A ZONE OF PREFERENTIAL ION HEATING EXTENDS TENS OF SOLAR RADIi FROM SUN

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ABSTRACT

The extreme temperatures and non-thermal nature of the solar corona and solar wind arise from an unidentified physical mechanism that preferentially heats certain ion species relative to others. Spectroscopic indicators of unequal temperatures commence within a fraction of a solar radius above the surface of the Sun, but the outer reach of this mechanism has yet to be determined. Here we present an empirical procedure for combining interplanetary solar wind measurements and a modeled energy equation including Coulomb relaxation to solve for the typical outer boundary of this zone of preferential heating. Applied to two decades of observations by the Wind spacecraft, our results are consistent with preferential heating being active in a zone extending from the transition region in the lower corona to an outer boundary 20-40 solar radii from the Sun, producing a steady state super-mass-proportional \( \alpha \)-to-proton temperature ratio of \( 5.2 - 5.3 \). Preferential ion heating continues far beyond the transition region and is important for the evolution of both the outer corona and the solar wind. The outer boundary of this zone is well below the orbits of spacecraft at 1 AU and even closer missions such as Helios and MESSENGER, meaning it is likely that no existing mission has directly observed intense preferential heating, just residual signatures. We predict that Parker Solar Probe will be the first spacecraft with a perihelia sufficiently close to the Sun to pass through the outer boundary, enter the zone of preferential heating, and directly observe the physical mechanism in action.

Subject headings: corona, solar wind, plasmas, turbulence, acceleration of particles, magnetic fields

1. INTRODUCTION

Observations of space over the last half century, including spectroscopic diagnostics of UV emission from coronal plasma and direct in situ sampling of solar wind by spacecraft have shed light on the non-thermal nature of heating in the corona and solar wind. Throughout the heliosphere, plasma is typically found in states other than local thermodynamic equilibrium, with relative drifts and unequal temperatures between species and anisotropic and otherwise non-Maxwellian velocity distribution functions commonly observed. Such non-thermal structure is indicative of mechanisms that selectively
couple to particles with particular velocities, charges or masses and preferentially heat different plasma species. One region in particular where our understanding of these mechanisms is incomplete is the inner heliosphere.

The visible 6000 K photosphere of the Sun is surrounded by a 1 – 10 MK solar corona that reaches many solar radii ($R_s$) into space before transitioning into the supersonic and ultimately super-Alfvénic solar wind. The temperature of the solar atmosphere rises to $10^5$ K within several hundred kilometers in the narrow transition region at the base of the corona. At around 0.1 – 0.3$R_s$, rising temperatures $T$ and falling densities $n$ are such that the frequency of Coulomb collisions, $\nu_c \propto n/T^{3/2}$, drops to the point that the coronal plasma becomes effectively collisionless, with electrons and individual ion species not persisting in a common local thermodynamic equilibrium. Ions become much hotter than electrons, and heavier ions achieve higher temperatures than the hydrogen that composes the majority of the coronal plasma (Esser et al. 1999; Landi & Cranmer 2009). Emission has been detected from steady non-flare coronal oxygen at $\sim 10^8$ K, a hundred times hotter than coronal hydrogen and more than six times hotter than the core of the Sun. Such unequal temperatures serve as a signature of preferential heating of different species in the corona. It is possible that preferential heating is occurring lower in the solar atmosphere, but the higher frequency of Coulomb collisions at lower heights would remove the signature of such heating. The relative temperatures of ion species are highly variable, with a statistical preference for either equal temperatures or equal thermal speeds corresponding to mass-proportional temperatures, but intermediate temperatures and super-mass-proportional temperatures are also observed. For example, a recent study suggested that coronal ions develop an equilibrium temperature $T_i/T_p \approx (4/3)m_i/m_p$ (Tracy et al. 2016). One of the most significant open challenges in solar and space physics is to unambiguously determine the physical processes responsible for this heating.

There are many plausible theories for the physical processes responsible for the extended and preferential heating of different ion species in the corona and solar wind, including resonant wave heating (Cranmer 2000; Hollweg & Isenberg 2002), velocity filtration (Scudder 1992), impulsive events including reconnection (Cargill & Klimchuk 2004; Drake et al. 2009), and stochastic heating by low-frequency Alfvénic turbulence (Chandran et al. 2010; Chandran 2010). Unambiguous identification of the dominant process is complicated by uncertainty in the nature of energy readily available for dissipation in the corona. For example, high frequency waves could escape from the photosphere through the transition region before being damped in the lower corona (Axford & McKenzie 1997). Alternately, MHD turbulence could be generated locally by the interaction between outward and reflected low frequency waves anywhere below the solar wind Alfvén point (Matthaeus et al. 1999). Recent work (Kasper et al. 2007; Chandran et al. 2013; Kasper et al. 2013) has shown that velocity moments of solar wind $H^+$ and $He^{2+}$ ions are consistent with both strong heating due to resonant absorption of Alfvén-cyclotron waves or stochastic heating. This heating could persist throughout the heliosphere or could occur only in a select region near the Sun, with the resultant non-thermal structure being reduced by infrequent Coulomb collisions as the solar wind expands (Kasper et al. 2008; Tracy et al. 2016).

In situ measurements over the last half-century, including those from the twin Helios spacecraft that approached to within 62$R_S$ from the Sun, show that the radial gradients of ion and electron temperatures are much more shallow than would be expected from a cooling solar wind undergoing adiabatic expansion (Hellinger et al. 2013). This evidence for ongoing radial heating in the inner heliosphere is not necessarily evidence for ongoing preferential ion heating of the type observed near the Sun. More detailed tests involving the correlation of particle distribution function structure and electromagnetic fields may be able to identify the energy source and distinguish between the proposed mechanisms (Klein & Howes 2016), but such tests need measurements of the plasma as the heating is occurring. It is therefore important to determine how far away from the Sun the preferential heating mechanism is active, and how close to the Sun a spacecraft must approach to directly resolve the process.

