On the Preferential Acceleration

and Heating of Solar Wind Heavy Ions

by

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Abstract

We investigate the feasibility of producing the observed velocities and temperatures of solar wind heavy ions ($v_h \approx v_p + V_A$, $T_h \approx A_h T_p$) by the resonant cyclotron interaction with left-polarized hydromagnetic waves. We set up a "most favorable case" scenario. The waves are parallel-propagating and dispersionless. The energy for the wave acceleration and heating is taken from saturated low-frequency Alfvén waves via a cascade to higher frequencies. This scenario is incorporated into a numerical solar wind code and agreement with observation is tested. We find that the resonant cyclotron interaction fails on at least three points, even in this most favorable case. First, we show that no preferential heating or acceleration is obtained for heavy ions unless the wave power spectrum is steeper than k^{-2}, which is not often observed. Secondly, we find that, when sufficiently steep spectra are assumed, the resonant cyclotron interaction is unable to produce helium or oxygen velocities in excess of a fraction of the Alfvén speed with respect to the protons inside 1 AU. The simultaneous wave interaction with the protons prevents the heavy ions from attaining flow speeds of $v_p + V_A$ and this effect is understood in the form of a limiting differential speed. This limiting speed depends on A/Z of the heavy ion. Solar wind iron ions, with larger A/Z than helium or oxygen, can reach $v_p + V_A$ in this model. Thirdly, the preferential wave heating is insufficient to produce $T_h \approx A_h T_p$ at 1 AU, and we find that T_h/T_p is negatively correlated with proton speed, the opposite of the observed behavior.

I. Introduction

The properties of solar wind ions heavier than H^+ have presented a puzzle to theorists for some time (see Neugebauer, 1981 a,b; Hollweg, 1981a for recent reviews). The observation that the heavy ion velocity, v_h, frequently exceeds the proton velocity, v_p, in fast solar wind (Formisano, et al., 1970; Bollea, et al., 1971; Asbridge, et al., 1974, 1976; Hirshberg, et al., 1974; Ogilvie, 1975; Bosqued, et al., 1976; Marsch, et al., 1982a; Ogilvie, et al., 1982) has not been satisfactorily explained. Equally puzzling is the fact that the heavy ion temperatures are also larger than that of the protons, with the temperatures often appearing proportional to the masses of the ions, i.e. $T_h = A_h T_p$, where A_h is the atomic weight of the heavy ion (Robbins, et al., 1970; Ogilvie and Zwally, 1972; Bosqued, et al., 1976; Ogilvie et al., 1980, 1982; Schmidt et al., 1980; Bochsler and Geiss, 1982)

The most abundant heavy ion, He^{++}, has been studied most extensively and positive correlations have been found between the differential speed of the alpha particles with respect to the protons, $v_\alpha - v_p$, and both the proton and Alfvén speeds (Asbridge et al., 1976; Neugebauer, 1981a; Grünwaldt and Rosenbauer, 1978; Marsch et al., 1982a). The differential speed is also seen to be positively correlated with the alpha-proton temperature ratio (Neugebauer and Feldman, 1979).

The correlation with the Alfvén speed, V_A, is usually interpreted to indicate that the mechanism producing the preferential acceleration is a wave-particle interaction with electromagnetic waves having phase speeds of V_A. This interpretation is strengthened by observations closer to the sun. Marsch et al. (1982a) report that inside 0.8AU, the alpha particles are often seen to flow at the Alfvén velocity with respect to the protons, prompting the image of the helium ions "surfing" on the waves. Furthermore, Isenberg and Hollweg (1982) have shown that preferential acceleration produced by wave dissipation is accompanied by preferential heating, implying a natural explanation of the relation between $v_\alpha - v_p$ and T_α/T_p.

Most recent theoretical efforts to explain the observations have concentrated on the resonant cyclotron interaction with parallel-propagating ion-cyclotron or high-frequency Alfvén waves. The low-frequency analogs of these waves are commonly observed in the solar wind (Belcher and Davis, 1971; Behannon, 1978). The interaction is known to be capable of accelerating heavy ions from $v_h \leq v_p$ to $v_h > v_p$ (Hollweg and Turner, 1978; Dusenbery and Hollweg, 1981, hereafter designated DH; Hollweg, 1981a). Additionally, the fact that the proton temperature in fast solar wind is enhanced in the perpendicular direction (Bame, et al., 1975; Marsch, et al., 1982b) is consistent with the expected result of resonant ion-cyclotron heating.

However, the efficacy of this mechanism as the explanation of the observed preferential heating and acceleration has yet to be determined. Hollweg and Turner (1978) investigated the resonant acceleration of a test population of alpha particles by nondispersive Alfvén waves. The helium was assumed to move in a background electron-proton solar wind which was unaffected by the resonant interaction or by the presence of the helium. Furthermore, the temperatures of all the species were predetermined, so wave heating was not considered. Under these assumptions, the authors were able to obtain $v_\alpha \approx v_p + V_A$ for a variety of models.

DH reported on a parametric study which calculated both heating and acceleration rates for the resonant interaction with both dispersive and nondispersive ion-cyclotron waves. The heavy ions were again treated as test particles which did not affect the wave field, and there was no attempt to incorporate the resonant effects into a solar wind model. This work indicated strong acceleration and heating of heavy ions by the resonant interaction, but found that there could be limits to the local acceleration. For waves satisfying a cold plasma ion-cyclotron dispersion relation, the sign of the ion acceleration became negative when $v_h - v_p$ reached a certain fraction of V_A depending on the temperature (see also McKenzie and Marsch, 1982). A similar limit was found for nondispersive waves interacting with warm anisotropic ion distributions where $T_\perp > T_\parallel$ (the subscripts '$\perp$' and '$\parallel$' indicate directions perpendicular to and parallel to the ambient magnetic field direction, respectively).

Another difficulty with the resonant mechanism pointed out by DH concerns the energy requirement. When alpha particles are observed to flow at $v_\alpha = v_p + V_A$ in the inner solar system (where V_A is as high as 150 km/sec), the kinetic energy flux of the population is a substantial fraction of the total flux of solar wind energy. To accelerate the observed numbers of alpha particles to the observed velocities requires the energy equivalent to at least 20% of the observed wave field. Further, the resonant ion cyclotron waves only constitute the high frequency portion of the wave spectrum while the bulk of the power is found at much lower frequencies. Thus, the observed wave power in the resonant regime is generally several orders of magnitude below what is needed.

So, the importance of this mechanism remains an open question. The energy problem can be circumvented, at least in principle, by invoking a nonlinear cascade of energy from low frequencies (as suggested by DH), but other questions must be dealt with. One is the simultaneous resonant interaction with the protons. It is necessary to show that the resonant waves produce <u>preferential</u> acceleration and heating of heavy ions with respect to protons. This is a point that neither of the above papers considered. Then, the local resonant interaction must be incorporated into a solar wind model to gauge the effects of the global inhomogeneities of the system. For instance, the Alfvén velocity in the solar wind decreases with increasing radius and a mechanism which produced $v_\alpha \simeq v_p + V_A/3$ at 20 R_S (R_S = solar radius) might be expected to exhibit $v_\alpha \simeq v_p + V_A$ at .3 AU if no other forces are present. There are, however, counteracting forces such as Coulomb collisions and non-resonant wave interactions, which tend to equalize the velocities of different populations. The relative effects of the various forces must be determined from a model calculation.

In this paper, we investigate a scenario which deals with these questions. This scenario should be considered as a "most favorable case" picture of the resonant acceleration of heavy ions by ion cyclotron waves. In this picture, the energy for the resonant acceleration is siphoned from the low-frequency part of the wave spectrum. We will hypothesize a saturation of the Alfvénic oscillations, i.e. a nonlinear dissipation which limits the power contained in the low frequency part of the wave spectrum. The saturation process is assumed to feed energy to higher frequencies by an unspecified non-linear cascade providing sufficient energy in the resonant range to accelerate the heavy ions.

In view of the fundamental limits inherent in resonant acceleration by dispersive ion-cyclotron waves, we consider the interaction with dispersionless, parallel-propagating left-polarized waves - essentially high-frequency Alfvén waves. Where DH showed that cold dispersive waves cannot resonate with helium moving faster than $\approx 0.2\ V_A$ with respect to the protons, the Alfvén waves can affect all particles with $v_\alpha - v_p \lesssim V_A$, so the most favorable resonant process is an interaction with high-frequency Alfvén waves.

Finally, these wave processes are incorporated into the three-fluid, supersonic, spherically-symmetric, time-independent solar wind equations which include the radial dependence of the wave phase speed as well as the effects of non-resonant wave forces, Coulomb collisions, solar rotation, and gravity. The heavy ions are not treated as test particles in this calculation. This is particularly important for helium since Helios measurements (Marsch, et al., 1982a) have shown that the helium population can carry a substantial fraction of the energy of the solar wind in the inner solar system. Consequently, we allow the heavy ion fluid to carry momentum and energy and to exchange these quantities with the other fluids and with the waves.

The numerical calculations show that the resonant cyclotron interaction can indeed produce some preferential acceleration and heating of the helium and heavier ions, and some tendency for mass-dependent heating. However, even for this "most favorable" case, the resonant ion-cyclotron interaction is incapable of accelerating the helium population to $v_\alpha \approx v_p + V_A$ at .3AU. Furthermore, for all but the most extreme values of the parameters, this mechanism cannot produce $v_\alpha \gtrsim v_p + V_A/2$ for the entire range between $10R_s$ and 1AU. Similar results are obtained for oxygen ions though the mechanism is not as constrained by the observations in these cases. Numerical results for iron ions are also discussed.

The next section of this paper contains a detailed discussion of the resonant wave-particle interaction and the derivation of the relevant equations for our "most favorable case". In Section 3, the computational model is described. Section 4 consists of a discussion of the computational results and their interpretation. Section 5 summarizes our major points and conclusions.

II Wave-Particle Interaction for High-Frequency Alfvén Waves

a) "Most favorable case" scenario.

With the intention of investigating the most favorable situation for resonant preferential acceleration and heating of heavy ions by left-polarized waves, we consider the interaction with nondispersive parallel - propagating waves. Essentially, these are high-frequency Alfvén waves propagating along the average magnetic field with a phase speed which is independent of the wave parameters.

The condition for cyclotron resonance between a particle and a circularly polarized parallel-propagating wave is that the Doppler-shifted wave frequency in the frame of the particle's guiding center is equal to the gyrofrequency of the particle. This is

$$\omega - k_{\parallel} v^{r}_{i\parallel} = \pm\Omega_i \tag{1}$$

where ω and $\underline{k}$ are the angular frequency and the wavevector, respectively, of the wave, $\Omega_i = q_i B_o/m_i c$ is the gyrofrequency of the particle and the superscript 'r' denotes the resonant velocity. In the solar wind, we will consider both waves and particles to be moving away from the sun, so $k_{\parallel}$ and $v^{r}_{i\parallel}$ are of the same sign. When the particles are positively charged, the upper sign in (1) is appropriate for left-polarized Alfvén or ion-cyclotron waves and the lower sign corresponds to the right-polarized magnetosonic or whistler mode. Equation (1) can be solved for the parallel velocity of a resonant particle in terms of the phase speed $\omega/k_{\parallel}$:

$$v_{i\parallel}^{r} = (\omega/k_{\parallel})\ (1 \mp \Omega_i/\omega) \tag{2}$$

It is clear that a particle can only resonate with the right-hand mode if it is already moving faster than the phase speed of the wave. Since we are concerned with accelerating particles from speeds below the Alfvén speed up to this speed, we will concentrate on the left-polarized ion-cyclotron waves from this point, taking the upper sign in (1) and (2). Possible consequences of the right-hand interaction will be discussed in Section 4.

Dispersive ion-cyclotron waves have a phase speed which goes to zero in the plasma frame as $\omega \rightarrow \Omega_p$ in a plasma dominated by protons. It can be seen that the resonant velocity (2) for a heavy ion then has a maximum which occurs when the resonant wave frequency is in the range $\Omega_h < \omega < \Omega_p$. This maximum velocity is always less than V_A. For an alpha particle in a cold proton plasma, this maximum is $v_{\alpha\parallel}^{r} \approx 0.2\ V_A$ for parallel waves (DH) and is only slightly larger for perpendicular propagation (McKenzie and Marsch, 1982). This means that a particle moving along the field at a speed faster than this maximum cannot resonate with dispersive ion-cyclotron waves and the resonant interaction with these waves cannot accelerate heavy ions to V_A.

