that is estimated from the observed anomalous resistance, we obtain $(\nu^i_{eff})^{-1} \simeq 1$ μsec, which is consistent with the observed speed of thermalization.

The elongation of the tail of the energy distribution of ions is thus attributed to the fast thermalization of the potential acceleration of ions due to the type A fluctuation with figh frequencies near the ion plasma frequency, while the heating of the cold component of ion distribution in Fig.3 can be explained by the nonresonant acceleration with the type A fluctuation and is described by Eq.(2-20).

The other acceleration processes related to the large amplitude waves in the magnetic fields should be also taken into account in the tail formation. Smith and Kaufman have showed that a part of trapped particles in the single large amplitude electrostatic wave propagating obliquely in the magnetic field execute a random walk in velocity space due to the overlapping of the harmonic resonance regions and stochastic acceleration of particles results[14]. The similar acceleration has been also shown to take place in the single wave propagating perpendicularly to the magnetic field[15].

The other possibility of the tail formation of the ion energy distirbution is due to the repetition of trapping and detrapping of ions in the wave traines propagating with the different angles to the magnetic field. A fraction of ions with velocities exceeding the phase velocity of these wave trains may be possible to gain their energies as accumulation of the velocity like Fermi's statistical acceleration, so that the mass dependent acceleration of the high energy tail observed in the experiments may be well explained.

As is seen in the above discussions, some possible processes producing the high energy ions may be considered, however, the enhanced fluctuation in the turbulent plasma always promotes the acceleration and thermalization of particles in any case , and these effects should be recognized as the most important role of turbulence in the further heating of a plasma with high ion temperature.

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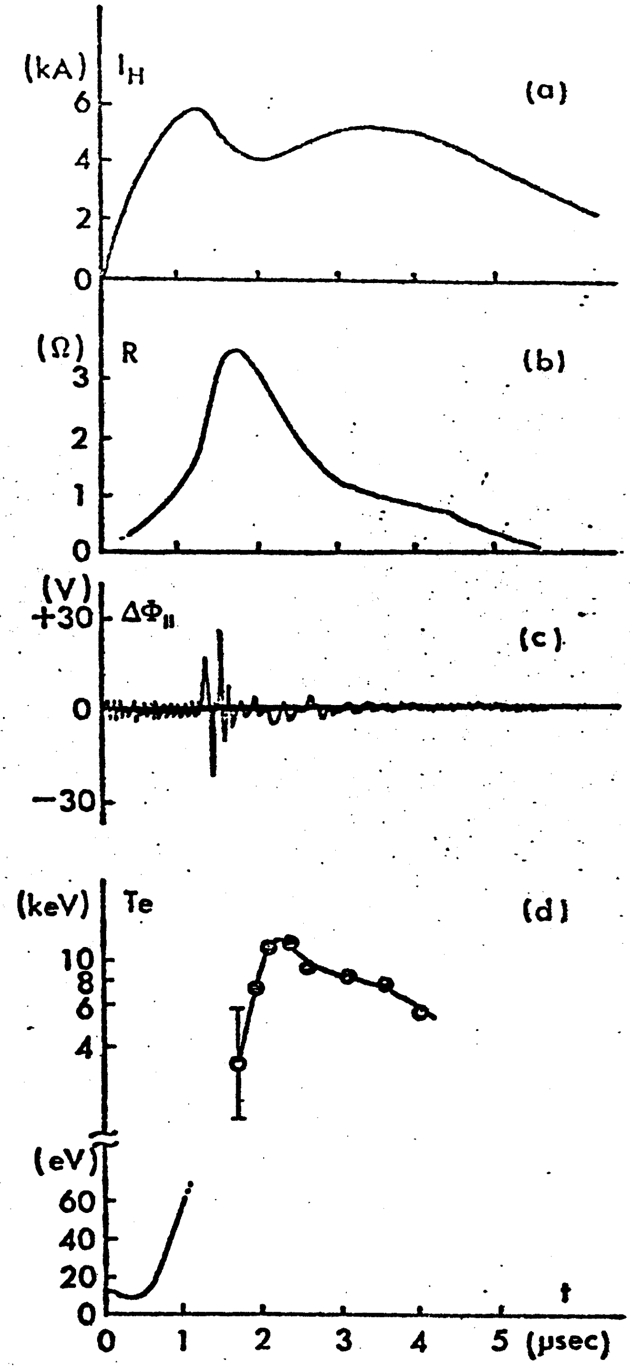

Fig.1:(a) The current of heating discharge, I_H.

(b) Resistance of the heated plasma column calculated from V_H, I_H and the inductance of the heating circuit.

(c) The floating potential difference between two probes set 2 mm apart each other along the axis of the machine.

(d) The electron temperature on the axis of the machine. Here, ruby laser scattering with an 8 channel detectors system was used for the temperature up to 70 eV, and X-ray measurement was used for higher temperature part.