In this paper, we address three related questions: Are unequal temperatures in the solar wind maintained by ongoing local preferential heating, or are they a leftover of heating that happened close to the Sun? Is faster solar wind further from local thermodynamic equilibrium than slow wind because only fast wind experiences preferential heating in the corona resulting in non-thermal structure? How far from the Sun does preferential ion heating continue? The purpose of this paper is to develop a technique for measuring how much time has elapsed since solar wind ions experienced preferential heating that was sufficiently strong to generate super-mass-proportional temperatures. We assume that there is a zone of preferential heating surrounding the Sun, starting at 0.1 – 0.3 $R_S$ as indicated by spectroscopic observations, and ending at an outer boundary $R_b$. Beyond $R_b$ any remaining non-preferential heating is weak and Coulomb collision dominates, leading to an exponential decay of the temperature ratio $T_i/T_p$ toward unity. Observational motivation for this work is presented in Section 2, with a model for the radial evolution of the temperature ratio between $H^+$ and $He^{2+}$ detailed in Section 3. The technique for measuring the outer boundary of the zone combining the derived model and in situ measurements from the Wind spacecraft is presented in Section 4. Values for the outer boundary, Section 5, are...
found to be within the perihelion of Parker Solar Probe, allowing for verification or falsification of our model and, potentially, the first in situ observation of the preferential heating mechanism.

2. OBSERVATIONS OF COULOMB THERMALIZATION

The observational basis of this work is an extensive set of measurements of solar wind plasma collected by the Solar Wind Experiment (SWE, Ogilvie et al. (1995)) and the Magnetic Field Investigation (MFI, Lepping et al. (1995)) instruments on the NASA Wind spacecraft. Wind was launched in late 1994 and has operated continuously in a variety of orbits passing through the solar wind near Earth, resulting in a comprehensive set of observations of solar wind in the ecliptic plane spanning nearly two decades. Solar wind H\(^+\) (protons) and He\(^{2+}\) (α particles) are measured by the SWE Faraday Cup instruments, which record a detailed three-dimensional measurement of the velocity distribution function (VDF) of the two ion species once every 90 seconds. We use a technique developed to extract anisotropic temperatures and differential flows for each species as first described by Kasper et al. (2002). This algorithm makes use of 3-second time resolution measurements of the vector magnetic field by the MFI flux gate magnetometers in order to determine the temperature of each ion species parallel and perpendicular to the local magnetic field. Approximate uncertainties in the resulting observations were documented in Kasper et al. (2006), which estimated a typical uncertainty in an ion temperature measurement of 8%. We follow the same data selection procedures described in Kasper et al. (2008), but with an additional 8 years of observations.

Previous work (Feldman et al. 1974; Neugebauer 1976; Livi et al. 1986; Kasper et al. 2008; Tracy et al. 2016) has demonstrated that Coulomb relaxation plays an important role in thermalizing the solar wind ions, removing non-thermal structure such as temperature anisotropy and temperature disequilibrium between species. The effect of this thermalization can be quantified by the estimated number of Coulomb thermalization times that have elapsed in the time it takes for the solar wind to expand from the corona to the observing spacecraft, a quantity often referred to as the Coulomb collisional age \(A_c\) (Salem et al. 2003; Kasper et al. 2008; Maruca et al. 2013). We will reserve \(A_c\) for a more precise calculation presented later in this paper, and introduce the Coulomb number \(N_C = \nu_{ab} R / U\) to indicate a rough approximation based only on observations at a spacecraft in interplanetary space with no accounting for propagation effects. In this equation \(R\) is the total distance of the spacecraft from the Sun, \(U\) is the speed of the solar wind, and \(\nu_{ab}\) is the characteristic rate for Coulomb interactions between two species \(a\) and \(b\); for \(N_C\) both \(\nu_{ab}\) and \(U\) are assumed to be constant. Throughout this paper we make use of the calculations of Hernandez et al. (1987) for the Coulomb interaction between two species with Maxwellian distribution functions, different temperatures, densities, and a net differential flow as discussed in more detail in the following section.

Column normalized distributions of three markers of non-thermal structure, \(T_{l,p}/T_{l,p}\), \(T_{\alpha}/T_{\alpha}\), and the α-proton drift velocity normalized by the Alfvén speed \(\Delta V_{\alpha p}/C_A\), are plotted as a function of solar wind speed \(U\) in the left panels of Fig. 1. As has been reported many times before, the fast solar wind is more non-thermal than slow solar wind (c.f. Marsch (2012)). In the right three panels, the same markers are plotted as a function of \(N_C\). As previously reported in Kasper et al. (2008), \(N_C\) is shown to also order the observations of these three non-thermal measures. The magnitudes of the properties are observed to decrease exponentially with both \(U\) and \(N_C\) for sufficiently large \(N_C\). For instance, the dependence of \(T_{\alpha}/T_{\alpha}\) on \(N_C\) is monotonic, with a single peak value of the temperature ratio for each value of \(N_C\), which decreases exponentially with large Coulomb number and a normal distribution of temperature ratios about that peak. The same cannot be said of the dependence of \(T_{\alpha}/T_{p}\) on \(U\); while they are also strongly correlated, the spread in each \(T_{\alpha}/T_{p}\) is larger and has multiple peaks for a given value of \(U\). In general, the distribution of \(T_{\alpha}/T_{p}\) for a given \(U\) is further from a normal distribution than the distribution for a given \(N_C\). The variation of \(T_{l,p}/T_{l,p}\) and \(\Delta V_{\alpha p}/C_A\) are more complex, possibly because they are more sensitive to kinetic microinstabilities and the effects of expansion, but even so these non-thermal features are washed away at sufficiently large \(N_C\).

