

Stochastic Heating in the Sub-Alfvénic Solar Wind

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Collisionless dissipation of turbulence is important for heating plasmas in astrophysical, space physics, and laboratory environments, controlling energy, momentum, and particle transport. We analyze Parker Solar Probe observations to understand the collisionless heating of the sub-Alfvénic solar wind, which is connected to the solar corona. Our results show that linear resonant heating through parallel-propagating cyclotron waves cannot account for turbulent dissipation in the sub-Alfvénic region, which observations suggest may dissipate turbulence at distances further from the Sun. Instead, we find that stochastic heating can account for the observed ion energization; however, because the dominant contributions arise from infrequent, large-amplitude events, turbulent intermittency must be explicitly incorporated. These observations directly connect stochastic heating via breaking of the proton magnetic moment with the intermittent and inhomogeneous heating of turbulence reported in many previous studies. Our identification of stochastic heating as a dynamic mechanism responsible for intermittent heating of the solar wind has significant implications for turbulent dissipation in the lower corona, other astrophysical environments, and laboratory plasmas.

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The collisionless heating of the solar wind and corona is a well-known phenomenon [1,2] with important implications for dissipation and heating of other astrophysical environments [3–7] as well as laboratory and fusion plasmas [8–10]. Remote observations of the inner corona reveal significant temperature anisotropy with $T_{\perp} > T_{\parallel}$ [11–14], a common feature of collisionless plasmas. Residual anisotropic heating is additionally present in the solar wind [15,16], resulting in nonadiabatic temperature profiles with distance [17–20]. These observations have led to significant interest in mechanisms that preferentially heat plasmas perpendicular to the magnetic field direction, which can be broadly separated into either quasilinear resonant electromagnetic field-particle interactions [1,21–23] or nonresonant mechanisms [24–29].

Recent observations from Parker Solar Probe (PSP) [30–33] have highlighted the importance of quasilinear (QL) heating through left-hand polarized electromagnetic

fields at ion-kinetic scales, which are likely quasiparallel ion-cyclotron waves ($\parallel$ -ICWs) [34–39]. Similar observations provide evidence for cyclotron heating in the 1 A.U. solar wind [40–43] and turbulent magnetosheath [44], suggesting that quasilinear resonant ICW interactions play an important role in collisionless dissipation in a range of environments. While there is a growing consensus regarding a role of $\parallel$ -ICWs in mediating turbulent heating [45], there is a long-standing debate regarding whether the observed $\parallel$ -ICWs are directly responsible for the dissipation of turbulence or if they are by-product of other mechanisms [46,47]. This important debate largely stems from the tendency of Alfvénic turbulence to transfer energy efficiently to high perpendicular wave number $k_{\perp}$ rather than to high parallel wave number $k_{\parallel}$ [48–54], leading to anisotropic fluctuations at small scales with $k_{\perp} \gg k_{\parallel}$; this type of cascade cannot directly produce $\parallel$ -ICWs. It is often thought that $\parallel$ -ICWs are generated via the $T_{\perp} > T_{\parallel}$ Alfvén or ICW instability [55], which cools the plasma [56]. Observations [31,40,57,58] show that the ion velocity distribution function (VDF) is roughly constant along the $\parallel$ -ICW diffusion contours [59], implying marginal stability and further complicating observation and interpretation of the dynamics.

If resonant interactions cannot supply the requisite levels of heating for turbulent dissipation, other, inherently nonlinear, processes must take an active role. Such mechanisms,

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often termed “stochastic,” occur in a variety of contexts: cosmic ray transport [60], diffusive shock acceleration [61], turbulent dissipation in the magnetopause [24], drift-wave turbulence [62,63], and energetic ions in tokamak experiments [64]. In the solar wind and corona, stochastic heating (SH) of ions, understood as a random walk in energy spurred by repeated nonresonant interactions with turbulent fluctuations [26,65,66], has been suggested as a viable plasma heating mechanism [67–72].

In contrast to the canonical formalism of QL theory [21], stochastic heating is highly sensitive to intermittency in the turbulence [73,74], which is understood as the broad, heavy-tailed, distribution of fluctuation amplitudes that intrinsically formed by turbulence [75–77]. Importantly, correlations between intermittency and heating are widely recognized [78–83], which suggests that QL resonant mechanisms that are canonically derived under nonintermittent conditions (e.g., [21]) cannot entirely explain turbulent heating. Constraining the interplay of QL mechanisms, which are important in understanding the observed wave populations [31,84], with intermittent nonresonant mechanisms, such as SH, is imperative in understanding dissipation [73].

Here, we use PSP [85] observations to investigate heating of the sub-Alfvénic solar wind, where the Alfvén speed exceeds the solar wind speed [86]. In contrast to previous results from larger heliocentric distances [31,32], our results show the plasma is locally unstable to $\parallel$ -ICWs [55,87] but with a net energy transfer that is far below the local turbulent cascade rate. Thus, $\parallel$ -ICW interactions are not energetically relevant in bulk heating of this sub-Alfvénic solar wind stream. While the $\parallel$ -ICWs cannot explain turbulent dissipation, we find SH rates [26,69] that are consistent with the turbulent cascade rate [88,89], if (and only if) we account for the intermittent nature of the turbulence. These observations suggest a connection between SH via magnetic moment breaking [26] and many studies that highlight intermittent and inhomogeneous turbulent heating [78–83]. Conversely, the fact that intermittency is required to attain SH rates that match the turbulent cascade implies that nonintermittent fluctuations may be irrelevant in heating. These results provide strong constraints on collisionless dissipation in Alfvénic turbulence with implications for coronal heating [46,90,91].