[Reference 10]

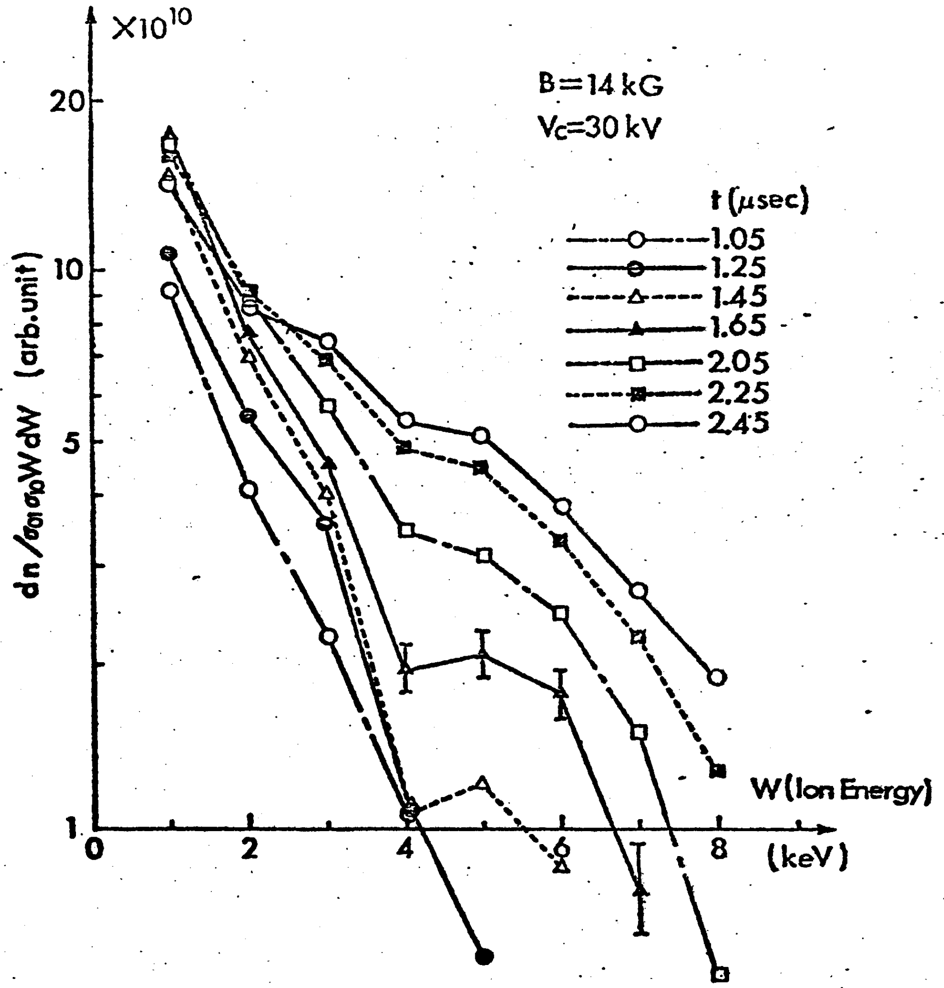

Fig. 2. Temporal evolution of the high energy part of the ion energy spectrum. The charging voltage of the capacitor for heating V_c is 30 kV.

[Reference: 19]

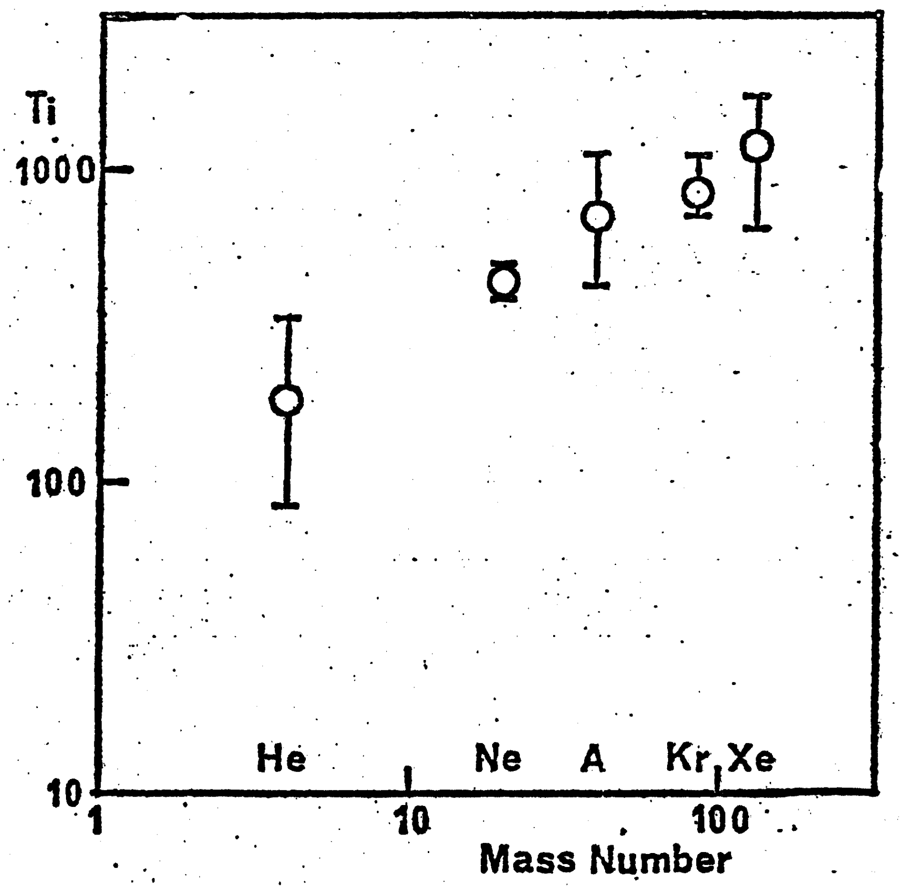

Fig. 3, Mass Dependence of Ion Temperature of Hot Component of Turbulently Heated Plasma Ions.
[Reference 12]