The exponential dependence of \(T_{\alpha}/T_{p}\) on \(N_C\) is consistent with a simple model for the radial evolution of the temperature ratio. Considering the thermalization of temperature differences in the absence of any effects other than Coulomb collisions, keeping \(T_p\) constant, and following Spitzer (1962), the time evolution of \(T_{\alpha}/T_{p}\) can be written as \(d(T_{\alpha}/T_{p})/dt = -\nu_{\alpha p} T_{\alpha}/T_{p}\), yielding a solution of \(T_{\alpha}/T_{p} \sim \exp\left[ -\int \nu_{\alpha p} dt \right]\). Under the oversimplifying but instructive assumption that \(\nu_{\alpha p}\) is constant, and that the appropriate dynamical time is the transit time from the Sun, allows the further simplification \(T_{\alpha}/T_{p} \sim \exp\left[ -N_C\right]\). This form is in good agreement with the solar wind observations, raising several interesting possibilities. First, the fact that this single formula fits all of the Wind observations across all solar wind speeds suggests that non-unity \(T_{\alpha}/T_{p}\) and preferential ion heating may not be restricted to fast solar wind. Perhaps all solar wind close to the Sun experiences strong preferential ion heating and develops a large \(T_{\alpha}/T_{p}\), and the apparent association between \(T_{\alpha}/T_{p}\) and \(U\) is simply due to the fact that the number of Coulomb collisions a parcel of solar wind experiences varies strongly with \(U\). Slower wind both takes longer to get to the spacecraft and tends to have a significantly higher \(\nu_{\alpha p}\),
resulting in a stronger suppression of non-thermal $T_\alpha/T_p$ which may be present closer to the Sun. One might counter that $T_\alpha/T_p$ and $N_C$ are both strongly correlated with speed or temperature, giving the false impression that $N_C$ regulates $T_\alpha/T_p$. However, Maruca et al. (2013) demonstrated that the temperature ratio has a stronger correlation with the number of Coulomb collisions than other solar wind parameters such as density, speed, and temperature.

To further show that $N_C$ and $T_\alpha/T_p$ are not simply dependent on $U$, we plot in Fig. 2 the mean value of the excess temperature ratio $\epsilon \equiv T_\alpha/T_p - 1$ as a function of both solar wind speed and $N_C$. While there is clearly a trend in the typical $N_C$ as a function of speed, the exponential drop in $\epsilon$ from high values of $\approx 4$ to less than unity happens at all observable speeds at $N_C \sim 0.7$, with slight dependence on speed. Even the slowest solar wind, with speeds less than 300 km s$^{-1}$, has high $\epsilon$ when $N_C$ is small. As captured in Figs. 1 and 2, observations of the solar wind are consistent with a hypothesis that all plasma close to the Sun experiences preferential heating of ions, even plasma that results in slow wind. The association of significant $\epsilon$ with faster wind speeds is simply due to the fact that slower wind in general has a higher collisional age, leading to a removal of the non-thermal structure by the time the plasma reaches 1 AU. This result is significant, because it implies that mechanisms that could produce non-thermal heating may be active both in slow and fast wind. In the following sections we use this theoretical framework to produce an estimate of how far from the Sun this heating occurred.

3. MODELING THE PREFERENTIAL HEATING ZONE
The clear exponential dependence of \( \epsilon \) on \( N_C \) in Figs. 1 and 2 is suggestive of the gradual thermalization due to Coulomb relaxation on a non-thermal plasma. Spitzer (1962) showed that non-thermal plasma relaxes to thermal equilibrium through a series of small-angle scattering of ions mediated by the Coulomb interaction. In the absence of any other processes, two species with a temperature difference \( \Delta T \) will come to equilibrium at a rate \( d\Delta T/dt = -\nu_c \Delta T \). Ignoring any \( T \) dependence in \( \nu_c \), we can rearrange this equation as \( d\Delta T/\Delta T = -\nu_c \), or integrating both sides and exponentiating,

\[
\Delta T = \Delta T_0 e^{-\int \nu_c \, dt}
\]

where we can define \( \Delta T_0 \) as the initial temperature difference and the collisional age \( A_c \) of the plasma as that integral over time of all Coulomb collisions experienced by the plasma since it began to relax

\[
A_c \equiv \int \nu_c \, dt \simeq N_C.
\]

We now develop a more sophisticated model for the behavior of \( T_0/T_p \), which improves upon the assumption that \( \nu_{op} \) is constant and that the correct dynamical time is the transit time from the center of the Sun to Earth at constant speed as used for \( N_c \). Such an approach was used in Maruca et al. (2013), which solved the ion temperature differential equations backwards in time to investigate the distribution of \( \epsilon \) near the Sun, finding that for radial distances of 0.1 AU \( \epsilon \) took on highly non-thermal values for all solar wind speeds. In this paper we do the opposite; we assume that the plasma is highly non-thermal near the Sun, with a large value of \( \epsilon \) below some radial boundary \( R_b \), and that the observed variation in \( \epsilon \) at 1 AU is subsequently determined solely by Coulomb relaxation. Values for \( R_b \) are then obtained from comparing models for radial solar wind behavior with in situ measurements at 1 AU.

We make the following key assumptions in the construction of our model, which are are illustrated schematically in Fig. 3:

- There is a zone in the inner heliosphere where the Coulomb collision frequency is sufficiently low and the ion heating rate, due to unspecified mechanisms, is sufficiently high to allow for preferential heating of ions. Based on spectroscopic observations of ion temperatures in the corona this zone begins just \( 0.2 - 0.3 \ R_s \) above the photosphere, but the outer extent of this zone is unknown.
- The preferential heating results in different ion temperatures, with \( \epsilon \) reaching an asymptotic value \( \epsilon_0 \) within the zone. Here we are motivated by the fact that the observed spread in \( \epsilon \) is very narrow for small \( N_c \).
- We assume that at some distance from the Sun the preferential heating falls off and quickly becomes negligible. We define this outer boundary of the zone as \( R_b \).
- Above \( R_b \), \( \epsilon \) decays exponentially as a function of the number of Coulomb collisions.