Nondispersive waves can interact with faster particles. Since we do not know the correct dispersion relation (which depends on the distribution functions of all the plasma constituents), we will consider the most favorable situation and assume the interacting waves to be parallel-propagating Alfvén waves, with a phase speed independent of wave parameters. There is some theoretical evidence (Gary, et al., 1975) that ion-cyclotron waves remain essentially dispersionless in a warm plasma, but our primary motivation in using nondispersive waves is that they seem capable, in principle, of accelerating heavy ions up to $v_h \simeq v_p + V_A$.

Another aspect of this mechanism to consider is the energy available in the resonant waves for the acceleration and heating of the alpha particles. As stated by DH, one needs about 20% of the total energy in the observed Alfvénic fluctuations to accelerate a number density of helium, $n_\alpha \simeq 0.04 n_p$, to $v_\alpha \simeq v_p + V_A$ at .3AU in the solar wind. Very little of the total power in the fluctuations is present at resonant frequencies and there is a serious problem in obtaining any substantial effect on heavy ions in the solar wind if one is restricted to using only the wave power observed in this range (see also Schwartz, et al., 1981).

There exists sufficient power in the low-frequency portion of the wave spectrum, however, and we propose a scenario which allows the resonant interaction to utilize this energy. We point out first that the low frequency waves are observed to dissipate as they propagate through the solar wind (Villante, 1980; Bavassano, et al., 1982a,b) though the mechanism is not understood. For the purpose of this investigation, we assume that a non-linear dissipation process saturates the wave activity at some relative amplitude, $\langle \delta B^2 \rangle / B_o^2$ = constant, where $\langle \delta B^2 \rangle$ is the total magnetic variance (integrated over the entire spectrum), and B_o is the magnitude of the average magnetic field. The energy lost to the wave field essentially comes from the large amplitude, low frequency end of the wave spectrum, but we further assume the operation of a cascade process which transports the energy from the saturated low-frequency range of the spectrum to the resonant range where it is picked up by the plasma through the resonant cyclotron interaction.

The details of the resonant interaction will be given by the quasi-linear expressions. The most favorable situation for the resonant interaction hypothesis would make the maximum amount of energy available to the resonant particles. Consequently, we assume that the resonant interaction is the only means of transferring the dissipated energy to the plasma and that the cascade is efficient enough to funnel all the power lost from the low-frequency waves into the resonant fluctuations.

The saturation and cascade processes will not be specified. We merely note that recent observations of fluctuation amplitudes suggest that the waves are saturated (Villante, 1980; Bavassano, et al., 1982a,b). Moreover, the typical power law character of observed wave power spectra has often suggested the operation of a non-linear "turbulent" cascade in wavenumber (Coleman, 1968; Barnes, 1981; Matthaeus and Goldstein, 1982). So we feel that the scenario outlined above is a plausible one. Finally, this saturation scheme (without the necessity of imposing a cascade hypothesis) has been used successfully by Hollweg (1978a) to model the high-speed flow of the solar wind.

b) Macroscopic equations

This scenario allows us to utilize the macroscopic results of Isenberg and Hollweg (1982) for multi-species acceleration and heating by Alfvén waves. In a time-independent system where the bulk fluid velocities of all species are directed along the average magnetic field (equivalent to the corotating frame in the solar wind), assuming all particle species to be thermally isotropic, they showed that the total plasma heating due to wave dissipation is related to the divergence of the wave action flux

$$\nabla \cdot (\omega \mathbf{S}) = - \sum_i Q_i \, (\omega/k_{||}) \, (\omega/k_{||} - v_{oi||})^{-1} \tag{3}$$

The flux of wave action for nondispersive Alfvén waves was found to be

$$\mathbf{S} = \frac{\langle \delta \mathbf{B}^2 \rangle}{4\pi} \frac{\omega}{k_{||}^2 V_A^2} \left(\frac{\omega}{k_{||}} - V_{CM} \right) \hat{\mathbf{b}} \tag{4}$$

where $\hat{\mathbf{b}} \equiv \mathbf{B}_o / B_o$. The volumetric heating rate of the i^{th} species is Q_i and its bulk velocity is $\mathbf{v}_{oi}$. The Alfvén speed is $V_A \equiv B_o \, (4\pi \sum_i \rho_i)^{-1/2}$ and the center of mass speed is $V_{CM} \equiv \sum_i (\rho_i v_{oi}) / \sum_i \rho_i$ where ρ is the density and the summations are over all species. The phase speed of these waves in a differentially flowing multi-ion plasma is

$$\omega/k_{||} = V_{CM} \pm \left[V_A^2 + V_{CM}^2 - \sum_i (\rho_i v_{oi}^2) / \sum_i \rho_i \right]^{1/2} \tag{5}$$

These results are valid for WKB Alfvén waves of arbitrary amplitude and arbitrary angle between $\mathbf{k}$ and $\mathbf{B}_o$. In the solar wind, most waves propagate out from the sun, in the same direction as the plasma, so we will take the upper sign in (5).

Note that the wave heating in a moving plasma is not given simply by the rate at which energy is lost by the waves. In general, the dissipation of waves adds momentum to the plasma as well as energy, so some of the wave energy must go into doing work. The conservation of wave action includes the effect of this additional energy loss.

The work done by these Alfvén waves in a steady solar wind can be related to the heating rate by the First Law of Thermodynamics. It is found, then, that the acceleration of a given plasma species $\underset{\sim}{a}_i$, due to the waves satisfies

$$\underset{\sim}{v}_{oi} \cdot \underset{\sim}{a}_i = \underset{\sim}{v}_{oi} \cdot \underset{\sim}{\nabla} \left[\frac{\langle \delta \underset{\sim}{B}^2 \rangle}{2B_o^2} \left(\frac{\omega^2}{k_{\parallel}^2} - v_{oi\parallel}^2 \right) \right] + \frac{Q_i}{\rho_i} \frac{v_{oi\parallel}}{\omega/k_{\parallel} - v_{oi\parallel}} \qquad (6)$$

In a spherically symmetric system, where all gradients are in the radial direction, (6) gives the change in flow energy of an ion species parallel to the field as

$$\frac{d}{dr} \left(\frac{1}{2} v_{oi\parallel}^2 \right) = \frac{d}{dr} \left[\frac{\langle \delta \underset{\sim}{B}^2 \rangle}{2B_r^2} \left(V_{phr}^2 - v_{oir}^2 \right) \right] + Q_i/\rho_i \left(V_{phr} - v_{oir} \right)^{-1} \qquad (7)$$

where the subscript 'r' denotes a radial component and $V_{phr} = (\omega/k_{\parallel})\ (B_r/B_o)$.

The first term on the right-hand-side of (7) is the non-dissipative wave pressure (Hollweg, 1974a,b, 1978b; Chang and Hollweg, 1976; Goodrich, 1978; McKenzie, et al., 1979) generalized to finite amplitude Alfvén waves with arbitrary directions of $\underset{\sim}{k}$ and $\underset{\sim}{B}_o$. This pressure force is an important source of solar wind acceleration. However, it acts on differentially flowing particle populations so as to equalize their bulk velocities, and therefore works against preferential acceleration. The second term in (7) is the acceleration which necessarily accompanies dissipative wave heating, and it is this term which will produce preferential acceleration resulting from the resonant interaction.

The scenario for the wave-particle interaction then proceeds as follows: for wave amplitudes below the saturation level, there is no dissipation, the heating rates are zero, and the waves evolve according to the conservation of wave action in a time-independent system

$$\underset{\sim}{\nabla} \cdot (\omega \underset{\sim}{S}) = 0 \tag{8}$$

In the expanding solar wind, (8) requires that $<\delta \underset{\sim}{B}^2>/B_o^2$ increase with increasing radius inside about 1AU. At some point, $<\delta \underset{\sim}{B}^2>/B_o^2$ will reach the designated saturation value and initiate the non-linear dissipation envisioned above. Beyond this radius, the wave amplitudes are given by the saturation condition and (3) determines the total wave heating. To proceed further, we must specify how this total wave heating is divided among the ion species, and this is determined by the microphysics of the resonant interaction.

c) Microscopic equations

The quasi-linear theory of the wave-particle interaction provides an expression for the rate of change of a particle distribution function $f_i(\underset{\sim}{v})$ due to the resonant interaction with a wave power spectrum $P(\underset{\sim}{k})$ (e.g. DH). The heating rate is then related to the second moment of $\partial f_i/\partial t$:

$$Q_i = 1/2\ \rho_i \int d\underset{\sim}{v} \left[(v_{\parallel} - v_{oi\parallel})^2 + v_{\perp}^2\right] \partial f_i/\partial t \tag{9}$$

For a bi-Maxwellian distribution interacting with a power-law spectrum of parallel-propagating left-polarized waves, the heating rate is

$$Q_i = \left(\frac{\pi m_i}{2\kappa T_{\parallel i}}\right)^{\frac{1}{2}} \rho_i \left(\frac{q_i}{m_i c}\right)^2 \int_o^\infty \frac{dk_{\parallel}}{k_{\parallel}} P_o \left(\frac{k_{\parallel}}{k_o}\right)^{-\gamma} \cdot$$
$$\cdot\ (\omega/k_{\parallel} - v_{oi\parallel}) \left[(1 - \Gamma)\ \Omega_i/k_{\parallel} + \Gamma(\omega/k_{\parallel} - v_{oi\parallel})\right] \cdot$$
$$\cdot\ \exp\left[-\frac{m_i}{2\kappa T_{\parallel i}} \left(\frac{\omega}{k_{\parallel}} - v_{oi\parallel} - \frac{\Omega_i}{k_{\parallel}}\right)^2\right] \tag{10}$$

where the power spectrum of magnetic fluctuations is taken to be $P(\underset{\sim}{k}) = P_o(k_{\parallel}/k_o)^{-\gamma}$, the temperature anisotropy is $\Gamma = T_{\perp}/T_{\parallel}$, and κ is Boltzmann's constant.

Dispersionless waves have a phase speed (5) independent of wave number so some terms can be factored out of the integral. The heating rate can be rewritten

$$Q_i = \left(\frac{\pi m_i}{2\kappa T_{\| i}}\right)^{\frac{1}{2}} \rho_i \left(\frac{q_i}{m_i c}\right)^2 P_o k_o^{\gamma} (\omega/k_\| - v_{oi\|})^{\gamma+2} \Omega_i^{-\gamma} \cdot$$

$$\cdot \int_o^\infty dx\, x^{\gamma-1} \left[(1-\Gamma)\, x + \Gamma\right] \exp\left[-C_i (x-1)^2\right] \quad (11)$$

where $x = \frac{\Omega_i}{k_\|}\left(\frac{\omega}{k_\|} - v_{oi_\|}\right)^{-1}$ and $C_i = \frac{m_i}{2\kappa T_{\| i}}\left(\frac{\omega}{k_\|} - v_{oi_\|}\right)^2$

If we assume that the integral is dominated by the resonance at $x = 1$ (i.e. $C_i \gg 1$), we find the proportionality relation

$$Q_i \propto \rho_i (A_i/Z_i)^{\gamma-2} (\omega/k_\| - v_{oi\|})^{\gamma+1} \left[1 + (\gamma-1)(\gamma-2\Gamma)/(4C_i) + \dots\right] \quad (12)$$

where Z_i is the ionic charge. In the numerical investigation to follow, we will assume that $T_{\| i}$ is small enough so that it is sufficient to keep only the leading term in (12). This is valid for the input conditions we will impose and it is likely to remain valid since the resonant interaction with ion cyclotron waves preferentially raises the temperatures perpendicular to the magnetic field. The calculations in DH showed that the parallel temperatures are even lowered in many cases by this interaction. Furthermore, their calculations indicate that when $v_{oh} - v_{op} \lesssim 0.5\, V_A$ (and $\gamma = 1.5$), the heating rate is essentially independent of temperature up to $2\kappa T_{\| h}/m_h \approx V_A^2$. Since the leading term in (12) does not depend on the temperature anisotropy, the effects of an anisotropic particle distribution will be neglected in the analysis to follow. This approximation and its implications will be discussed further in Section 4.