SH model—Studies of SH in Alfvénic turbulence show strong sensitivity to the fluctuation amplitude [74,92–94]. The efficiency of SH is quantified via a stochasticity parameter $\epsilon_u = \delta u / v_\perp$ [26], where δu is the turbulent fluctuation amplitude at a particle’s gyroradius ρ_i and the perpendicular speed $v_\perp = \rho_i \Omega_i$ with $\Omega_i = qB/m_i$. Chandran *et al.* (hereafter, *C + 2010* [26]) estimate the SH energy diffusion coefficient of low- β Alfvénic turbulence as

$$D_E = c_1 v_\perp \delta u^3 \Omega_i e^{-c_2/\epsilon_u} \quad (1)$$

with constants c_1 and c_2 of the order of unity and a corresponding velocity-space diffusion coefficient

$D_\perp = D_E / v_\perp^2$. The exponential in Eq. (1) results in extreme sensitivity to ϵ_u such that SH is irrelevant for $\delta u \ll v_\perp$ [74]. While D_E is a function of $v_\perp$, *C + 2010* specifically focus on SH at the thermal gyroradius $\rho_{th} = v_{th\perp} / \Omega_i$ noted as $Q_{SH}^{\rho_{th}}$, with

$$Q_{SH}^{\rho_{th}} = c_1 \frac{\delta u_{\rho_{th}}^3}{\rho_{th}} e^{-c_2/\epsilon_{\rho_{th}}}, \quad (2)$$

where $v_{th\perp}$ is the perpendicular ion-thermal speed, $\delta u_{\rho_{th}}$ is the fluctuation amplitude at the thermal gyroscale, and $\epsilon_{\rho_{th}} = \delta u_{\rho_{th}} / v_{th\perp}$. Observational studies have used values of $c_1 \approx 0.75$ and $c_2 \approx 0.34$ [71,72], based on test particle simulations [26]. Alternatively, matching the SH rate Q_{SH} to the turbulent cascade rate ϵ^T [67] or via numerical simulations [68] can constrain the coefficients.

While the *C + 2010* expression for SH in Eq. (2) considers only heating localized at the thermal-perpendicular speed, more general SH formalisms [69,74] express SH as a $v_\perp$ -dependent diffusion of a gyrotropic VDF, with

$$Q_{SH} = \frac{m_i}{2} \int v_\perp^2 \frac{\partial f}{\partial t} d^3v = -2\pi m_i \int_0^\infty v_\perp^2 D_\perp \frac{\partial F(v_\perp)}{\partial v_\perp} dv_\perp, \quad (3)$$

where the reduced VDF is $F(v_\perp) = \int f(v_\perp, v_\parallel) dv_\parallel$ and the last term in Eq. (3) is obtained by integration by parts. We also define $Q_{SH} = \int dQ_{SH}$. An ion with perpendicular velocity $v_\perp$ has a gyroradius $\rho = v_\perp / \Omega$, and, thus, $F(v_\perp)$ can equivalently be considered a distribution of gyroradii corresponding to perpendicular speeds via $\rho_i = v_\perp / \Omega_i$. The quantity δu in Eq. (1) corresponds to the turbulent amplitude at the gyroradius of the particles with $v_\perp$ and, thus, can also be considered a function of $v_\perp$. While scale-dependent SH has been applied in both analytic theory and simulation [69,74], we presently apply the scale-dependent approach to *in situ* observations. Importantly, extending D_E strictly based on *C + 2010* is valid if $v_\perp > v_{th\perp}$, as these particles are heated by fluctuations at scales $\rho > \rho_{th}$, and the magnetohydrodynamics electric field is appropriate [69]. It is technically not valid to extend it to lower $\rho < \rho_{th}$ since the scaling of the electric field fluctuations changes for scales below the thermal ion gyroradius (e.g., the Hall field); however, as we will show, minimal heating occurs when $\rho < \rho_{th}$.

Methods—We calculate a scale-dependent fluctuation amplitude of the magnetic field δB using five-point increments [95,96] over a range of 88 time lags τ , ranging from $0.85 \text{ ms} < \tau < 7.23 \text{ s}$, corresponding to frequencies $0.07 < f < 586 \text{ Hz}$ with $f = 1/2\tau$ [97]. The modified Taylor hypothesis [98], which is appropriate for large-amplitude outward-propagating Alfvénic fluctuations, determines the corresponding spatial scales and is necessary in sub-Alfvénic wind where $V_A > V_{sw}$ [99,100]:

$$2\pi f = k_\parallel V_A + k_\parallel V_{sw} \cos \theta_{VB} + k_\perp V_{sw} \sin \theta_{VB}. \quad (4)$$

The Doppler shift equation can be used to connect perpendicular fluctuations with $k_{\perp}\rho = 1$ to f as $2\pi f = k_{\perp}V_{\text{sw}}\sin\theta_{VB}$, which enables conversion from δB to δu assuming Alfvénic turbulence and a transition to the kinetic-Alfvén wave (KAW) mode at ρ_{th} such that $\delta u = \delta B/B_0 v_{\text{ph}}^{\text{kaw}}$, where $v_{\text{ph}}^{\text{kaw}} = v_A\sqrt{1 + k_{\perp}^2\rho_{\text{th}}^2}$ [101]. This correction is valid for frequencies $\omega < \Omega_i$, corresponding to the largest sub-ion kinetic scales, and for highly imbalanced turbulence, such as this stream [102,103].