We acknowledge that this model makes several critical simplifications, each of which merits further investigation. For example, \( R_b \) may vary with time, solar wind type, or level of solar activity. The preferential heating in practice will not shut off completely at \( R_b \), and it would be worthwhile to investigate the impact of a more gradual evolution. Finally, we know that the steady state \( \epsilon \) in solar wind with low \( A_c \) is a function of other plasma properties, such as differential flow and plasma \( \beta \) (Kasper et al. 2013), and has a non-negligible spread for a given set of parameters. Nonetheless, for the purposes of this paper, where we aim to determine if the mean value of the observed temperature excess can be described using a fixed outer boundary, and differentiate between a boundary in the lower corona, interplanetary space, or somewhere in between, this model is sufficient.

To model the excess temperature, we start with an energy equation for \( T_p \) and \( T_s \)

\[
\frac{dT_s}{dr} = (\gamma - 1) \left[ \frac{T_s \, dn_s}{n_s \, dr} - \frac{Q_s}{n_s k_B U} \right] - \sum_{s'} \frac{\nu_{ss'}}{U} (T_s - T_{s'}),
\]

which includes the effects of expansion, input heating, and collisional relaxation. The Coulomb coupling between the ion species is governed by the frequency of energy-changing collisions between the two species \( \nu_{ss'} \), and the input heating is parameterized by the heat in-
put rate $Q_s$ in ergs s$^{-1}$ cm$^{-3}$. This form of the adiabatic energy equation, found for example in Cranmer et al. (2009), assumes a steady wind with radially dependent speed of $U(r)$. Given this form of radial temperature evolution, and assuming the collisional coupling is dominantly between the protons and $\alpha$ particles, the radial change in $\epsilon$ is

$$\frac{d\epsilon}{dr} = \frac{1}{T_p} \frac{dT_p}{dr} - \frac{T_p}{T_p^2} \frac{dT_p}{dr}$$

$$= \left(\frac{\gamma - 1}{T_p} \left[ \frac{T_p d\epsilon}{d\rho} - \frac{T_p d\rho}{\rho} - \frac{Q_{\alpha}}{T_p n_\alpha k_B U} + \frac{T_p Q_p}{T_p n_p k_B U} \right] \right) - \left(\frac{T_p - T_\rho}{T_p} \left[ \frac{\nu_{\alpha\rho}}{U} + \frac{\nu_{\alpha\rho} T\alpha}{U T_p} \right] \right).$$

(4)

To model the excess temperature ratio beyond $R_b$, we assume that either $Q_s = 0$ for both ion species, or that any remaining heating affects both species equally. This allows us to relate $Q_p$ and $Q_{\alpha}$ by

$$Q_{\alpha} = Q_p \frac{n_\alpha T\alpha}{n_p T_p}$$

(5)

which upon insertion into Eqn 4 allows us to neglect the input heating terms. We further assume that both ion species follow the same radial density profile

$$n_s(r) \propto n_0 r^{-\xi}$$

(6)

leading to the cancellation of the $dn_s/dr$ terms in Eqn. 4. With these two assumptions, we have the simplified expression

$$\frac{d\epsilon}{dr} = -\epsilon \left[ \frac{\nu_{\alpha\rho}}{U} + \frac{\nu_{\alpha\rho}}{U} (\epsilon + 1) \right] = -\frac{\nu_{\alpha\rho}}{U} \left[ \epsilon (1 + F) + \epsilon^2 F \right]$$

(7)

where we have employed standard expressions for the Coulomb collision frequency between two Maxwellian distributions,

$$\nu_{ss'} = 4\pi q_s^2 q_{s'}^2 \ln \Lambda n_{s'} m_s \mu_{ss'}$$

(8)

presented in Hernandez et al. (1987) with reduced mass ratio $\mu \equiv m_s m_{s'}/(m_s + m_{s'})$ and the combined thermal speed $w_{ss'}^2 = 2T_s/m_s + 2T_{s'}/m_{s'}$. To write the ratio of collision frequencies in terms of the mass density ratio

$$\frac{\nu_{\alpha\rho}}{\nu_{\alpha\rho}} = \frac{n_\alpha m_\alpha}{n_p m_p} \equiv F.$$

(9)

As Tracy et al. (2015) recently demonstrated, for all heavy ions in the solar wind including He$^{2+}$ the dominant coupling via Coulomb collisions is with H$^+$. As $\nu_{\alpha\rho}$ depends on both $T_\alpha$ and $T_p$, separating $\epsilon$ and $\nu_{\alpha\rho}$ as necessary for a solution to Eqn. 7 necessitates the construction of a ‘reduced’ collision frequency which only depends on a single temperature

$$\tilde{\nu}_{ss'} = 8\pi q_s^2 q_{s'}^2 \ln \Lambda n_{s'} m_s^3 w_{ss'}^2$$

$$= \frac{2\nu_{ss'}}{1 + m_s/m_{s'}} \left(1 + \frac{T_s m_{s'}}{T_{s'} m_s}\right)^{3/2}$$

(10)

where the single species thermal speed is $w_s^2 = 2T_s/m_s$. Using this reduced collision frequency, we separate Eqn. 7 into terms which do and do not depend on $\epsilon$, resulting in a differential equation of the form

$$\int_{R_b}^{R_w} \frac{2}{5} \left[ 1 + \frac{(\epsilon + 1)}{4} \right] \frac{3/2}{(1 + F) + \epsilon^2 F} \frac{d\epsilon}{\nu_{\alpha\rho}(r)} = -\frac{A_c}{U(r)}$$

(11)

where our solution depends on integration from the outer boundary of the zone of preferential heating $R_b$ to the radius of the observer $R_w$.