One final point remains to be made about this scenario for preferential acceleration and heating of heavy ions. If these preferential effects are to be produced solely by the resonant interaction with left-polarized high-frequency Alfvén waves, it is necessary that $Q_h/\rho_h > Q_p/\rho_p$, where Q_i is given by the leading term in (12). For $v_{oh\parallel} = v_{op\parallel}$, it is clear that the acceleration and heating process will only act preferentially on the heavy ions, $A_h/Z_h > 1$, when the wave spectrum is steeper than k^{-2}. Unfortunately, observed wave spectra do not appear to be steep enough (Behannon, 1978; Bavassano, et al., 1982a). Typical spectral indices for $\omega<\Omega_p$ (in the plasma frame) are on the order of $\gamma \approx 1.5$, with occasional observations as high as $\gamma \approx 2$ but no higher. The spectra do steepen to $\gamma \approx 4$ at high frequencies, but it is the range $\omega<\Omega_p$ (in the plasma frame) which determines the heating for the case of similar ion speeds.

In view of this discrepancy, we include one more assumption in the "most favorable case" picture: that $\gamma>2$. This is admittedly a questionable assumption. There are no observations at 20 R_S where the bulk of the interaction will take place in our model, but indications from radio scintellation measurements are that the spectra flatten as one approachs the sun, rather than steepening as required (Woo, 1981). Still, the resonant ion-cyclotron mechanism for preferential heating will not work without this assumption, so the "most favorable case" must incorporate it.

The scenario to be considered is now complete. When the waves have reached saturation amplitude, the total plasma heating is given by (3).

The relative wave heating of each species is determined by the leading term in (12) so the heating rate of each species can be found. Finally, the bulk acceleration of each species due to the waves is specified by (7). This wave interaction will be incorporated into the numerical solar wind model detailed in the next section.

III Computational Model

a) Assumptions

To test the effectiveness of our "most favorable case" heavy ion acceleration by high-frequency Alfvén waves, we incorporate the interaction scenario into a numerical solar wind calculation which uses the following assumptions:

1) The wave-averaged system is time-independent and spherically symmetric. Therefore, all bulk quantities have $\partial/\partial t = 0$ and vary only in the radial direction.

2) The calculations will be performed in the heliospheric equator and in the frame rotating with the sun. In this frame, all fluid velocities are parallel to the average magnetic field and the only large-scale electric field is the usual radial charge-separation field.

3) Since we are not concerned here with the origin of the solar wind per se, we will start the calculation at 10 solar radii (R_S) from the center of the sun, which is outside the sonic critical points for all parameters of interest. We will require that the parameter values at 10 R_S be such as to give reasonable agreement with observations given the simplifications made. This requirement allows us to choose proton input values which are consistent with those suggested by the wave-driven high-speed stream model of Hollweg (1978a).

4) The calculation will deal with three fluid species: electrons, protons, and a heavy ion species. We assume isotropic distributions for all species, characterizing the bulk properties by a flow velocity, $\underset{\sim}{v}_0$, and a temperature T. In the explicit calculation of the quasi-linear wave heating rate and of the Coulomb collision terms, the distributions are taken to be

Maxwellians, but this stricter assumption is not needed elsewhere.

Considerable particle anisotropies are observed in the solar wind, but the assumption of isotropy is in keeping with our "most favorable case" intent, at least as concerns the ions. Ion anisotropies will most seriously affect the wave-particle interaction, but the ion-cyclotron interaction will tend to produce $T_\perp > T_\parallel$ which reduces the heating (and acceleration), as can be seen from (12). Reduced heating as a consequence of $T_\perp > T_\parallel$ is also seen in DH. They found that the interaction of a pancake distribution with left-polarized waves results in deceleration and cooling for much of the parameter space. The effect of particle anisotropies on the resonant acceleration and heating will be discussed further in the next section.

5) The heavy ion species is allowed to affect the flow; that is, they are not test particles. They interact with the other species by Coulomb collisions, participate in determining the spiral field angle, the phase speed and evolution of the waves, and they take a finite amount of energy out of the waves when they are accelerated and heated. This aspect of the calculation is particularly important in the case of He^{++}.

By starting the computation outside the corona, however, we do not deal with the problem of heavy ion abundance in the solar wind. We will assume that some process such as Coulomb friction pulls the heavier ions out of the corona (Geiss, et al., 1970; Joselyn and Holzer, 1978; Hollweg, 1981b; Owocki, 1982) so that, at 10 R_S, the concentrations are such that subsequent motion produces observed values of n_h/n_p at 1 AU.

6) The electron density and bulk velocity are given by the requirements of quasi-neutrality and zero current flow. We neglect the electron inertial

terms and the thermal force (Braginskii, 1965) so that the motion of the electron gas is governed only by pressure forces, the electric field, and Coulomb interactions with the ions. The electron energy equation includes a collisionless heat flux (Hollweg, 1976), $\underset{\sim}{q}_e = 1.5\, n_e \kappa T_e v_{oe\parallel} \hat{\underset{\sim}{b}}$. Ion heat fluxes are neglected.

7) In the inertial frame, the center-of-mass velocity is taken to be radial, and the average magnetic field is carried out from the sun at this velocity. The tangent of the spiral angle of the average field is then

$$W \equiv B_\phi / B_r = - r\Omega_S / V_{CMr} \tag{13}$$

where $B_\phi{}^2 + B_r{}^2 = B_o{}^2$, Ω_S is the angular rotational velocity of the sun, and r is heliocentric distance. In this work, $\underset{\sim}{V}_{CM}$ is defined in the corotating frame, but it follows from this radial assumption that the center-of-mass velocity in the inertial frame is simply V_{CMr}.

8) Following the discussion of the previous section, the wave field is taken to consist of WKB, nondispersive Alfvén waves propagating along the average magnetic field. The wave power spectrum is a power law $\sim k^{-\gamma}$ with a total variance of $\langle \delta \underset{\sim}{B}^2 \rangle$. The propagation, dissipation and interaction with the solar wind plasma proceed according to the scenario and equations given previously.

b) Equations

We solve the steady-state, spherically-symmetric solar wind fluid equations in the corotating frame for the independent variables v_{opr}, v_{ohr}, T_e, T_p, T_h as functions of r.

The continuity, momentum, and energy equations for an ion species, i, under the above assumptions are

$$n_i v_{oir} r^2 = \text{constant} \tag{14}$$

$$\tfrac{1}{2} \rho_i \frac{d}{dr} \left[(1 + W^2) v_{oir}^2\right] = \rho_i(\Omega_S^2 r - GM_S/r^2) - \frac{d}{dr} (n_i \kappa T_i)$$

$$+ n_i q_i E_r + \rho_i \frac{d}{dr} \left[\frac{\langle \delta B^2 \rangle}{2B_r^2} (V_{phr}^2 - v_{oir}^2)\right]$$

$$+ Q_i (V_{phr} - v_{oir})^{-1}$$

$$+ (1 + W^2) \sum_{j \neq i} K_{ij} (v_{ojr} - v_{oir}) f(\eta_{ij}) \tag{15}$$

$$\frac{3}{2} n_i v_{oir} \kappa \frac{dT_i}{dr} - v_{oir} \kappa T_i \frac{dn_i}{dr} = Q_i + \sum_{j \neq i} 3K_{ij} \frac{\kappa(T_j - T_i)}{m_i + m_j} \exp(-\eta_{ij}^2)$$

$$+ (1 + W^2) \sum_{j \neq i} K_{ij} \frac{m_j}{m_i + m_j} (v_{ojr} - v_{oir})^2 f(\eta_{ij}) \tag{16}$$

where M_S is the solar mass, G is the gravitational constant, and E_r is the radial electric field. Q_i is taken equal to zero unless the waves are saturated. It is then determined by (from (3))

$$\frac{1}{r^2} \frac{d}{dr} (r^2 \omega S_r) = - \sum_i Q_i V_{phr} (V_{phr} - v_{oir})^{-1} \tag{17}$$

combined with the leading term in equation (12).

Equation (15) is obtained by taking the component of the momentum equation parallel to the average magnetic field and multiplying by $(\hat{r}\cdot\hat{b})^{-1} = \sqrt{1 + W^2}$. The terms on the right-hand-side of (15) are, respectively, the centrifugal force in the corotating frame, the gravitational force, the thermal pressure force, the force from the radial electric potential, the non-dissipative Alfvén wave pressure, the dissipative wave force, and the Coulomb friction. The right-hand-side of (16) contains the effects of heating due to wave dissipation, the energy exchanged with other species through Coulomb collisions, and the Joule heating from these collisions. The Coulomb terms are summed over all species $j \neq i$, including the electrons. The symbols in the Coulomb terms are defined by

$$K_{ij} = \frac{16\pi^{\frac{1}{2}}}{3} n_i n_j (q_i q_j)^2 \left(\frac{m_i m_j}{m_i + m_j}\right)^{\frac{1}{2}} \left[\frac{m_i + m_j}{2\kappa (m_i T_j + m_j T_i)}\right]^{3/2} \ln\Lambda \quad (18a)$$

$$n_{ij} = \left(\frac{1 + W^2}{2\kappa}\right)^{\frac{1}{2}} \left|v_{ojr} - v_{oir}\right| \left(\frac{T_i}{m_i} + \frac{T_j}{m_j}\right)^{-\frac{1}{2}} \quad (18b)$$

$$f(x) = 3/(2\, x^3) \left[\pi^{\frac{1}{2}}/2 \;\mathrm{erf}(x) - x \exp(-x^2)\right] \quad (18c)$$

where the Coulomb logarithm has been taken to be $\ln\Lambda = 21$. Except for the terms involving Q_i, equations (14) - (18) are equivalent to the system of equations written by McKenzie, et al. (1979); some typographical errors appearing in their paper have been corrected. The Coulomb collision terms are equivalent to those given by Burgers (1960).

The electron equations, subject to assumption (6), are

$$q_e n_e = \sum_i q_i n_i \tag{19}$$

$$q_e n_e v_{oer} = \sum_i q_i n_i v_{oir} \tag{20}$$

$$\frac{d}{dr}(n_e \kappa T_e) = n_e q_e E_r + (1 + W^2) \sum_i K_{ei} (v_{oir} - v_{oer}) f(\eta_{ei}) \tag{21}$$

$$3 n_e v_{oer} \kappa \frac{dT_e}{dr} - v_{oer} \kappa T_e \frac{dn_e}{dr} = \sum_i 3K_{ei} \frac{\kappa(T_i - T_e)}{m_i + m_e} \exp(-\eta_{ei}^2) + (1 + W^2) \sum_i K_{ei} \frac{m_i}{m_e + m_i} (v_{oer} - v_{oir})^2 f(\eta_{ei}) \tag{22}$$

where here, the summations are over all ion species. The momentum equation (21) is used to determine the electric field. The first term of the energy equation (22) includes the divergence of the collisionless electron heat flux.

With (21) inserted in (15) through the electric field, the equations (14) and (15) for each of the ion species, along with (22), form a set of five first order differential equations for the five variables: two radial velocities and three temperatures. Specification of the values of these quantities at the inner radius, taken to be 10 R_S, will determine the solution. Since 10 R_S is beyond any critical points of these equations for the parameters of interest, it is sufficient to simply propagate the solution outward from this position. The results of this calculation will be presented and discussed in the next section.

IV Numerical Results and Discussion

In this section, we will consider the question of whether the resonant cyclotron interaction with high-frequency Alfvén waves, in the "most favorable case" scenario we have constructed, is capable of accelerating and heating the observed heavy ion populations in the solar wind to the observed conditions. In this investigation, we will deal primarily with alpha particles as the heavy ion species, since observations of this species are most extensive and the constraints are therefore strongest.