Relating fluctuations at a spacecraft frequency f to wave vector $\mathbf{k}$ in the sub-Alfvénic wind requires significant care. While effects from $k_{\parallel}$ are not needed for KAW conversion with usual dispersion relations [101,104], contributions from $k_{\parallel}$ are significant in converting observations at lag τ to spatial scales in the sub-Alfvénic wind [98]. Assuming a spectrum of critically balanced turbulence [49,53], $k_{\perp}\delta u = k_{\parallel}V_A$, uniquely relates fluctuations at lag τ to a specific $k_{\perp}$ (see End Matter); we note the KAW conversion from δB to δu is not affected by this model. Then, requiring $k_{\perp}\rho = 1$ associates each τ with a single $v_{\perp} = \Omega_i/k_{\perp}$ (computed for each interval). We can then construct $\epsilon_u(f, t)$ from the increment estimate of the amplitude and $v_{\perp}$ allowing for an estimate of D_E in Eq. (1). We note that the importance of the $k_{\perp}\rho = 1$ is an extension of the original C + 2010 formalism, whereas recent extensions of SH suggest that timescales (rather than spatial scales) may be most important [105,106].

Turbulence is inherently intermittent [82], which is important to the effectiveness of SH [73] through locally enhancing ϵ_u [74], which can counter the exponential suppression in Eq. (1). The importance of intermittency also arises in the third-order dependence of $D_E \propto \delta u^3$, which indicates an inherent sensitivity to large-amplitude, rare fluctuations. In calculating SH rates, we consider two different estimates of D_E . The usual estimate, $D_E^{\text{rms}}(f)$, is constructed from the rms $\delta u^{\text{rms}} = \sqrt{\langle \delta u^2 \rangle}$ amplitude over a time interval, which averages out intermittent effects. In the second, D_E is calculated individually for each increment $\delta u(t, f)$, giving an ensemble $D_E(t, f)$, which includes the effects of intermittency. The time average $\bar{D}_E = \langle |D_E| \rangle$ then includes intermittency in δu . This method amounts to treating Q_{SH} and D_E as functions of the random variable δu with average quantities integrated over the distribution $\langle X \rangle = \int X(\delta u)P(\delta u)d\delta u$.

Data and results—We study a stream from April 8, 2021 spanning 10:00:00–14:00:00, which was identified as the first prolonged encounter of PSP with sub-Alfvénic solar wind [86]. Several studies have investigated the turbulent properties of this interval [102,103,107], which has $\beta \approx 0.05$ and is highly imbalanced [103], suggesting applicability of the C + 2010 SH mechanism.

The interval is broken into nonoverlapping subintervals lasting 55.92 s ($\approx 1500\Omega_i$), which is sufficient to calculate SH rates well into the inertial range. Increment analysis is applied to each interval to construct $\bar{D}_E$ and D_E^{rms} . In each subinterval, the $\parallel$ -ICW spectrum is computed by filtering

out wavelet coefficients without strong left-hand polarization. The $\parallel$ -ICW heating rate $Q_{\parallel\text{ICW}}$ can be empirically calculated from the $\parallel$ -ICW spectrum and the gradients of the VDF [21,31,32], which we obtain via the PSP SPAN-i instrument [108,109] at a 3.5 s cadence. We apply drifting-bi-Maxwellian fits [31,110,111] to each VDF in each subinterval. The heating rates Q_{SH} , $Q_{\text{SH}}^{\text{RMS}}$, and $Q_{\parallel\text{ICW}}$ are then computed (see End Matter for further discussion).

Figure 1(a) shows the measured Q_{SH} and $Q_{\text{SH}}^{\text{RMS}}$ for the drifting bi-Maxwellian fits to the VDF from $\bar{D}_E$ and D_E^{rms} . Dashed lines show median estimates of heating; the median Q_{SH} rate is similar to the median turbulent cascade rate ϵ^T computed via local energy transfer (LET [89]; see End Matter). The absolute value of $Q_{\parallel\text{ICW}}$ is plotted in red when $Q_{\parallel\text{ICW}} > 0$ (heating) and blue when $Q_{\parallel\text{ICW}} < 0$ (emission). Throughout, SH dominates over $\parallel$ -ICW heating, with $|Q_{\parallel\text{ICW}}| \ll Q_{\text{SH}}$. Overall, there is net $\parallel$ -ICW emission, suggesting that waves are generated by the A/ICW instability. Oscillating positive and negative $Q_{\parallel\text{ICW}}$ suggests the plasma is marginally stable with respect to the A/ICW instability.

Figure 1(b) shows probability distributions of the various heating rates. The SH rates computed from D_E^{rms} are more than an order of magnitude lower than those computed from $\bar{D}_E$, which includes intermittency. Our means of