Expanding the numerator and performing typical $u$-substitutions, known integral identities, and arithmetic manipulations yields a closed form expression for the left-hand side of Eqn. 11:

$$\int_{R_b}^{R_w} \frac{2}{5} \left[ 1 + \frac{(\epsilon + 1)}{4} \right] \frac{3/2}{(1 + F) + \epsilon^2 F} \frac{d\epsilon}{\nu_{\alpha\rho}(r)} = \frac{\sqrt{5} + \epsilon w}{10 F} \left[ \frac{1}{2} \ln \left( \sqrt{1 + \frac{\epsilon w}{w}} + 1 \right) - \frac{1}{2} \ln \left( \sqrt{1 + \frac{\epsilon w}{w}} - 1 \right) \right]$$

$$+ \frac{(4F - 1)^{3/2}}{10F^{3/2}} \left[ \text{arctanh} \left( \sqrt{\frac{\sqrt{5} + \epsilon w}{\sqrt{4F - 1}} \right) - \text{arctanh} \left( \sqrt{\frac{\sqrt{5} + \epsilon w}{\sqrt{4F - 1}} \right) \right]$$

(12)

where $\epsilon_0$ and $\epsilon_w$ are the excess temperature ratio at the outer boundary of the zone of preferential heating and at 1 AU respectively.

A solution for $A_c$, the right-hand side of Eqn. 11, re-
quires a description for the radial evolution of the reduced collision frequency $\tilde{\nu}_{\alpha p}$ which depends on the radial structure of $n_p, U,$ and $T_p$, as well as a value for the boundary $R_b$. From Eqn. 10, $\tilde{\nu}_{\alpha p}$ varies as $n_p T_p^{-3/2}$. Both $n_p$ and $T_p$ fall off with distance from the Sun, so it is expected that $\tilde{\nu}_{\alpha p}$ should increase substantially closer to the Sun. Using the radial variations found in Helios observations (Hellinger et al. 2013), we take $T_p \propto r^{-\delta}$, $U \propto r^{-\sigma}$ and $n_p \propto r^{-2} U(r)$. From these scalings, $\tilde{\nu}_{\alpha p}$ as a function of the measured collision rate at 1 AU $\tilde{\nu}_{\alpha p}(R_w)$ may be written as

$$\tilde{\nu}_{\alpha p} = \tilde{\nu}_{\alpha p}(R_w) \frac{2+\sigma-1.5\delta}{2+2\sigma-1.5\delta}$$

and thus the collisional age integral is expressed as

$$A_c = \int_{R_w}^{R_b} dR \frac{\tilde{\nu}_{\alpha p}(R_w)}{U_w} \left[ \frac{1}{1+2\sigma-1.5\delta U_w} \frac{R_b}{R_w} \left( \frac{R_b}{R_w} \right)^{-1-2\sigma+1.5\delta} \right].$$

Note that care must be taken in evaluating this equation, as a singularity appears for $-2\sigma + 1.5\delta = 1$. We note the assumed scaling relations used for $\tilde{\nu}_{\alpha p}$ may not be accurate especially close to the Sun as temperatures in the corona are lower than extrapolations from the Helios trend lines close to the Sun. We will offer a post hoc justification that $R_b$ is sufficiently far from the Sun where these scaling relations serve as accurate descriptions.

Combining Eqns. 12 and 14 into Eqn. 11, one can produce a transcendental expression that relates $\epsilon_w$ to measured quantities $F, n_p, U, T_p$, fixed parameters $\delta, \sigma$, and free parameters $R_b, \epsilon_0$. Note that rather than determining the parameters that result in $\epsilon_w$ approach zero for high $A_c$, we allow our solution to relax to a residual $\epsilon_1$, which is treated as a free parameter in our modeling, to account for the fact that $\epsilon \simeq 0.2 - 0.3$ has been reported even in the case of high $A_c$ for both Wind/SWE observations of helium and hydrogen temperatures (Maruca & Kasper 2013; Kasper et al. 2013) and for heavier ions (Tracy et al. 2016). It is an open question whether this residual $\epsilon_1$ is indicative of continuing preferential heating acting in interplanetary space at a much reduced weaker level compared to in the zone of preferential heating, or if it is indicative of an instrumental measurement error in temperature ratios. The asymptotic value of the temperature excess below $R_b$ in the zone will therefore be $\epsilon_0 + \epsilon_1$. This modeled value $\epsilon_w$ will be compared to observed values of $\epsilon$ in the fashion described in the following section in an effort to indirectly measure the extent of the zone of preferential heating.

4. DETERMINATION OF ZONAL BOUNDARY USING THE WIND DATA SET

We now describe our procedure for solving for the outer boundary $R_b$ of the zone of preferential heating by comparing our model predictions for the excess temperature ratio with observations of the solar wind by the Wind spacecraft. The model is a function of solar wind speed, density, temperature, mass density ratio $F$, and spacecraft location for each measurement, along with five global free parameters: the boundary height $R_b$, the steady state excess temperature ratio within the zone $\epsilon_0 + \epsilon_1$, the residual excess at 1 AU $\epsilon_1$, and the radial exponents of solar wind speed $\sigma$ and temperature $\delta$. We use observations of the radial dependence of solar wind properties from the Helios mission to guide our choice of exponents, since as we will show, our values for $R_b$ are closer to Helios perihelion than they are coronal heights where there are spectroscopic measurements. Since there are different reported values for the radial temperature exponent $\delta$ in the literature (Marsch et al. 1982; Hellinger et al. 2013), we consider four values of $\delta, 0.7, 0.8, 0.9, 1.0$ that cover the reported range. Those same studies have also shown that $\delta$ may be a weak function of speed, so the observations were analyzed in separate 25 km s$^{-1}$ intervals in solar wind speed. While Fig. 2 clearly shows the same relaxation of $\epsilon$ with increasing Coulomb collisions continues at least to 650 km s$^{-1}$, we have limited our analysis to the range 300 – 425 km s$^{-1}$ in order to have good observational statistics at high and low $A_c$ with 25 km s$^{-1}$ interval size. For this work, we keep the solar wind speed exponent fixed at $\sigma = 0$. Observational studies have found negligible dependencies of the solar wind speed on radial distance, with a preference toward a shallow increase in $U$ with $r$ (Hellinger et al. 2011, 2013). As the speed and temperature exponents only appear as a linear combination in the expression for $A_c$, exploring the dependence of $R_b$ on $\delta$ also provides direct insight into the effects of changing $\sigma$.