We start with a "benchmark" solution for the helium case, using what we feel is a plausible set of initial parameter values at 10 R_S. We take values for the protons, electrons, and waves as suggested by the wave-driven high speed stream model of Hollweg (1978a):

$$v_{opr} = 300 \text{ km/s} \qquad T_e = 10^6 \text{ }^oK$$

$$T_p = 5 \times 10^4 \text{ }^oK \qquad B_o = 1600 \ \gamma$$

$$n_p = 5000 \text{ cm}^{-3} \qquad <\delta \underline{B}^2>/B_0^2 = 0.2 \tag{23a}$$

at 10 R_S. The waves will saturate when $<\delta \underline{B}^2>/B_0^2 = ½$, and the spectral index of the wave power will be taken as $\gamma = 4$. To test the ability of our most favorable case to preferentially accelerate and heat helium ions, we will start a typical concentration of helium at parameter values equal to those of the protons:

$$n_\alpha / n_p = 0.04$$

$$v_{o\alpha r}/v_{opr} = T_\alpha / T_p = 1 \tag{23b}$$

at 10 R_S.

The velocities and temperatures from this benchmark case are shown in Figure 1. Figure 1a is a plot of the ion radial velocities, along with $v_{opr} + V_{Ar}$, where V_{Ar} is the radial component of the Alfvén velocity. The ion velocities, starting together, remain together until the waves saturate at 19 R_S. At this point, the resonant interaction begins to operate and the alpha particles are preferentially accelerated. However, it is clear from the figure that the calculated alpha velocity exceeds the proton velocity by only a fraction of the Alfvén speed. Figure 1b displays the electron, proton, and alpha temperatures and the behavior here is similar. The ion temperatures evolve together until the waves saturate. The resonant interaction then heats both species strongly with preferential enhancement of T_α, but T_α / T_p at 1 AU is only 2.9, less than the "desired" value of 4.

Since the benchmark parameters do not work, we investigate the effect of modifying these parameters. We find that it is quite difficult to change the character of the benchmark solution, in particular to produce $v_\alpha = v_p + V_A$.

Increasing the initial wave amplitude (or equivalently, decreasing the imposed saturation level) does not improve the situation and indicates further conflicts with observation. Since the waves accelerate the protons as well as the helium, the result of increased wave energy is to create faster solar wind speeds, with both ion species flowing faster and the differential speed essentially unchanged. This development is shown in Figure 2, where we plot the values of v_{opr} and $\Delta v_{\alpha p} \equiv v_{o\alpha r} - v_{opr}$ at 1 AU, along with $\Delta v_{\alpha p}$ at 60 R_S, as functions of the initial mean square wave amplitude, $<\delta \underset{\sim}{B}^2>_i$, with all other parameters as in (23). The benchmark case is indicated by the vertical dashed line at $<\delta \underset{\sim}{B}^2>_i = 5 \times 10^{-5}$ Gauss^2. As $<\delta \underset{\sim}{B}^2>$ at 10 R_S is increased from the benchmark value, the flow speeds are enhanced, but the differential speeds are only slightly larger. For comparision, the Alfvén speeds in the benchmark case are V_{Ar} (60 R_S) = 103 km/s, V_{Ar} (1 AU) = 29.8 km/s, and these values increase by a factor ~ $\sqrt{2}$ from the left edge of Figure 2 to the right. It is clear that the inability of this model to produce larger differential speeds does not result from insufficient wave power. Furthermore, the strong correlation of $\Delta v_{\alpha p}$ with v_{opr} which is observed at 1 AU (Asbridge, et al., 1976; Grünwaldt and Rosenbauer, 1978; Neugebauer, 1981a) is only vaguely approximated in these results.

The effect of increased wave power on the ion temperatures is shown in Figure 3. Here, we plot the proton and helium temperatures at 1 AU for the same cases used in Figure 2. To aid comparison of these results with observations, we have plotted the values against the proton velocity (see Figure 2). The ion temperatures increase with v_p (and with the initial wave power) as expected, but T_α/T_p decreases dramatically. This trend in T_α/T_p is opposite to the observed positive correlation at 1 AU (Feynman, 1975; Bosqued, et al., 1977; Neugebauer, 1981a; see Neugebauer, 1981b for compiled observational results appropriate to Figures 2 and 3).

Investigation of other solutions shows that increasing the initial values of T_α, v_α, or γ fails to yield differential speeds equal to the Alfvén speed from this calculation. The effects of the wave dissipation are so pronounced that the initial velocities and temperatures of the ions are essentially irrelevant (if they are not so extreme as to make the dissipative interaction unnecessary), and the final relative ion parameters are determined by the dissipative interaction alone.

The failure of the resonant cyclotron interaction to produce the observed differential speeds can be explained qualitatively by invoking this dominance of the dissipative effects. We find that this interaction possesses a natural limit to the differential speeds, due to the simultaneous action of the waves on the protons.

Assume that the flow dynamics of all ion species are completely dominated, in some region, by the dissipative wave forces. Then, from (6), the acceleration of a given ion species along the magnetic field is

$$a_i = \frac{Q_i}{\rho_i\left(\frac{\omega}{k_{\parallel}} - v_{oi_{\parallel}}\right)} \tag{24}$$

Now consider the ratio of the wave acceleration of a heavy ion species to that of the protons, using the first term in (12) to give the relative heating rates. This gives

$$a_h/a_p = (A_h/Z_h)^{\gamma-2}(\omega/k_{\parallel} - v_{oh\parallel})^{\gamma}\ (\omega/k_{\parallel} - v_{op\parallel})^{-\gamma} \tag{25}$$

Again, we note that the waves will not preferentially accelerate a heavy ion species with respect to the protons unless $\gamma>2$. In addition, this mechanism will only preferentially accelerate an ion species up to a velocity such that $a_h/a_p = 1$. Further wave acceleration will act equally on both species and will not increase the differential speed Δv_{hp}. To estimate this critical differential speed, take $\omega/k_{\parallel} \simeq v_{op\parallel} + V_A$, and set the ratio (25) equal to one. The critical differential speed is then

$$\Delta v_{hp}^{*} \simeq V_A\left[1 - (Z_h/A_h)^{(\gamma-2)/\gamma}\right] \tag{26}$$

and this is a limiting value for preferential acceleration of heavy ions by cyclotron resonance with Alfvén waves. For the benchmark case of alpha particles and $\gamma = 4$, this limiting differential speed is only $\Delta v_{\alpha p}^* = .29\ V_A$ and it is clear from Figure 1a that this is too small to yield $v_\alpha = v_p + V_A$ at any subsequent point inside 1 AU. The critical differential speed (26) is also a weak funtion of γ, which explains why increasing γ in the calculation is not very effective in obtaining larger differential speeds. As $\gamma \to \infty$ for alpha particles, the critical differential speed approaches a maximum of $V_A/2$.

The critical differential speed in (26) was obtained by using the leading term in equation (12) which is equivalent to assuming very low parallel temperatures. This assumption is stretched somewhat at the temperature peaks of the benchmark solution (since the calculation uses isotropic temperatures), so we return to the full heating rate in (12) and ask what effect the correction term there has on the limiting speed (26). If $\gamma>1$, we find that, to lowest order, the Δv_{hp}^* will be larger (smaller) if ε is positive (negative), where

$$\varepsilon = (\gamma - 2\Gamma_h)\ T_{\parallel h}/A_h - (\gamma - 2\Gamma_p)\ T_{\parallel p}\ (Z_h/A_h)^{2(\gamma-2)/\gamma}$$

This is an involved expression, but if the ion temperatures are proportional to the ion masses, the anisotropies are equal, and $\gamma>2$, then the limiting differential speed will be lowered if the anisotropy is large enough: $\Gamma>\gamma/2$. Since resonant ion-cyclotron heating generally increases Γ, we expect that warm plasma corrections to (26) will not result in significantly larger differential speeds. However, this point clearly needs further study.

Given the existence of this critical differential speed, it is easy to understand why the calculation is so insensitive to the initial helium parameters. Increasing v_α at 10 R_S simply means that the critical speed is reached sooner, resulting in less preferential heating and smaller helium temperatures. Any reasonable increases in T_α at 10 R_S are either eliminated by the Coulomb interaction before saturation or are simply swamped by the overwhelming wave heating.

However, the critical differential speed in (26) has been derived by neglecting the action of thermal pressure forces compared to the dissipative acceleration. This is valid for most cases of interest, but it is possible to produce numerical results in which Δv_{hp} substantially exceeds the critical differential speed through excessive pressure forces. This can be done by lowering the initial heavy ion velocity, as seen in Figure 4. In this "overshoot" case, the helium velocity at 10 R_S was reduced to one-half of the proton velocity, with the helium density adjusted to retain an abundance ratio of 4% near 1 AU. All other input parameters are given by (23). Initially, Coulomb collisions act to equalize the ion velocities and produce some heating of the ions. However, at saturation the helium is still moving slower than the protons. The wave heating rate is a strongly increasing function of the difference between the wave phase speed and the bulk speed of the ion. Thus, after saturation the helium ions are heated and accelerated to an even greater extent than in the other calculations. The helium quickly overtakes the protons and the wave acceleration returns to normal, but by this point the helium has been so strongly heated that the pressure force causes the helium to overshoot the critical speed. In this case, however, the helium only reaches $v_\alpha = v_p + 2V_A/3$ at 1 AU and still exhibits $T_\alpha < 4T_p$

there. This result could be pushed further, but it seems unlikely that the helium would lag the protons by as much as 150 km/s at 10 R_S, and the temperatures involved make the approximation for the heating rates of questionable accuracy. Still, the means by which heavy ions are propelled out of the corona are not well known and it is possible that further study of this overshoot phenomenon will prove valuable.

Thus, with the possible exception of this overshoot case, we find that the resonant interaction with ion-cyclotron waves cannot accelerate or heat the solar wind helium population to the observed values with respect to the protons, even in the most favorable case. The fundamental difficulty is that the cyclotron resonance also acts on the solar wind protons, and the preferential interaction with helium is simply not preferential enough.

Equation (26) also predicts that other heavy ions will find limiting differential speeds which depend on their mass-per-charge ratio. This is verified by the calculations displayed in Figures 5 and 6 for oxygen and iron ions, respectively. In these calculations, the proton and wave parameters were input as for the benchmark case (23), but the heavy ion abundance ratios were lowered to values comparable to those given by Bame, et al. (1979) and Mitchell, et al. (1981): $n_O/n_p = 5 \times 10^{-4}$ and $n_{Fe}/n_p = 5 \times 10^{-5}$. The calculation for each ion charge state was made separately and these abundances were taken for each case. At these densities, the oxygen and iron ions are essentially test particles, though in principle they are still allowed to affect the protons. We also caution the reader that the oxygen and iron results do not include helium, so the heavy ions are possibly given more energy than a full multi-ion calculation would yield. Thus, the velocities and temperatures shown in Figures 5 and 6 should perhaps be considered as upper limits. The results of these calculations at 60 R_S and 1 AU are also shown in Table 1, along with the benchmark results for He^{++}.

The two most abundant oxygen ions, O^{+6} and O^{+7}, have mass-per-charge ratios only slightly larger than He^{++}. Consequently, the velocities in Figure 5a are only slightly higher than the helium benchmark flow velocities. As expected from (26), the O^{+6} speeds are also marginally higher than those for O^{+7}. The oxygen temperature ratios are significantly lower at 1 AU than the $T_O/T_p \approx 16$ which is often observed, and the ratios display the same type of anticorrelation with proton speed as shown in Figure 3 for helium. These results are in conflict with the observations of oxygen in high speed solar wind (Ogilvie, et al., 1980, 1982; Schmidt, et al., 1980) which show $v_O \approx v_\alpha > v_p$, $T_O/T_p \approx 16$, and the same positive correlation of T_O/T_p with v_p as seen for helium.

The iron ions have a higher critical differential speed (26), and indeed, they do reach speeds of $v_p + V_A$ in this calculation. However, the ratios T_{Fe}/T_p fall well below the mass proportionality level that is observed (Mitchell, et al., 1981; Bochsler and Geiss, 1982), and the iron temperature ratios are also anticorrelated with v_p.