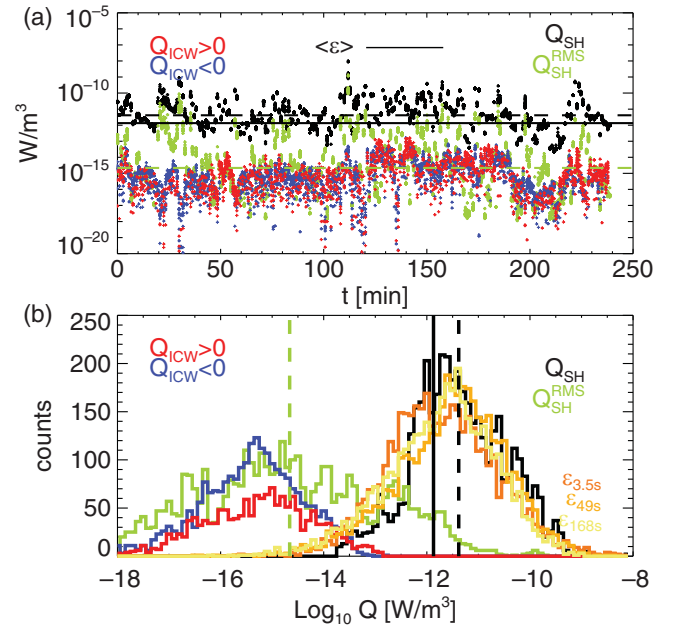


FIG. 1. (a) Heating rates as a function of time: stochastic heating Q_{SH} (black), nonintermittent stochastic heating $Q_{\text{SH}}^{\text{RMS}}$ (green), ICW emission $Q_{\parallel\text{ICW}} < 0$ (blue), and ICW heating $Q_{\parallel\text{ICW}} > 0$ (red). Median values for Q_{SH} and $Q_{\text{SH}}^{\text{RMS}}$ are shown as dashed lines. The solid black line shows the median local-energy transfer (LET) estimate $\langle \epsilon \rangle$ of the turbulent cascade rate. (b) Distributions of heating rates of data shown in the top panel; additionally, distributions of ϵ from LET estimate at lags corresponding to inertial range scales with lags at 3.5, 49, and 168 s are shown (orange, gold, and yellow, respectively).

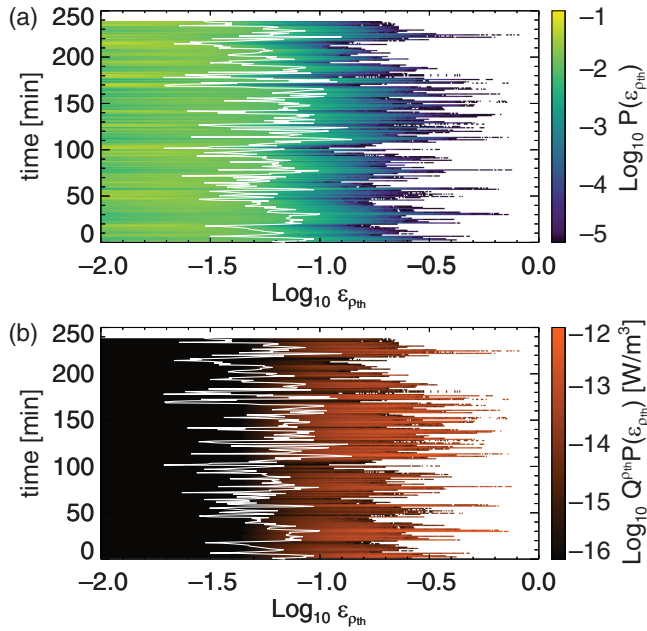


FIG. 2. (a) Spectrogram showing distribution of normalized amplitude $P(\epsilon_{\rho_{th}})$ evaluated at the thermal proton gyroscale $\epsilon_{\rho_{th}} = \delta u_{\rho_{th}}/v_{th\perp}$ in each interval. (b) Distribution of normalized amplitudes weighted by SH rate $P(\epsilon_{\rho_{th}})Q_{SH}^{\rho_{th}}$, showing that heating occurs in the tails. The white contour shows the largest 25% level of $P(\epsilon_{\rho_{th}})$.

including intermittency via averaging over D_E differs from that previously applied in numerical work of [74] but reveals the significant enhancement to the SH rate due to intermittent fluctuations [73].

To further highlight the role of intermittency in SH, Fig. 2(a) shows the probability distribution of the stochasticity parameter measured at the thermal gyroscale $P(\epsilon_{\rho_{th}})$ in each interval (the single scale is chosen for illustrative purposes). Figure 2(b) shows $P(\epsilon_{\rho_{th}})$ weighted by the heating rate at the thermal gyroradius [Eq. (2)], i.e., $P(\epsilon_{\rho_{th}})Q_{SH}^{\rho_{th}}$, which demonstrates the level of heating associated with a given $\epsilon_{\rho_{th}}$. The heating very clearly occurs in the intermittent, large-amplitude tails of the distribution [82], which have occurrence probabilities $\lesssim 0.1\%$. The upper quartile is shown in white.

Figure 3(a) shows the power spectral density from the SCAm data converted to Alfvénic units for a single subinterval; the KAW correction to the spectra is additionally illustrated. We also plot the scale-dependent dQ_{SH} . Heating is observed to peak in the transition range [112–114], which has been interpreted as a signature of dissipation [115–117]. Figure 3(b) shows an estimate of the moving spectral index of the turbulent spectra, α : Again, it is clear that SH becomes effective at frequencies associated with spectral steepening. Figure 3(c) shows a spectrogram of dQ_{SH} . A white contour is shown outlining the range of scales for which the spectral index is very steep with $\alpha < -4$. Over the entire interval, it is clear that the SH occurs at the range of scales corresponding to transition