For each range in solar wind speed and assumed value of $\delta$ and $\sigma$ we now determine the best fit values of our three free parameters $\epsilon_0, \epsilon_1$, and $R_b$ by minimizing the $\chi^2$ per degree of freedom difference between the model and the observations, weighted by an error estimate. It might seem like the easiest way to conduct this analysis would be to directly compare the predicted and observed $\epsilon$ for every individual observation in a given speed interval. We found that this was unreliable, as Fig. 2 shows there is a very strong preference for a particular collisional age at a given speed. In order to avoid biasing the analysis due to the most common age, we instead histogram our observations into bins in $A_c$, calculate the mean $\epsilon$ in each bin, and compare those means to the model value.

We start with an initial guess for $R_b, \epsilon_0$, and $\epsilon_1$. We found that the following analysis is highly insensitive to those initial guesses, but fitting a simple exponential curve to the data provided a good initial guess that speeds up the calculations. For each measurement of $\epsilon_w$ we then calculate an initial collisional age $A_c^i$ using the current values for $R_b, \epsilon_0$, and $\epsilon_1$. We also calculate an
overall average mass density ratio $F$ using the selected data. The measured $\epsilon$ are then binned as a function of the calculated $A_c$, with the range of $A_c$ and the resolution of our bins set up beforehand so there are always at least 1,000 individual measurements per bin. We use the average over all the selected data. We then calculate a prediction for $\epsilon_w$ for each bin using the transcendental expression resulting from Eqns. 11, 12 and 14. We do not want to use the uncertainty in the mean for each interval in $A_c$ for this analysis because the high correlation between speed and $A_c$ shown in Fig. 2 would strongly bias the best fit to the handful of intervals with the bulk of the observations. Instead we identified a constant error estimate based on the observed spread of $\epsilon$ at very high $A_c$. We found that at high $A_c$ the majority of epsilon observations are normally distributed with a width of 7%, and used this value as the error estimate at all $A_c$. The non-linear best fit is thus calculated by variation of $R_b$, $\epsilon_0$, $\epsilon_1$ and $\epsilon_b$ and iteration of the above binning routine, producing an estimate for the global minimum of $\chi^2/dof$ along with the best fit values and one-sigma uncertainties for the three free parameters for each interval of $U$ and value of $\delta$.

Fig. 4 illustrates the process and results for three solar wind speed intervals $U = 300 - 325, 350 - 375, 400 - 425$ km s$^{-1}$, using $\delta = 0.7, 0.6, 0.8$ respectively. Our simple model of a zone of preferential heating is able to predict the mean excess helium temperature to within 8-9% with a $\chi^2/dof = 1.6 - 1.9$. Typical values for $R_b$ are tens of solar radii from the Sun.

5. DISCUSSION

Following the procedure outlined in the previous section, the best fit values for $R_b$ as a function of solar wind speed and temperature power law exponent are calculated and shown in Fig. 5. For the value of $\delta = 0.75$, matching the observations of slow wind reported in Hellinger et al. (2013), the outer edge of the boundary ranges from 15 to 40 $R_S$ from the Sun’s surface for varying solar wind speed. The zone boundary is a decreasing function of $\delta$, and as shown in Fig. 6 we see that there is a simple linear relationship that allows one to correct $R_b$ for different assumptions of $\delta$, with the dependence of the boundary value on $\delta$ fairly independent of speed. Averaging over all the trends shown in Fig. 6, $R_b$ drops by 8.8 $R_S$ for every 0.1 increase in $\delta$. Physically, the faster temperature falls off with distance from the Sun, the smaller the preferential heating zone. Similarly, a linear relation between $U$ and $R_b$ can be found approximately satisfying $R_b \propto 0.1U$, not shown.

Perhaps the most significant implication of the inferred zone is that the preferential heating does not persist throughout the heliosphere, and thus can not be measured locally by spacecraft at 1 AU. Attempts to locally differentiate between proposed mechanisms for the preferential heating will rely on missions such as Parker Solar Probe (Fox et al. 2015) and Solar Orbiter (Müller et al. 2013) that will make measurements of particles and fields in the near-Sun region of the heliosphere (Kasper et al. 2015; Bale et al. 2016). Several key radial distances for the Parker Solar Probe mission are shown in Fig. 5. All but one of the predicted $R_b$ are below the starting distance for Parker Solar Probe science observations. By the end of the mission Parker Solar Probe will cross all but
one of the predicted values of $R_b$, allowing direct measurement of the region where the preferential heating is predicted to occur.

We note that our assumption that $R_b$ is sufficiently far away to employ radial scalings of $T_p$ and $U$ measured by Helios has been justified post hoc. Had $R_b$ been on the order of a few solar radii, a more sophisticated model for the radial dependence of the collision frequency would have become necessary. Additional modifications to this model, such as the inclusion of the effects of temperature anisotropy or other non-thermal features (Hellinger 2016) on the collision frequency may have a quantitative effect on the position of $R_b$, but we expect these effects to be small. Additionally, we will be able to improve our model by using measurements of the radial dependence of $n$, $U$, and $T$ by Parker Solar Probe to determine the accuracy of power-law extrapolations from Helios observations into the near-Sun environment.

When evaluating this model, one must address the nature of the energy input beyond $R_b$. While the structure of this model and the best-fit values of $R_b$ indicate that the preferential heating is limited to a region close to the Sun, this does not necessarily imply that no heating persists beyond this region. We know that ion heating of some level extends out to 1 AU and beyond (Cramer et al. 2009), but the rate most likely drops with distance. In the radial model of Chandran et al. (2011) for example, the heating rate is high until about $20 R_s$ and then it falls off as a power law. Our model allows for such heating as long as the heat input per particle for the $\alpha$ particles satisfies Eqn. 5. The persistence of a residual, non-zero $\epsilon$ even for high-$A_\alpha$ plasma, commented on in Maruca et al. (2013), may be an indication of a small amount of preferential heating of ions beyond $R_b$, or may be an instrumental limitation. Characterization of this residual excess temperature ratio will be left to future work.