We note, finally, that the temperature ratios with respect to helium, T_h/T_α, are more closely matched to the mass ratios than the equivalent values with respect to the protons (see Table 1). This is an intriguing result, indicating an intrinsic tendancy to mass-proportional heating by wave dissipation. It can be shown through an analysis similar to that used to derive the critical differential speed (26), that if the waves accelerate all the heavy ions to the same speed, the ions will be heated in proportion to their masses. This result is independent of the expression for the dissipative heating, Q_i, requiring only that the acceleration and heating be related by (24) and that the interaction take place quickly. In this picture, however, the protons, essentially carrying the waves, are not accelerated to the same degree as the other ions and so do not enter into this mass-proportionality as we see from our model. Of course, in our model the ions are not all accelerated to the same speed, but rather to their critical speeds which depend on A_i/Z_i, so the mass-proportionality is not exact. However, we can expect that a wave dissipation process which does accelerate all heavy ions to the same speed, as is observed (Ogilvie, et al., 1982) will necessarily produce the observed mass-proportional temperatures as well.

In this study, we have concentrated on the interaction with left-polarized Alfvén waves. A recent calculation by Marsch, et al. (1982c) has included effects of an interaction with right-polarized magnetosonic waves, and they find that their model can produce $v_\alpha \simeq v_p + V_A$ inside 30 R_S if they hypothesize enough power in the waves.

There are a number of differences between their model and the work presented here. The model of Marsch, et al. (1982c) allows for anisotropic (bi-Maxwellian) particle distributions, and uses dispersive waves with a phase speed in the plasma frame of

$$\omega/k_{\parallel} = V_A \left\{ (1 \mp \omega/\Omega_p) \left[1 + 4\pi \; (p_\perp - p_\parallel)/B_0^2 \right] \right\}^{\frac{1}{2}} \tag{27}$$

where the minus (plus) sign refers to left- (right-) polarized waves and $p_\perp$, $p_\parallel$ are the proton pressures perpendicular to and parallel to the average field. They do not include a process equivalent to our saturation-and-cascade scenario, so the required wave power at resonant frequencies is much larger than we have used. As a consequence of starting the calculation with large amounts of resonant wave power, the plasma is accelerated and heated almost immediately, over a distance of several hundredths of a solar radius rather

rather than the 1-10 R_S interaction length we obtain.

Since they use dispersive waves, Marsch, et al. (1982c) can only accelerate helium ions to differential speeds $\Delta v_{\alpha p} \lesssim 0.2\ V_A$ by the interaction with left-polarized ion-cyclotron waves, verifying statements to that effect in DH, McKenzie and Marsch (1982) and Section 2 of this paper. Subsequently, the plasma flows into regions of decreasing Alfvén speed while the distribution functions evolve in the expanding heliospheric field. If the helium has been heated enough in the parallel direction, and if the Alfvén speed decreases enough, the faster helium ions will move into cyclotron resonance with any right-polarized MHD waves that are present. These waves can further heat and accelerate the resonant ions. With equal intensities of right- and left-polarized waves, Marsch, et al.(1982c) produce results with $\Delta v_{\alpha p} \gtrsim V_A$ well within 30 R_S, apparently demonstrating that cyclotron resonant interactions can explain the Helios observations of helium.

However, there may be a problem with this interpretation. Inserting the phase speeds (27) into (2), one obtains expressions for the parallel velocity needed by an ion in order to cyclotron resonate with waves of either polarization. These two expressions are never equal: the maximum velocity for left-polarized resonance is always less than the minimum velocity for right-polarized resonance. For a helium ion, this "resonance gap" is very large, being always greater than $2V_A \left[1 + 4\pi(p_\perp - p_\parallel)/B_o^2\right]^{\frac{1}{2}}$. At 30 R_S, for isotropic distributions, this gap can be greater than 600 km/s, corresponding to the thermal speed of a helium gas with a parallel temperature of $T_\parallel > 174$ million degrees. (These numbers are even larger if $p_\perp > p_\parallel$, as would be expected from cyclotron heating by left-polarized waves.) A particle moving with a parallel velocity in this gap is not in resonance with either mode, so it is clear that the resonant interaction with the left-polarized waves cannot directly produce particles in resonance with right-polarized waves. If particles traveling at the maximum velocity for left-polarized resonance subsequently move into a region of low enough Alfvén speed, they can become resonant with the right-polarized waves, but Alfvén speeds this low are typically found only beyond 1 AU.

It is therefore puzzling that Marsch, et al. (1982c) find such a strong resonance with right-polarized waves inside 30 R_S, where the Alfvén speed is so large. We suspect that their assumption of a Maxwellian distribution in parallel velocity is responsible. This assumption insures that there are at least a small number of particles in resonance with the right-polarized waves at the start of the calculation. When these few particles are accelerated and heated by the right-polarized waves, they are interpreted as representing a widening of the entire velocity distribution and this widening is interpreted as an increase in the number of resonant particles. Thus, the assumption that the distribution remains Maxwellian can artifically propagate particles across the resonance gap. Eventually, the entire distribution may be propelled to high speeds by this numerically induced process. This picture is only a conjecture, as yet without proof, but we feel this point needs further study.

V Summary and Conclusions

We have investigated numerical solutions to the spherically symmetric, steady-state solar wind equations including acceleration and heating through a resonant interaction with left-polarized waves. The equations for the resonant interaction model a scenario which we feel is a "most favorable case" for preferential acceleration and heating of heavy ions. Briefly, we assume dispersionless, parallel-propagating high-frequency Alfvén waves which saturate when the magnetic fluctuations become comparable with the background field. The energy dissipated in this saturation process is assumed to cascade to resonant wavenumbers where it is picked up by the plasma through the quasi-linear resonant cyclotron interaction.

We find that, even in this most favorable case, the resonant cyclotron interaction is not capable of accelerating and heating the alpha particle population in the solar wind to the observed speeds and temperatures. We first point out that the quasi-linear interaction does not yield preferential heating or acceleration for heavy ions unless the wave power spectrum is steeper than k^{-2}. Wave spectra this steep are rarely seen in the solar wind. Moreover, even if sufficiently steep spectra are assumed, we find that the resonant process in the case of helium, though acting preferentially on the helium, still accelerates protons to such an extent that the observed differential speeds, $\Delta v_{\alpha p} \approx V_A$, are not produced. We have succeeded in understanding this result analytically, in terms of a critical differential speed (Equation 26) for preferential acceleration. We also find that this most favorable case is unable to explain either the values of T_α/T_p, or the T_α/T_p - v_p correlation, which is frequently observed in high speed streams.

The critical differential speed depends on the mass-per-charge ratio of the ion. Solar wind oxygen (primarily O^{+6} and O^{+7}), with similar ratios to that of He^{++}, is affected in a similar manner to the helium. Iron ions in the solar wind have higher mass-per-charge ratios, and therefore reach higher differential speeds than helium or oxygen, if the wave spectra are steep enough to produce any preferential acceleration. This process is not able to produce the observed mass-proportional temperature ratios, $T_h/T_p \approx A_h$ at 1 AU. However, some tendency towards mass-proportional heating is found among the heavy ions (excluding protons), yielding $T_h/T_\alpha \approx A_h/A_\alpha$ at 1 AU.

In conclusion, we have shown that even a "most favorable case" model of the resonant cyclotron interaction with left-hand waves is not capable of producing the observed preferential acceleration and heating of heavy ions in the solar wind.

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Figure Captions

Fig. 1 Benchmark case for helium. (a) Radial flow speeds of protons (thick line) and alpha particles (thin line). Dashed line is the radial component of $\underset{\sim}{v}_p + \underset{\sim}{V}_A$. (b) Temperatures of protons (thick line), alpha particles (thin line), and electrons (dashed line).

Fig. 2 Proton radial speed at 1 AU and differential speeds at 60 R_S and 1 AU as functions of the initial wave power. Vertical dashed line indicates the location of the benchmark case.

Fig. 3 Ion temperatures and temperature ratio as functions of proton radial speed at 1 AU, obtained by varying the initial wave power. Vertical dashed line indicates the location of the benchmark case.

Fig. 4 Overshoot case for helium. Radial flow speeds (a) and temperatures (b) as in Fig. 1.

Fig. 5 Oxygen. Radial flow speeds (a) and temperatures (b) for protons, O^{+6}, and O^{+7}.

Fig. 6 Iron. Radial flow speeds (a) and temperatures (b) for protons, Fe^{+8}, Fe^{+10}, and Fe^{+12}.