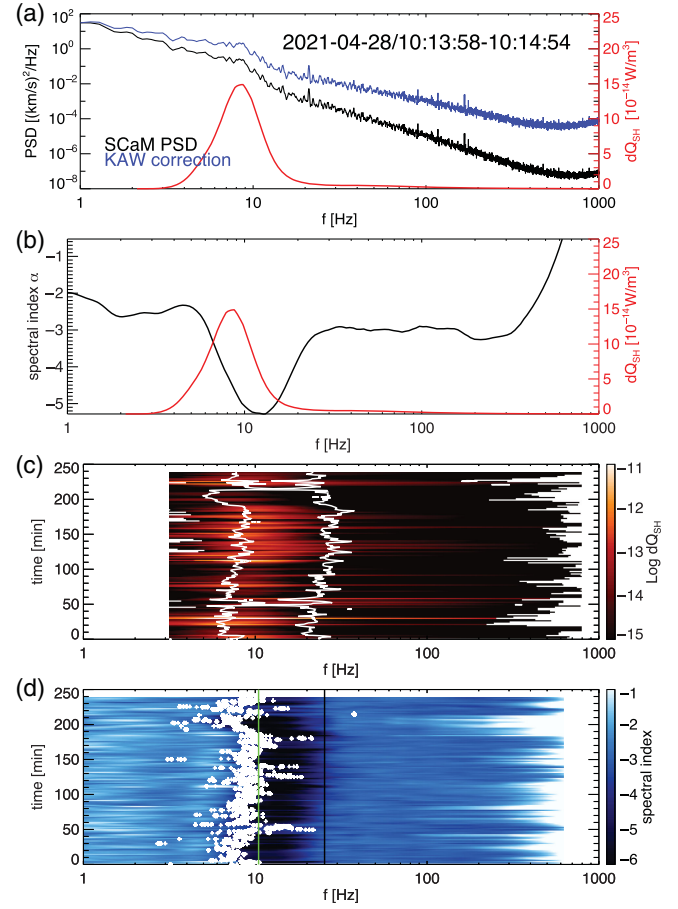


FIG. 3. (a) Power spectral density (PSD) constructed from PSP FIELDS magnetic field from April 2021, 28, 10:13:58–10:14:54 [118,119]; PSD is normalized using Alfvénic speed and the KAW normalization has been applied in blue; differential heating $dQ_{SH}(f)$ is shown in red. (b) Moving window spectral index α computed from (non-KAW-normalized) PSD. (c) Spectrogram showing dQ_{SH} as a function of frequency and time; a white contour shows where spectral index of fluctuations has $\alpha < -4$. (d) Moving window spectral index α as a function of frequency and time; white dots show the f with the maximum dQ_{SH} . The black line shows average $f = V_{sw}/2\pi\rho_i$; the green line shows average $f = V_{sw} \sin \theta_{VB}/2\pi\rho_i$.

range steepening. Figure 3(d) shows the estimate of the spectral index; white dots mark the location of the peak SH value, which occurs in the low end of the transition range. Figure 3(d) also shows the frequency corresponding to the average gyroscale $f = V_{sw}/2\pi\rho_i$ as well as the average $f = V_{sw} \sin \theta_{VB}/2\pi\rho_i$. The small amplitudes below the transition range decrease the effectiveness of SH via the exponential suppression in Eq. (1): Moreover, the fact that Q_{SH} is comparable to the overall cascade rate suggests that SH contributes to the steepening, with a considerable fraction of the turbulent cascade power heating the ions at these scales. These results indicate that SH may be more effective at scales larger than the gyroradius [105], contrasting studies that focus on SH at the thermal gyroscale [26,70,71,74] (see End Matter).

Discussion—Stochastic mechanisms are important to a variety of plasma processes and environments [24,25,61–64,94,120] including heating in the solar wind and corona. By extending the C + 2010 SH formalism [26] through including both scale dependence [69] and intermittency [73], we demonstrate that heating via SH is largely consistent with observed turbulent cascade rates, as measured through local energy transfer [88,89]. Our focus on Alfvénic turbulence is motivated by widely realized observations of strong Alfvénicity of turbulent fluctuations in the inner heliosphere [121–124]; however, stochastic scattering via other coherent structures, e.g., pressure balance structures, magnetically dominated intermittent current sheets [125], and balanced turbulence [126], may also contribute to heating. Our results indicate that stochastic magnetic moment breaking by intermittent structures in imbalanced turbulence is a viable heating mechanism in the sub-Alfvénic solar wind and provides important insight into coronal heating.

PSP studies of turbulent heating at larger radial distances found net $\parallel$ -ICW resonant heating on the order of the turbulent cascade rate [31–33]. Our present results indicate that, in contrast, in the sub-Alfvénic wind $\parallel$ -ICWs are close to marginal stability and, overall, cool the plasma. Further work is needed to understand the transition, distribution, and difference between intervals where ICW heating is important [31–33] versus those where it is negligible: The sub-Alfvénic nature of the present interval may indicate a radial dependence on heating mechanisms, with $\parallel$ -ICWs becoming more important in super-Alfvénic regions. Further studies should investigate this transition, which may be the result of radial variations in the turbulence or β , and whether it supports signatures of preferential heating near the Alfvén surface [127].

Importantly, we show empirically that the SH mechanism can provide the relevant heating levels only when intermittent effects are included; this point has been noted in theoretical work and simulations [73,74], though previous observational studies of SH [67,70–72] have not included intermittency. Our observations suggests that nonintermittent fluctuations may not be sufficiently large to drive SH. Our results support the body of work that invokes dissipation by intermittent coherent structures to explain observed inhomogeneous heating [79,80] that is associated with discontinuities [75,78,83,128–133]. It is well understood that intermittency significantly enhances cascade rates [134,135], such that turbulent dissipation must inherently be sensitive to intermittency. Our observed levels of SH, enabled by intermittency, provide an important connection of the dynamic process of magnetic moment breaking [25,26,65,94,136] to the large body of work that empirically connects intermittency with collisionless heating but where the specific dynamic mechanisms remain unknown. Indeed, our results in Fig. 2 suggest that SH may be suppressed for the vast majority of nonintermittent fluctuations.