We think the most plausible interpretation of our results is that $R_b$ is linked to the Alfvén critical point $R_A$, the radial distance where the solar wind transitions from being sub-Alfvénic to super-Alfvénic, and that the zone of preferential heating is simply the volume of space below the Alfvén point where reflected waves traveling back towards the Sun interact with escaping waves to enhance the turbulent cascade and allow it to transport significant energy in the form of intense and counter-propagating fluctuations down to ion kinetic scales. Typical predicted values for $R_A$ lie between 10 and 30 $R_s$ (Verdini et al. 2012; Perez & Chandran 2013), consistent with our findings for $R_b$. Since all sunward directed Alfvénic fluctuations generated below $R_A$ are trapped below $R_A$, it is natural to expect that intense reflection driven turbulence will be stronger in this region Verdini & Velli (2007).

As a specific example of a preferential ion heating mechanism that would be active below $R_A$, the presence of counter-propagating ion kinetic scale fluctuations consider Kasper et al. (2013), which showed that when $A_\alpha$ is small the dependence of $T_\alpha/T_p$ on plasma $\beta$ and normalized differential flow speed $\Delta V_{\alpha p}/C_A$ is consistent with heating by counter-propagating Alfvén ion-cyclotron waves (AIC, kinetic scale Alfvénic fluctuations propagating in opposite directions along the local magnetic field) which are significantly more efficient at heating $He^{2+}$ relative to $H^+$ and in fact all other ions heavier
than $H^+$. In the presence of a spectrum of AIC waves propagating in a single direction, ions heavier than $H^+$ are slowly heated as resonant scattering diffuses them in velocity space about the phase speed of the waves. If counter propagating waves are introduced, $He^{2+}$ and heavier ions can scatter off waves traveling in opposite directions, permitting a more general diffusion in velocity space and a far more rapid and preferential heating. A unified explanation of the observations could be as follows. Everywhere below the Alfvén point, counter-propagating Alfvén waves are present and strongly preferentially heat ions heavier than $H^+$. This heating is first apparent $0.1 - 0.3 R_S$ above the surface of the Sun when the Coulomb collisions are no longer able to prevent temperature differences from emerging. From that height up to the Alfvén point all ions are heated by these counter-propagating waves and diffuse in phase space to reach an equilibrium temperature excess. Suddenly at the Alfvén point reflected waves are not able to travel back towards the Sun, and the power in counter-propagating waves drops significantly, shutting off the counter-propagating AIC mechanism. The temperature excess developed below the Alfvén point then decays through Coulomb relaxation to the level observed at an interplanetary spacecraft. This sharp drop in heating at $R_A$ is consistent with our assumption that heating stops at $R_b$ and could help explain why our simple model for heating with distance fits the observations so well. In this framework the reason Kasper et al. (2013) could only see their correlations with AIC predictions in low $A_e$ solar wind is because they were never directly observing the heating in action, but instead a signature frozen into solar wind ions that crossed $R_A$, and therefore $R_b$, days earlier. Finally, the small residual temperature excess seen for ions in high $A_e$ plasma (Kasper et al. 2013; Maruca et al. 2012; Tracy et al. 2016) could then be due to the weaker heating of ions by AIC fluctuations traveling predominantly in one direction.

Another intriguing possibility is that $R_b$ could correspond to the distance recently identified in DeForest et al. (2016) where a transition from relatively steady and laminar radial flow to sheared and turbulently mixed flow is remotely observed. Using an improved analysis of remote observations from the Heliospheric Imager on STEREO, the authors identified a region $\sim 40 - 80 R_S$ from the solar surface in which the smooth radial expansion of the slow solar wind appears to fragment and break up. As with the Alfvén critical point, the transition in solar wind flow structure reported by DeForest et al. (2016) could signify another boundary in the solar wind that separates different levels of fluctuations and dominant heating mechanisms. Of course it is also possible that the breakup in smooth flow seen in the images is simply a manifestation of the Alfvén point itself, or another height where the Alfvén mach number crosses some value and the plasma becomes unstable.

Lastly, we can use the values for $\epsilon_0 + \epsilon_1$ to look at the implied excess temperature ratio back in the zone of preferential ion heating for direct comparison with coronal heating theories and with other observations. Fig. 7 shows the coronal excess temperature of helium relative to hydrogen as a function of solar wind speed and temperature exponent $\delta$, with the symbols for $R_b$ at different $\delta$ using the same color scheme as Fig. 6. Here we see another significant signature of the physical process responsible for preferential ion heating in the inner heliosphere. Our analysis shows an average temperature excess of about $4.2 - 4.3$, meaning that that helium is $5.2 - 5.3$ times hotter than hydrogen, independent of solar wind speed and our assumption for $\delta$. This excess would appear to be a highly significant and model-independent result. Any theory of heating in the corona or extended solar wind should be able to produce this steady state excess temperature. Finally, we note the dashed red lines on the figure, which were developed by taking the temperature dependency in low $A_e$ or collisionless solar wind reported by Tracy et al. (2016) for heavy ions in the solar wind, but evaluated for the mass of helium. We find that within the error reported in that analysis, the implied steady state helium temperature excess in the corona is consistent with the mass dependence of heavy ions in the solar wind when selecting collisionless wind that presumably indicates the coronal values. We therefore propose that our results, combined with those of Tracy et al. (2016), are consistent with the idea that there is a preferential ion heating mechanism acting in a zone of non-thermal heating, that acts on all ions, and extends out tens of solar radii from the Sun. Within
this zone ion temperatures reach a steady state ratio of
\( T_i/T_p = (4/3)m_i/m_p \), independent of solar wind speed.

6. CONCLUSION

We have examined the temperature ratio of fully ionized He\(^{2+}\) and H\(^+\) in the solar wind and its dependence on Coulomb collisional age in order to solve for the location of an outer boundary of an apparent zone of preferential ion heating in the inner heliosphere. Using millions of observations from the Wind spacecraft in concert with a physically motivated model for the excess temperature ratio, we are able to construct a best fit value for the outer boundary of this region, which falls between \(20 \sim 40 R_S\) with some variation with solar wind speed and radial temperature scalings. The restricted radial extent of this region would frustrate attempts to identify preferential heating mechanisms using measurements at 1 AU, but future missions including Parker Solar Probe will provide measurements both within and outside this region, allowing for the novel measurement of the mechanisms that lead to the non-thermal heating of solar wind minor ions.