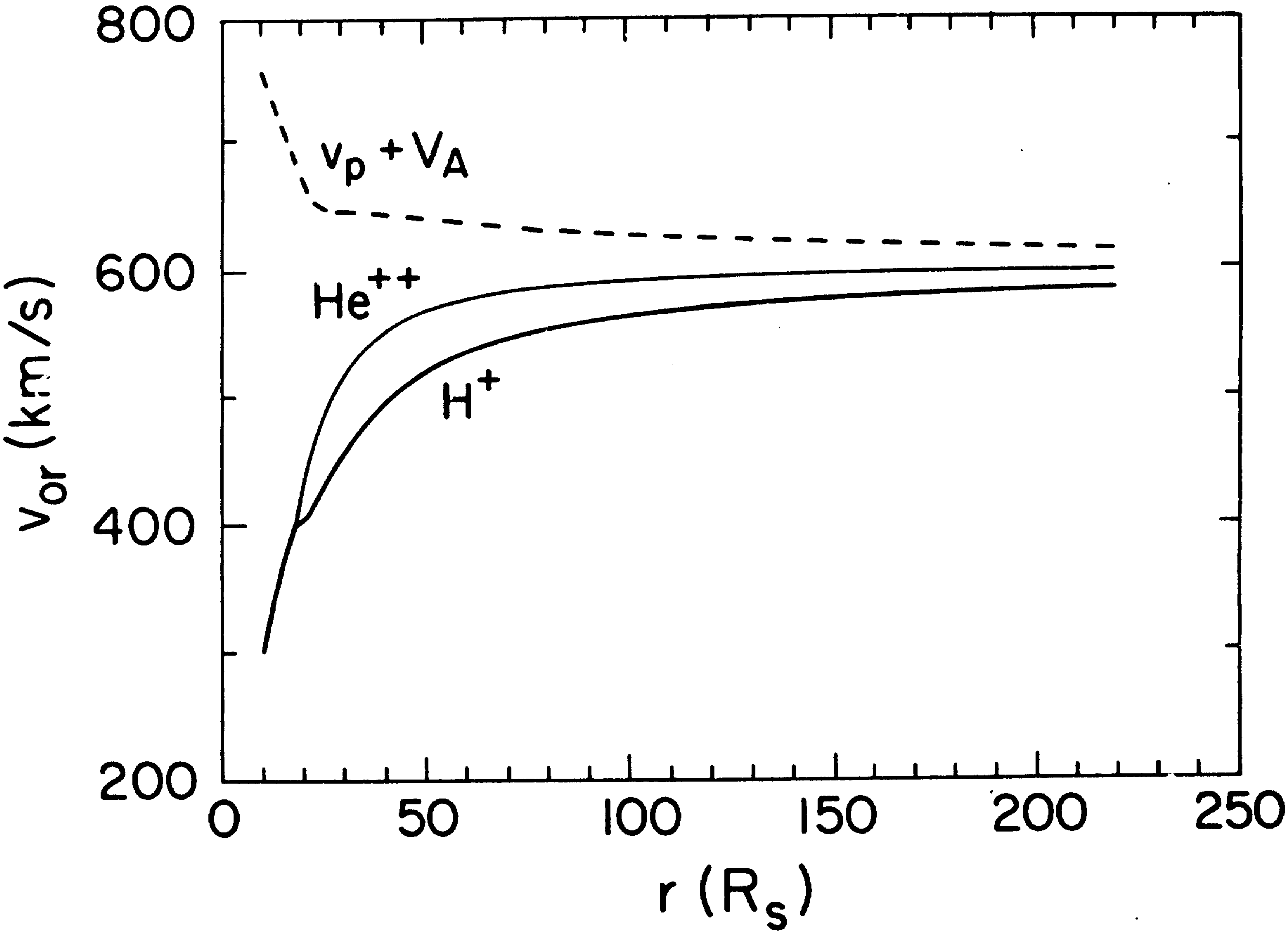

Fig 1a

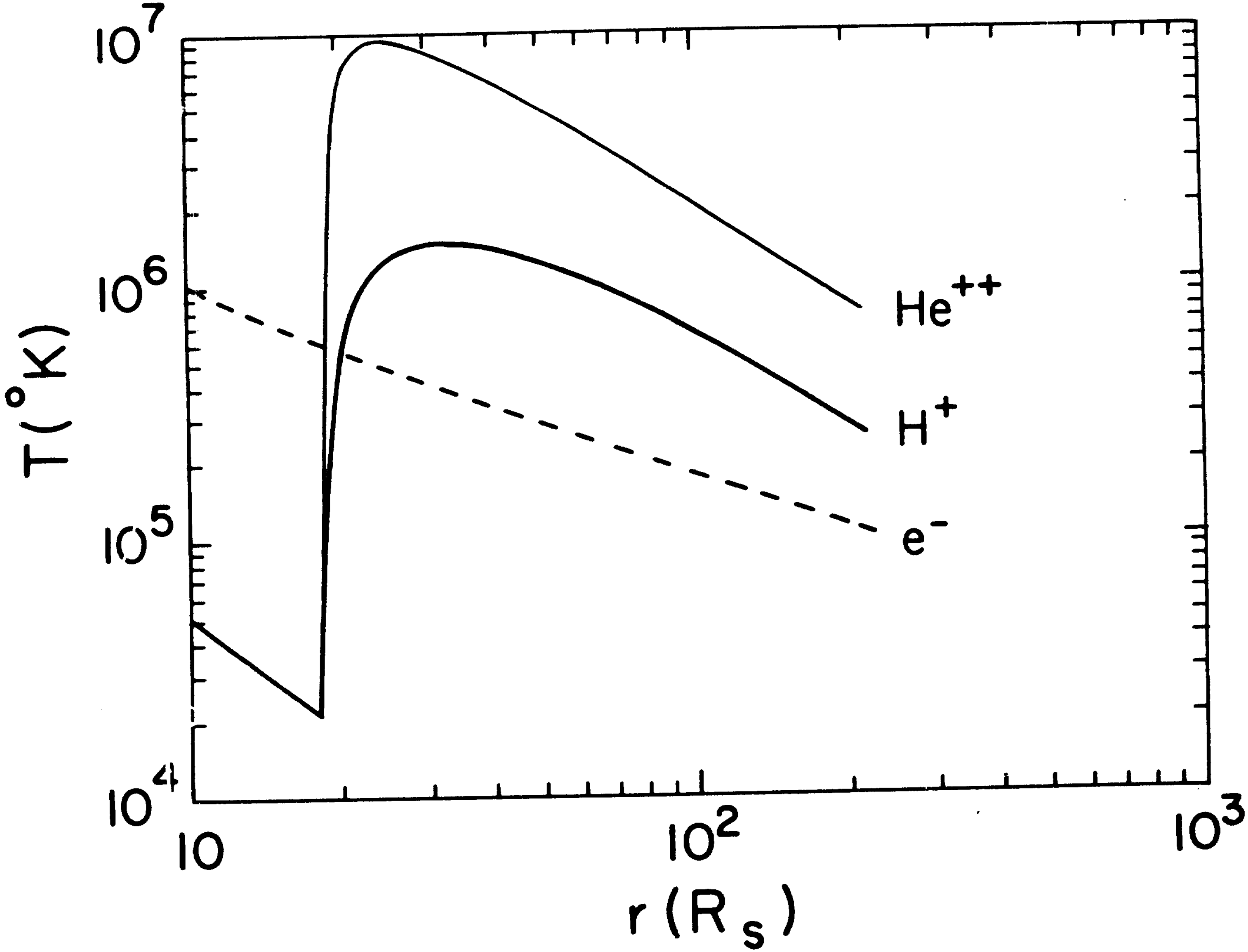

Fig 16

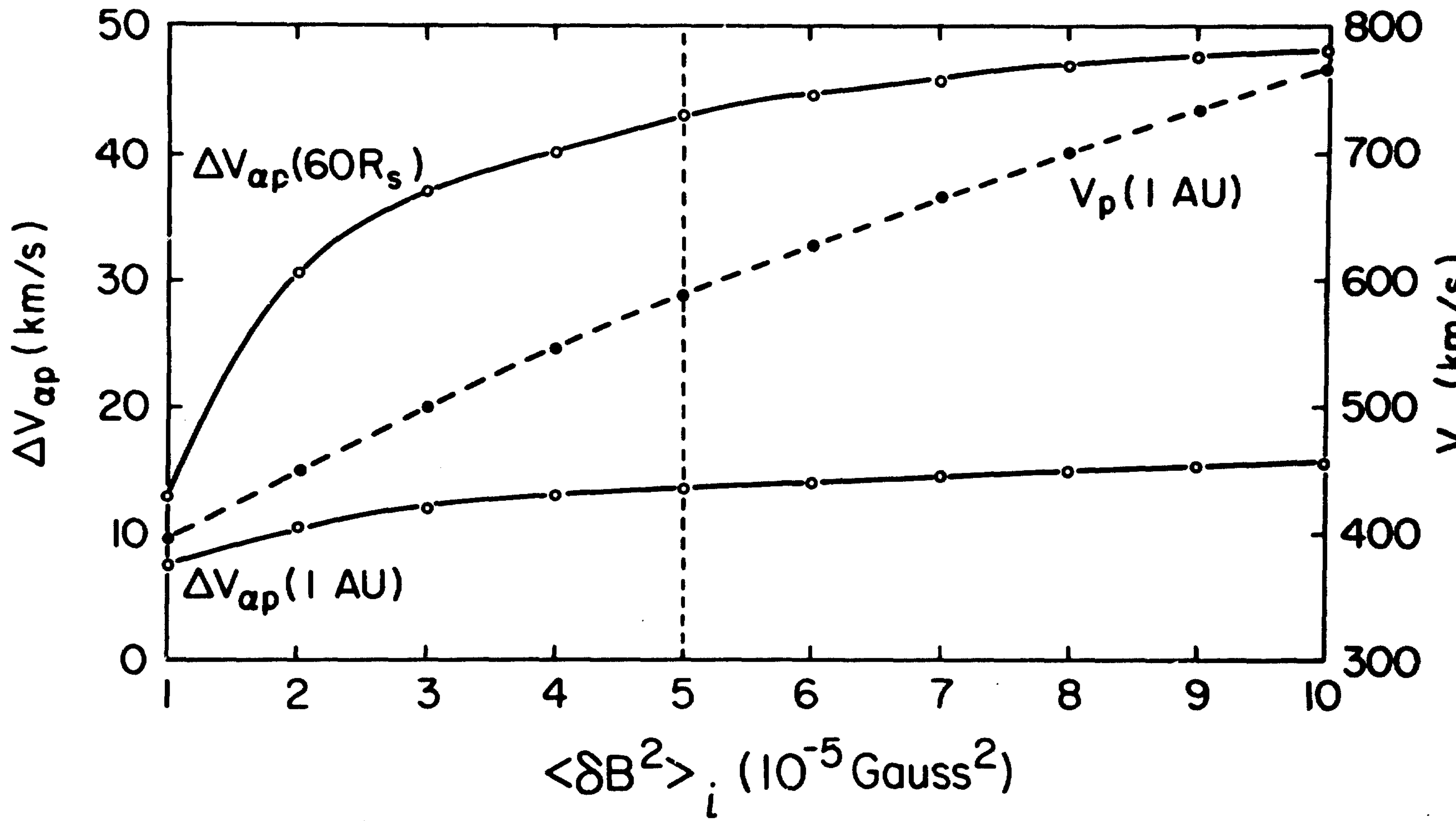
$\Delta V_{\alpha p}$ (km/s)
V_p (km/s)
$\Delta V_{\alpha p}(60R_s)$
V_p(1 AU)
$\Delta V_{\alpha p}$(1 AU)
$<\delta B^2>_i$ (10^{-5} Gauss^2)
0
10
20
30
40
50
300
400
500
600
700
800
1
2
3
4
5
6
7
8
9
10

Fig 2

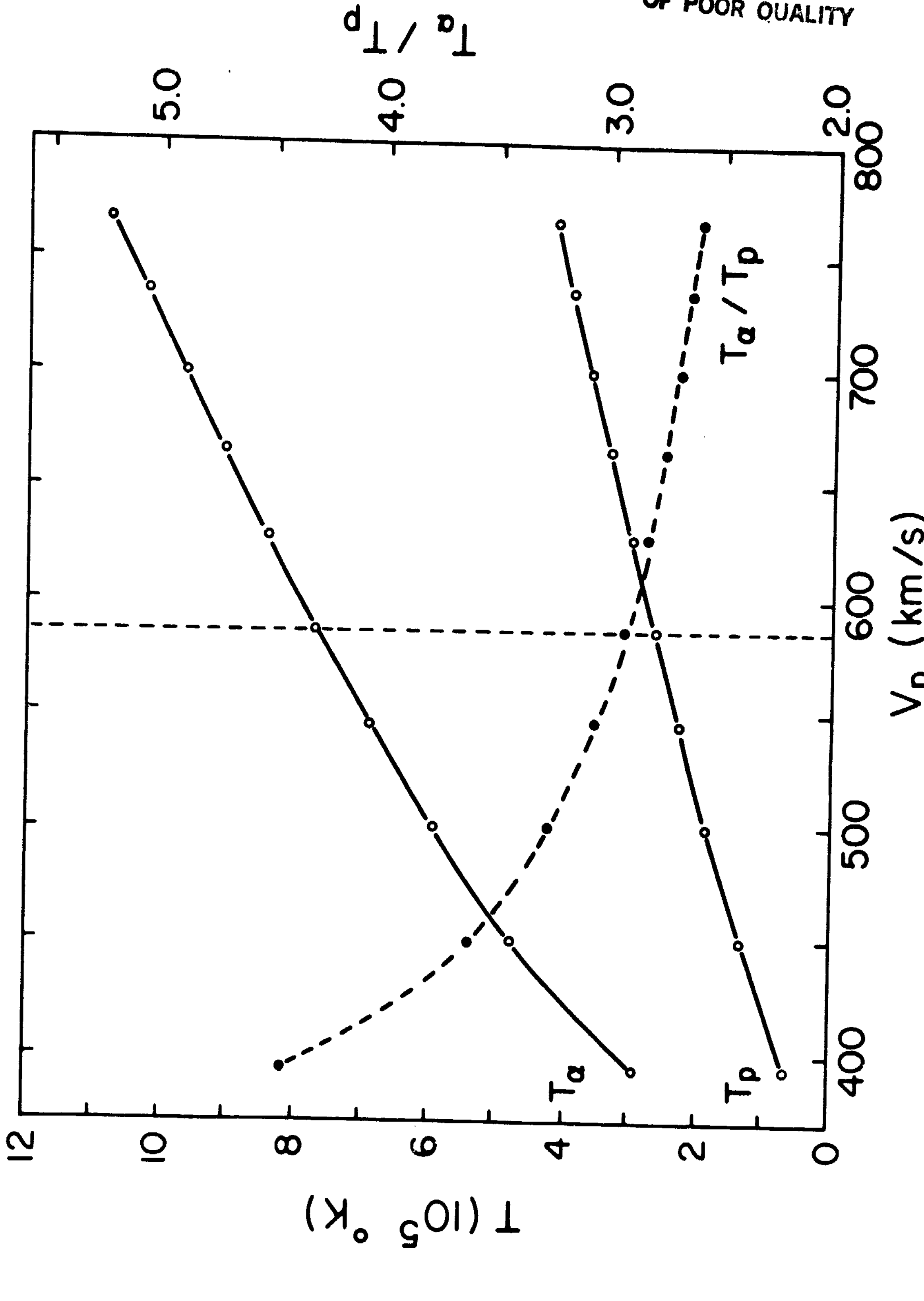
T_α / T_p
5.0
4.0
3.0
2.0
800
700
600
500
400
T_α / T_p
T_α
T_p
12
10
8
6
4
2
0
T (10^5 °K)
V_p (km/s)