The interplay of linear and nonlinear processes is a contentious issue [82,137–143] that is fundamental to

understanding the turbulent mechanisms that shape astrophysical [7,144] and laboratory [10] plasmas. Our results clearly support the role of intermittency in turbulent dissipation. Recently, [105] proposed that SH and QL heating by oblique ICWs are related processes (with a similar D_E) that exist, respectively, within limits of balanced and imbalanced turbulence; additionally, they predict that heating in imbalanced turbulence should peak above the transition range, consistent with our observations, and in contrast with previous focus on SH at the thermal gyroscale [67,70,71,73,74]. Further work is needed to explore similarities of cyclotron resonant and stochastic processes. Importantly, QL modeling [46,47,105,126] does not incorporate the effects of intermittency, which we show is likely necessary to reconcile observed turbulent cascade and heating rates. These insights are particularly interesting in the light of simulations suggesting that SH should primarily occur in balanced, low cross helicity turbulence, rather than the highly imbalanced, Alfvénic turbulence observed in this interval [126], which could be subject to a helicity barrier-type [47,145] mechanism that may result in oblique cyclotron resonant heating. Understanding intermittent effects on QL cyclotron interactions [73] would help clarify this interplay. Additionally, while our turbulent cascade rates agree with the SH levels, empirical measurements of D_E [146,147] and investigating α -particle signatures of these mechanisms [126,148–150] may help distinguish between (or, indeed, unify) QL processes and SH.

These results highlight several subsequent avenues for continued investigation. ICW dissipation has been shown to affect intermittency in sub-ion kinetic scale turbulence [45], and it is unclear whether a dominant SH mechanism would have similar effects. In this same vein, studying the stochastic heating of electrons by kinetic scale current sheets [81,151] may prove useful in understanding electron-scale dissipation and intermittency [113,152,153]. Because of the relative infrequency of magnetic reconnection observed by PSP [154], which is likely a result of the Alfvénic nature of inner-heliospheric turbulence [155], we have mostly glossed over reconnection as a dissipation mechanism [27,29]. However, studying the interplay between SH and magnetic reconnection may be an important avenue in understanding the heating of the lower corona, where the relative importance of reconnection heating [156–158] and dissipation of Alfvénic turbulence [90,91,159,160] is still a matter of debate.

In conclusion, there is growing consensus that intermittent heating is fundamental to turbulent dissipation and collisionless dissipation. Our results provide a clear connection between inhomogeneous heating via intermittent dissipation and the dynamics associated with violation of the first adiabatic invariant, namely, particle gyromotion, by turbulence. These results, which connect intermittent heating to Alfvénic turbulence, may help unify understanding of collisionless heating across relatively disparate phenomenological descriptions of turbulence.

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Data availability—The data that support the findings of this article are openly available [108,109,118,119].

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End Matter

Survey of the interval—We consider the interval from PSP’s eighth perihelion encounter from April 8, 2021 from 10:00:00–14:00:00 initially reported as the first encounter with the solar corona [86]. We use the merged PSP search coil and magnetometer (SCaM) data [119], which was sampled at 2 kHz. Figure 4(a) shows a survey of the interval: the magnetic field in S/C coordinates; the magnitude of the field maintains a constant amplitude, which is consistent with large-amplitude Alfvénic turbulence [124]. Figure 4(b) shows the SPANi velocity measurements in the S/C coordinate system and frame. Figure 4(c) shows the total proton $\beta = v_{th}^2/V_A^2$ as well as the perpendicular and parallel contributions; uniformly, $\beta \ll 1$, in agreement with the limits under which the SH diffusion in Eq. (1) is derived [26,65]. Figure 4(d) shows the Alfvénic Mach number U_{sw}/V_A , where U_{sw} is the solar wind speed computed in the heliocentric inertial frame. Figure 4(e) shows θ_{BV} , the angle between the magnetic field and the solar wind velocity in the spacecraft frame for each measurement (the angle is taken between 0 and 90°); the mean angle is $\langle \theta_{VB} \rangle = 25^\circ$. Figure 4(f) shows the reduced magnetic helicity σ_B [161],

computed via a wavelet transform. The red band around 8–10 Hz indicates left-hand circularly polarized ion-scale waves thought to be ICWs [37,38].

Wavelets—We use a sixth-order Morlet wavelet transform [162,163] to identify intervals with circular polarization via σ_B . The full 3D vector flux gate magnetometer measurement (sampled at 292.969 samples per second) is used to compute the wavelet transform in field aligned coordinates, $\tilde{\mathbf{B}}(t, f) = (\tilde{B}_{\perp 1}, \tilde{B}_{\perp 2}, \tilde{B}_{\parallel})$. We follow methods developed in [31,32] to measure QL ||-ICW heating rates based on *in situ* observations of VDFs and wave populations. The reduced magnetic helicity [161] is computed from the wavelet transform as

$$\sigma_B(f, t) = -2 \operatorname{Im}(\tilde{B}_{\perp 1} \tilde{B}_{\perp 2}^*) / (\tilde{B}_{\perp 1}^2 + \tilde{B}_{\perp 2}^2) \quad (\text{B1})$$

and is used to find highly circularly polarized fluctuations, which are likely ion-scale waves propagating parallel to the mean magnetic field and away from the Sun. The expression for Q_{ICW} is found in