We can now answer the three questions proposed in the Introduction. The large unequal ion temperatures seen in situ by spacecraft in the solar wind are not maintained by ongoing local and strong preferential heating. Instead they are a leftover of heating that happened closer to the Sun. Solar wind at all speeds appear to experience strong preferential heating within our proposed zone, and the only reason fast wind appears more non-thermal than slow wind in interplanetary space is due to the large difference in Coulomb collisions that transpire as the solar wind travels from the outer boundary of the zone to the observing spacecraft. The strong preferential ion heating seen close to the Sun in spectroscopic observations continues \(20 \sim 40 R_S\) from the Sun before dropping off, perhaps due to a lack of counter-propagating Alfvénic fluctuations. It is possible that the residual temperature excess observed in interplanetary space indicates that a weaker form of preferential heating is active outside of the zone, but it is only able to produce temperatures that are different by tens of percent.

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REFERENCES

Axford, W. I., & McKenzie, J. F. 1997, in Cosmic Winds and the Heliosphere, ed. J. R. Jokipii, C. P. Sonett, & M. S. Giampapa, 31
Bale, S. D., Goetz, K., Harvey, P. R., et al. 2016, Space Sci. Rev., doi:10.1007/s11214-016-0244-5
Cargill, P. J., & Klimchuk, J. A. 2004, ApJ, 605, 911
Chandran, B. D. G. 2010, ApJ, 720, 548
Chandran, B. D. G., Dennis, T. J., Quataert, E., & Bale, S. D. 2011, ApJ, 743, 197
Chandran, B. D. G., Li, B., Rogers, B. N., Quataert, E., & Gemmrich, K. 2010, ApJ, 720, 503
Chandran, B. D. G., Verscharen, D., Quataert, E., et al. 2013, ApJ, 776, 45
Cramer, S. R. 2000, ApJ, 532, 1197
Drake, J. F., Cassak, P. A., Shay, M. A., Swisdak, M., & Quataert, E. 2009, ApJ, 700, L16
Eser, R., Fineschi, S., Doebrycka, D., et al. 1999, ApJ, 510, L63
Feldman, W. C., Asbridge, J. R., & Bame, S. J. 1974, J. Geophys. Res., 79, 2319
Fox, N. J., Velli, M. C., Bale, S. D., et al. 2015, Space Sci. Rev., doi:10.1007/s11214-015-0211-6
Helliger, P. 2016, ApJ, 825, 120
Helliger, P., Matteini, L., Štverák, Š., Trávníček, P. M., & Marsch, E. 2011, Journal of Geophysical Research (Space Physics), 116, A09105
Helliger, P., Trávníček, P. M., Štverák, Š., Matteini, L., & Velli, M. 2013, Journal of Geophysical Research (Space Physics), 118, 13731
Hernandez, R., Livi, S., & Marsch, E. 1987, J. Geophys. Res., 92, 7723
Hollweg, J. V., & Isenberg, P. A. 2002, J. Geophys. Res., 107, 1147
Kasper, J. C., Lazarus, A. J., & Gary, S. P. 2002, Geophys. Res. Lett., 29, 20
Kasper, J. C., Lazarus, A. J., & Gary, S. P. 2008, Physical Review Letters, 101, 261103
Kasper, J. C., Maruca, B. A., Stevens, M. L., & Zaslavsky, A. 2013, Physical Review Letters, 110, 091102
Kasper, J. C., et al. 2006, J. Geophys. Res., 111, 3105
Klein, K. G., & Howes, G. G. 2016, ApJ, 826, L30
Lei, D., & Cranmer, S. R. 2009, ApJ, 691, 794
Lepping, R. P., Acna, M. H., Burlaga, L. F., et al. 1995, Space Sci. Rev., 71, 207, 10.1007/BF00751330
Livi, S., Marsch, E., & Rosenbauer, H. 1986, J. Geophys. Res., 91, 8045
Marsch, E. 2012, Space Sci. Rev., 172, 23
Marsch, E., et al. 1982, J. Geophys. Res., 87, 35
Maruca, B. A., Bale, S. D., Sorriso-Valvo, L., Kasper, J. C., & Stevens, M. L. 2013, Physical Review Letters, 111, 241101
Maruca, B. A., & Kasper, J. C. 2013, Advances in Space Research, 52, 723
Maruca, B. A., Kasper, J. C., & Gary, S. P. 2012, ApJ, 740, 137
Marsi, W., H., Zank, G. P., Oughton, S., Mullan, D. J., & Dimitruk, P. 1999, ApJ, L93
Müller, D., Marsden, R. G., St. C., C., & Gilbert, H. R. 2013, Sol. Phys., 285, 25
Neugebauer, M. 1976, J. Geophys. Res., 81, 78
Ogilvie, K. W., Chornay, D. J., Fritzemeier, R. J., et al. 1995, Space Sci. Rev., 71, 55
Perez, J. C., & Chandran, B. D. G. 2013, ApJ, 776, 124
Salem, C., Hubert, D., Lacombe, C., et al. 2003, ApJ, 585, 1147
Scudder, J. D. 1992, ApJ, 398, 299
Spitzer, L. 1962, Physics of Fully Ionized Gases (Interscience)
Tracy, P. J., Kasper, J. C., Raines, J. M., et al. 2016, Physical Review Letters, 116, 255101
Tracy, P. J., Kasper, J. C., Zurbuchen, T. H., et al. 2015, The Astrophysical Journal, 812, 170
Verdini, A., Grappin, R., Pinto, R., & Velli, M. 2012, ApJ, 750, 133
Verdini, A., & Velli, M. 2007, ApJ, 662, 669

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