FIG 3

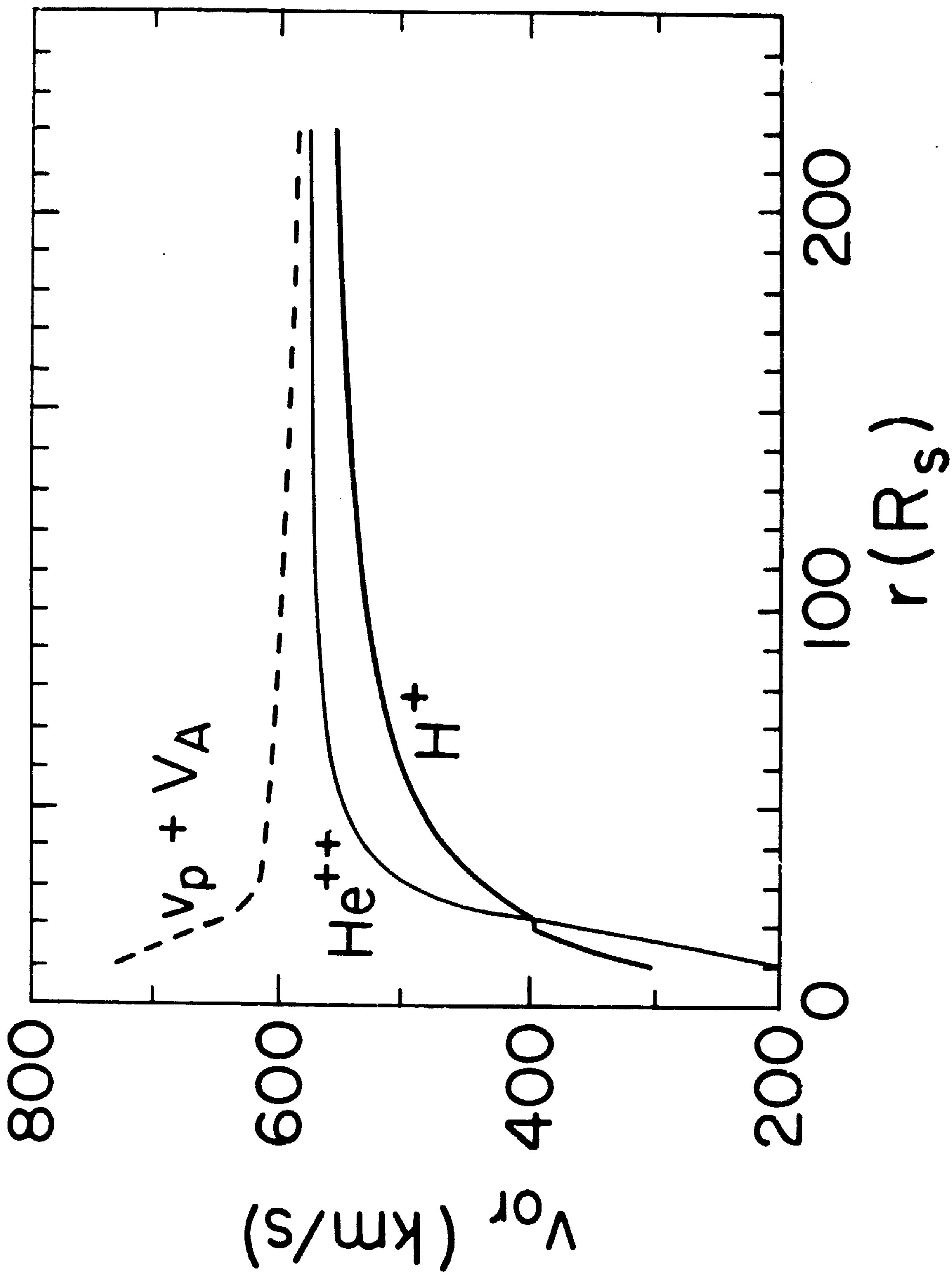

FIG 4a

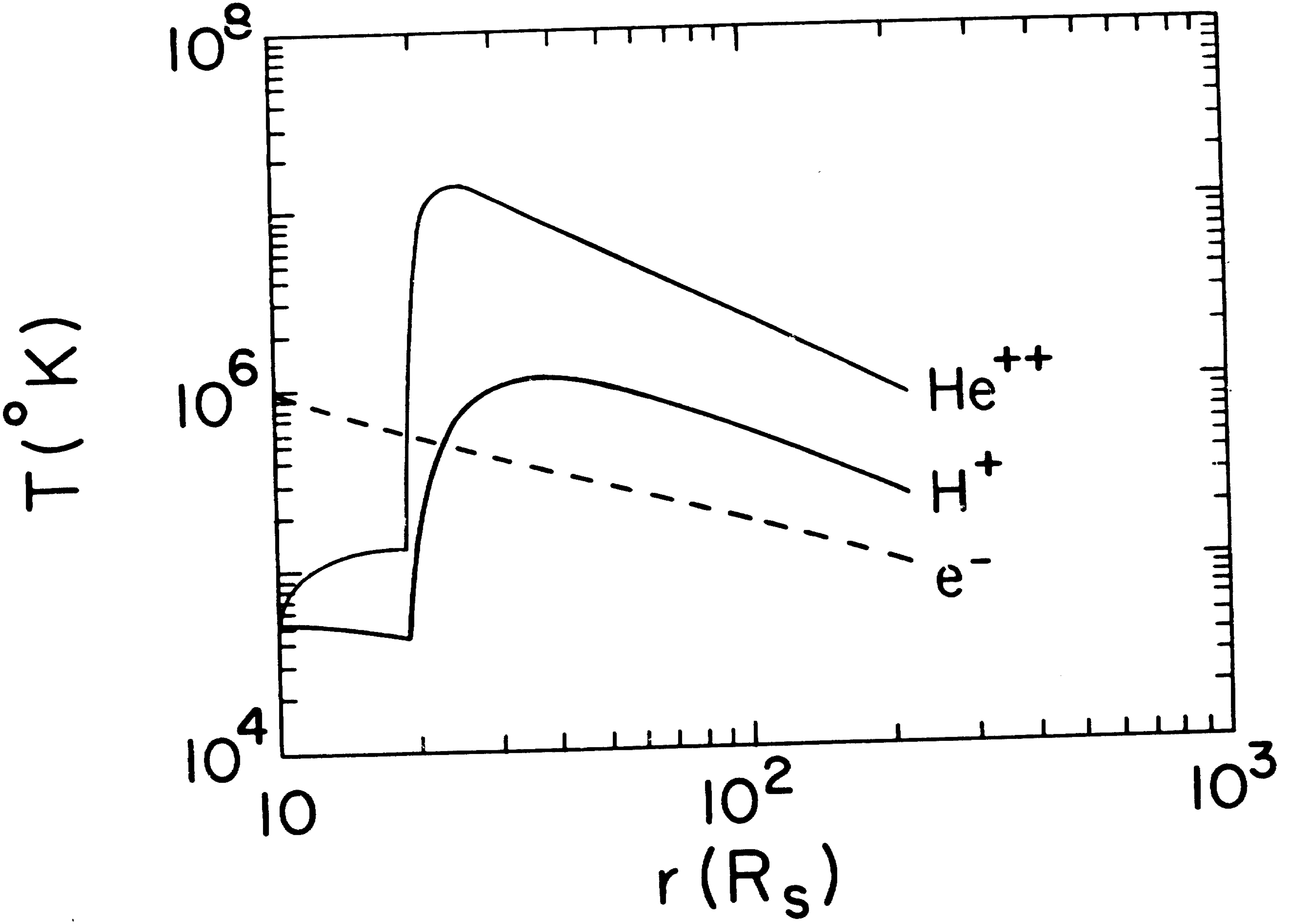

Fig 4b

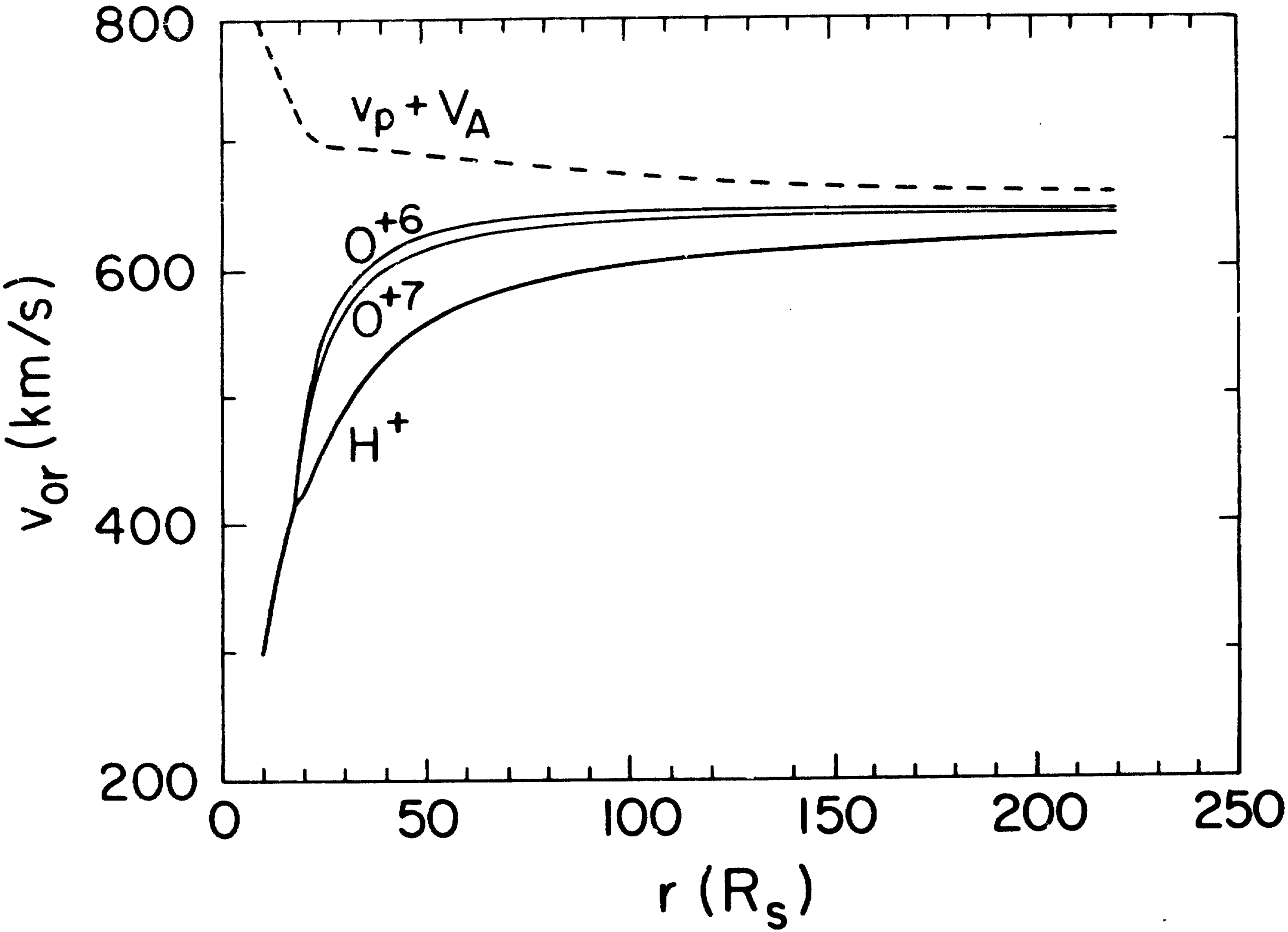

Fig. 5a

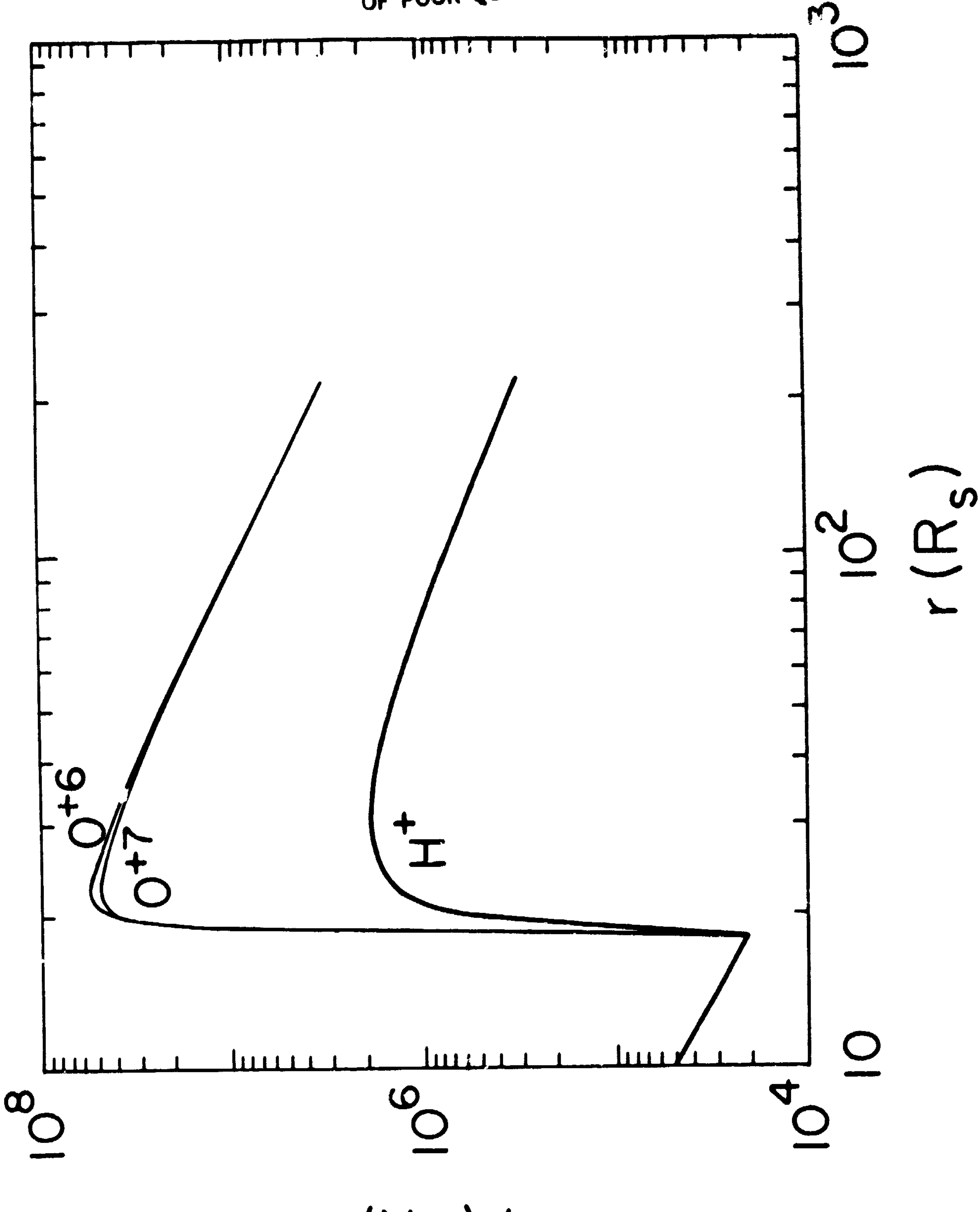

Fig 5b

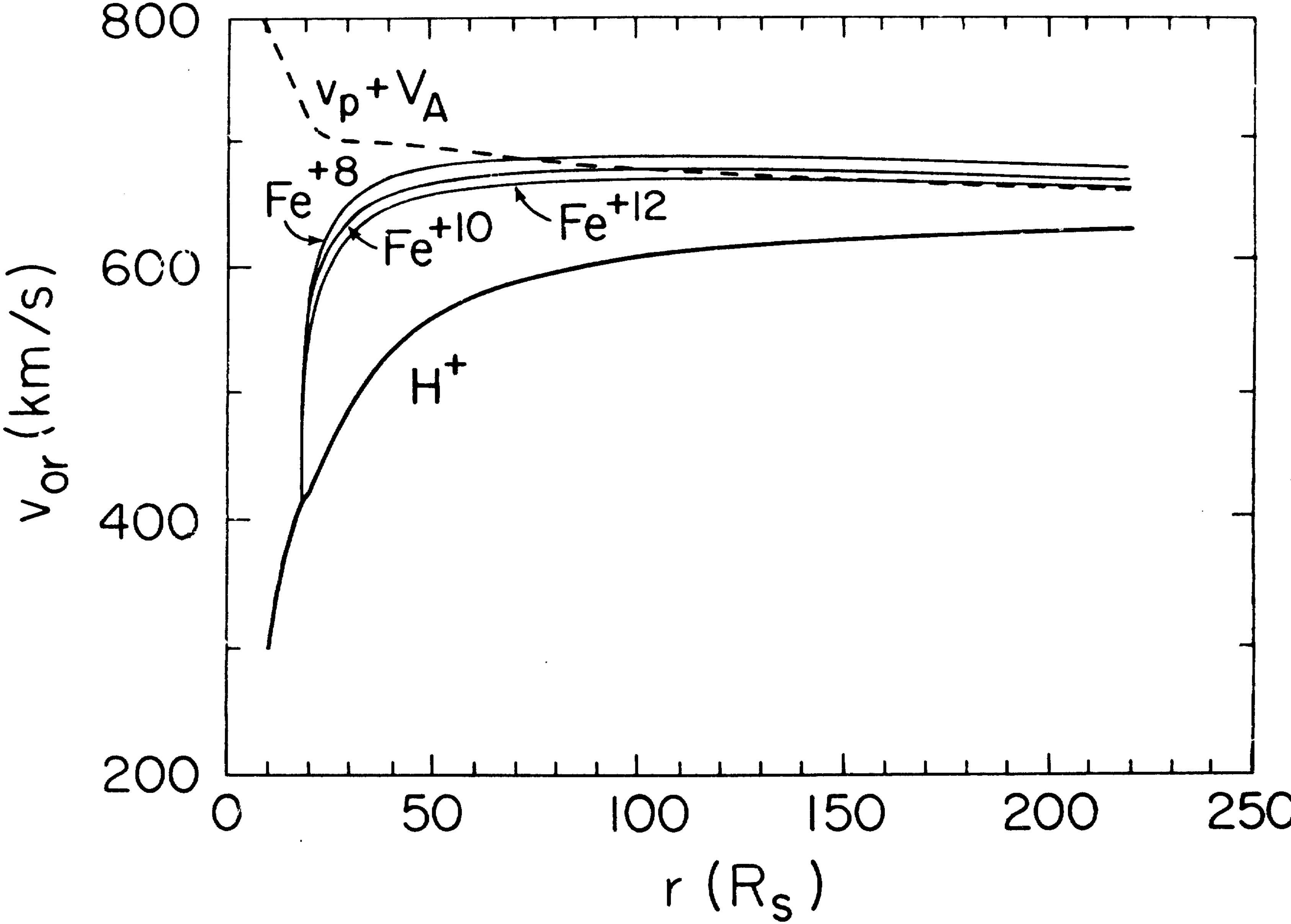
800
600
400
200
0
50
100
150
200
250
v_{or} (km/s)
r (R_s)
$v_p + V_A$
Fe^{+8}
Fe^{+10}
Fe^{+12}
H^+

Fig 6a

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OF POOR QUALITY

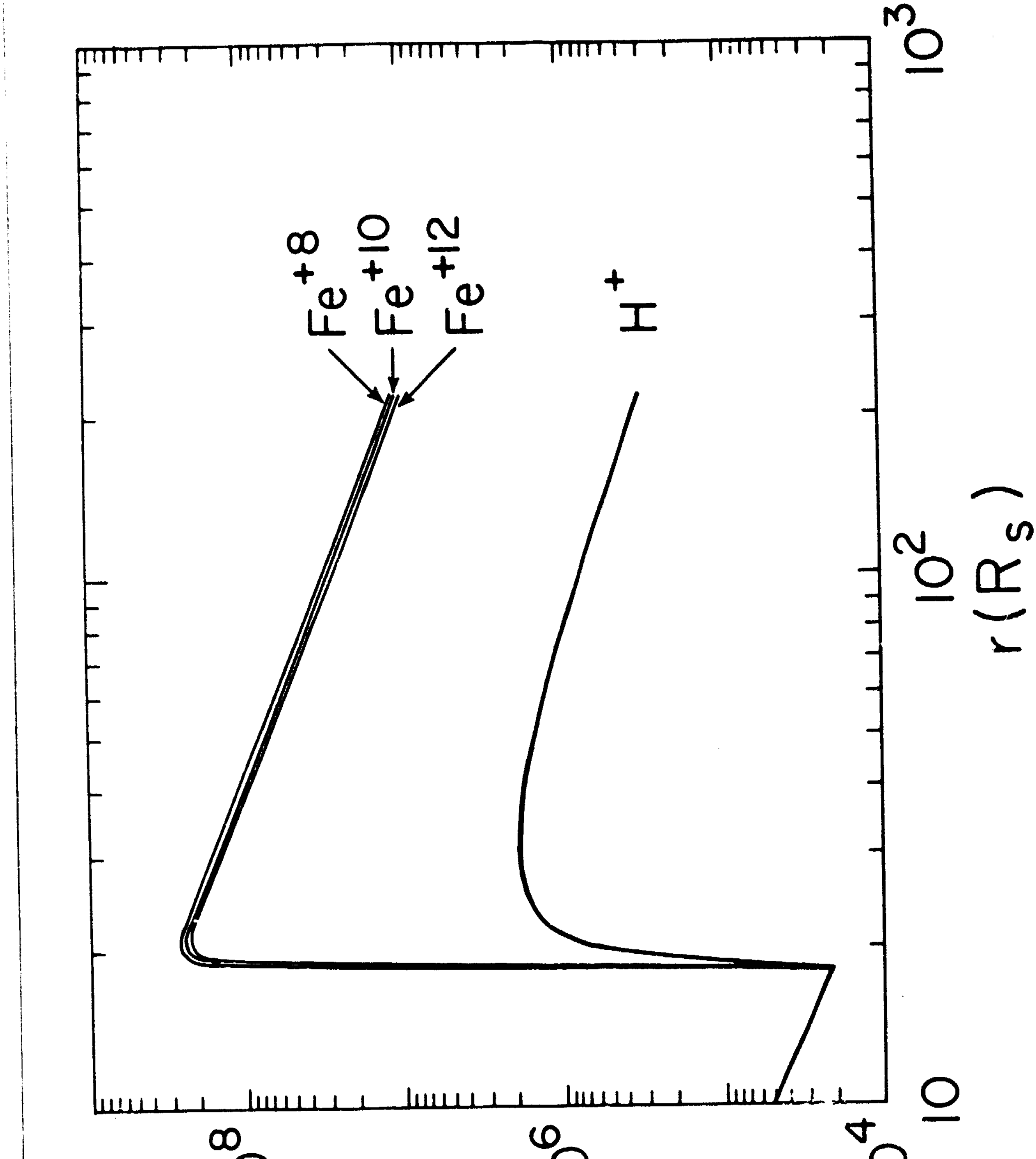

Fig 6b

TABLE I

Numerical Results for Various Ions

	$\Delta v_{hp}/V_A$		T_h/T_p		T_h/T_α	v_p(1 AU)	T_p(1 AU)	V_A(1 AU)
	60 R_S	1 AU	60 R_S	1 AU	1 AU	(km/s)	(10^5 °K)	(km/s)
He^{++}	.418	.452	3.68	2.92	-	587.	2.63	29.9
O^{+6}	.543	.616	13.9	10.5	4.21	629.	3.09	33.1
O^{+7}	.470	.517	14.0	10.7	4.29	629.	3.08	33.1
Fe^{+8}	.937	1.41	49.8	37.1	14.9	631.	3.10	33.3
Fe^{+10}	.847	1.19	48.8	35.9	14.5	631.	3.10	33.3
Fe^{+12}	.775	1.02	48.1	35.0	14.1	631.	3.10	33.2