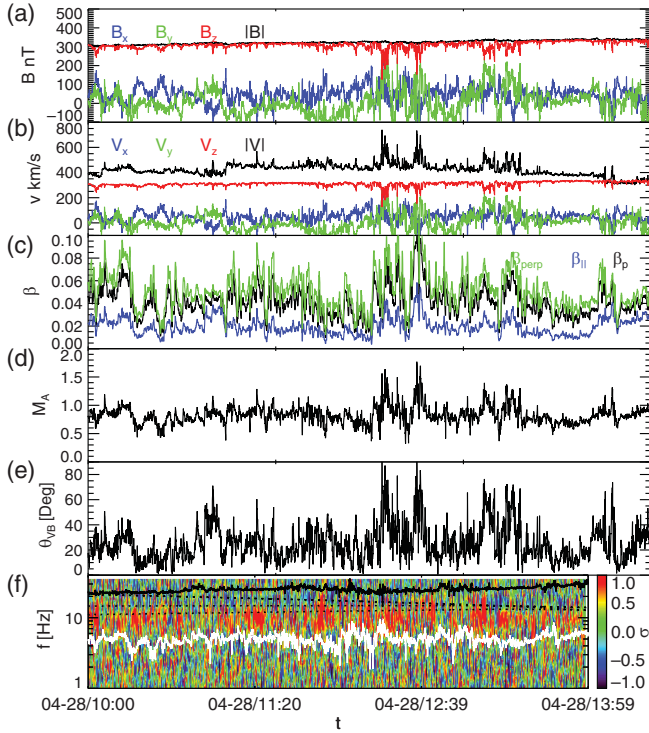


FIG. 4. (a) PSP FIELDS magnetic field in spacecraft coordinates on April 2021, 28 from 10:00:00–14:00:00. (b) PSP SPANi measurements of flow velocity in S/C coordinates. (c) Plasma β , total value as well as perpendicular and parallel decomposition, are shown. (d) Alfvénic Mach number. (e) Angle between magnetic field direction and solar wind flow direction θ_{VB} . (f) Reduced magnetic helicity as a function of frequency and time; dashed lines show spacecraft reaction wheels, and white and black lines show kd_i and $k\rho_i$, respectively.

Eq. (A1) in [32]; the reader is guided to [31,32] for reproducing Q_{ICW} here.

The spectrum of ICW is determined by filtering highly circularly polarized coefficients with $\sigma_B > 0.7$. We compute drifting bi-Maxwellian fits to the SPANi VDFs, assuming a core (c) and drifting beam (b) bi-Maxwellian populations

$$f(v_{\perp}, v_{\parallel}) = \sum_{j=b,c} \frac{n_j \pi^{-3/2}}{w_{j,\perp}^2 w_{j,\parallel}} \exp \left[-\frac{v_{\perp}^2}{w_{j,\perp}^2} - \frac{(v_{\parallel} - v_0)^2}{w_{j,\parallel}^2} \right] \quad (\text{B2})$$

using nonlinear least square fitting methods [31,110,111]. We explored radial basis function interpolation to account for non-Maxwellian effects, but found no significant changes [32].

Five-point increments—Two-point increments cannot resolve the steep scaling of the transition range [97,164], and, thus, we use five-point increments [95] to compute fluctuation amplitudes as

$$\delta B_i(t, \tau) = \Delta_5 \text{pt} B \quad (\text{C1})$$

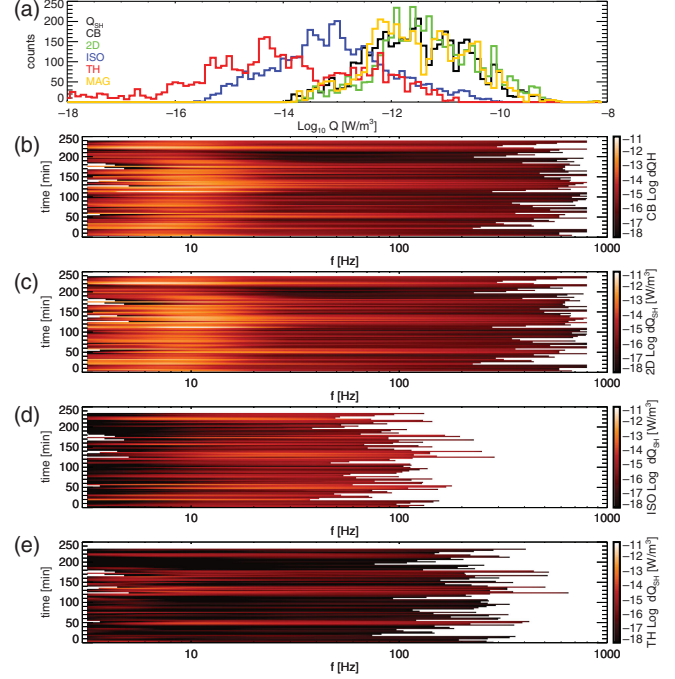


FIG. 5. (a) Probability distribution of heating rates for various assumptions of spectral anisotropy: CB (black), 2D (green), isotropic (blue), and standard Taylor (red). (b)–(e) Respective estimates of dQ_{SH} for assumptions (CB, 2D, ISO, TH) of spectral anisotropy.

$$= \frac{1}{\sqrt{35}} [B_i(t - 2\tau) - 4B_i(t - \tau) + 6B_i(t) \quad (\text{C2})$$

$$- 4B_i(t + \tau) + B_i(t + 2\tau)], \quad (\text{C3})$$

where i refers to the vector components. The KAW correction converts each increment δB_i to a δv_i .

Only two components of the search coil magnetometer (SCM) are available such that the total instantaneous amplitude is computed as $\delta u(f, t) = \sqrt{3/2}(\delta u_v^2 + \delta u_w^2)$, where subscripts v and w denote the functional SCM axes. The rms D_E^{rms} is computed from the time averaged rms $\delta u^{\text{rms}}(f) = \sqrt{\langle |\delta u(f)|^2 \rangle}$. Alternatively, $D_E(t, f)$ is computed through $|\delta u(f, t)|$ in Eq. (1). Evaluating both D_E^{rms} and $\bar{D}_E$ in a scale-dependent manner additionally requires determining $v_{\perp}(f)$. A primary observational complication in this Letter is associating $\delta u(f, t)$ with the appropriate $v_{\perp}(f)$.

Our results show that heating mostly occurs near the ion-transition range, which is well resolved by the flux gate magnetometer (MAG), for which 3D vector measurements are possible. To verify the results from the two-axis SCAm data, we additionally repeated the analysis with the MAG data. No significant changes were found in the total heating rate, as can be seen in Fig. 5(a).

Modified Taylor hypothesis—The Doppler shift equation is generally given by

$$2\pi f = \omega(\mathbf{k}) + \mathbf{k} \cdot \mathbf{V}_{\text{sw}}. \quad (\text{D1})$$

The Taylor hypothesis is obtained simply by dropping the $\omega(\mathbf{k})$ dependence. The modified Taylor hypothesis in Eq. (D2) should be appropriate for the large-amplitude Alfvénic fluctuations, propagating with $\omega = k_{\parallel} V_A$ [98–100]. The hypothesis is valid when the amplitude of the subdominant Elsasser mode is less than the spacecraft's perpendicular speed in the heliocentric inertial frame $|\delta \mathbf{z}^-| \ll V_{sc\perp}$ and additionally $|\delta \mathbf{z}| \ll |\delta \mathbf{z}^+|$ [98], which are generally satisfied in the stream under study.

The condition $k_{\perp} \rho = 1$ corresponds to $k_{\perp} = \Omega_i / v_{\perp}$. Following the modified Taylor hypothesis, with the assumption of critical balance $k_{\parallel} V_A = \delta u \Omega_i / v_{\perp}$, we find

$$v_{\perp} = \frac{\Omega_i}{2\pi f} \left[\delta u \left(1 + \frac{V_{\text{sw}}}{V_A} \cos \theta_{VB} \right) + V_{\text{sw}} \sin \theta_{VB} \right], \quad (\text{D2})$$

where we use $\delta u^{\text{rms}}(f)$ to convert from f to $v_{\perp}$.

We also consider a purely two-dimensional (2D) perpendicular spectra with $k_{\parallel} = 0$:

$$v_{\perp}^{2D} = \frac{\Omega_i}{2\pi f} V_{\text{sw}} \sin \theta_{VB};$$

an isotropic spectra $k_{\perp} = k_{\parallel}$, which yields

$$v_{\perp}^{\text{iso}} = \frac{\Omega_i}{2\pi f} (V_A + V_{\text{sw}} \cos \theta_{VB} + V_{\text{sw}} \sin \theta_{VB});$$

and the standard Taylor hypothesis (TH)

$$v_{\perp}^{\text{TH}} = \frac{\Omega_i}{2\pi f} V_{\text{sw}}.$$

Figure 5 shows results using various definitions of $v_{\perp}$ based on different assumptions of the spectra. Assuming an entirely perpendicular 2D spectra yields similar heating rates as a CB spectra; the use of isotropic or a standard TH gives much lower heating rates. Figures 5(b)–5(d) show the corresponding dQ_{SH} calculations for the CB, 2D, isotropic, and standard TH analysis. The CB is the same as in Fig. 3 but replotted on the same scale for comparison with the other spectra. The clear similarities in dQ_{SH} between the CB and 2D data, and obvious differences in comparison with isotropic spectra and standard TH approaches, suggest that our analysis should hold as long as the turbulent spectra are strongly perpendicular.

LET—To estimate the total energy cascade rates ϵ^T , we use the local-energy transfer [88,89,135]. Standard Politano-Pouquet [165] estimates of the magneto-hydrodynamic cascade rate require mixed third-order moment of increments of the Elsasser variables $\Delta \mathbf{z}^{\pm} = \Delta \mathbf{v} \pm \Delta \mathbf{B} / \sqrt{\mu_0 \rho_0}$ given as

$$\langle \epsilon^{\pm} \rangle = -\frac{3}{4} \langle |\Delta \mathbf{z}(t, \tau)^{\pm}| \Delta \mathbf{z}(t, \tau)^{\mp} \rangle \frac{1}{V_{\text{sw}} \tau}. \quad (\text{E1})$$

The unaveraged $\epsilon^T = (\epsilon^+ + \epsilon^-)/2$ can be interpreted as an estimate of the local cascade rate associated with individual increments [89]. The distributions in Fig. 1 show distributions of the local cascade rate at three lag values (3.5, 49, and 168 s) corresponding to the inertial range over the entire 4-h interval. The distributions are roughly consistent with the dissipation associated with intermittent stochastic heating. Intermittency has been shown to directly contribute to nonhomogeneity of the turbulent cascade rate [135], which is similar to our present results that highlight the importance of intermittency in stochastic heating.