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CHARGE STATES OF SOLAR ENERGETIC IRON: NONEQUILIBRIUM CALCULATION WITH SHOCK-INDUCED ACCELERATION

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ABSTRACT

Observed charge states of solar energetic ions by *Solar, Anomalous, and Magnetospheric Particle Explorer* and *Advanced Composition Explorer* from recent large solar events in 1992 and 1997 show an energy-dependent feature that apparently has not been seen in earlier events, e.g., in 1978–1979. The mean charge state of iron, and to a lesser extent other heavy ions, was observed to increase with energy from ∼0.5 to ∼50 MeV nucleon^{-1}. We present a possible explanation for this new feature and discuss potential implications for models of particle acceleration. Our calculation suggests that this energy-dependent feature is a result of the charge-changing and energy-changing processes taking place concurrently and can be absent when the processes act over different timescales.

Subject headings: acceleration of particles — Sun: flares — Sun: particle emission

1. INTRODUCTION

The observed ionic charge states of solar energetic particles (Luhn et al. 1984, 1985; Leske et al. 1995; Mason et al. 1995; Tylka et al. 1995; Oetliker et al. 1997; Mazur et al. 1999; Möbius et al. 1999) are considered to be sensitive probes of the coronal plasma conditions, e.g., its mean electron density and temperature. This is largely based on the notion of a "frozen-in" seed population. According to this notion, the measured charge state distribution is established via charge-changing processes followed by acceleration and transport processes, which do not alter the distribution in any significant way. The use of equilibrium ionization temperatures to infer the coronal plasma temperature from the measured mean charge suggests different temperatures for different ion species (Luhn & Hovestadt 1985; Oetliker et al. 1997).

The charge-changing processes are due to ion-electron collisions and photoionization in the hot coronal plasma. In large (or gradual) solar events, observational evidence points to an acceleration by large-scale shocks in the outer corona and in the solar wind plasma (e.g., Lin 1987; Reames 1990; Boberg, Tylka, & Adams 1996). Propagation effects, due to the lower temperature and density of the solar wind, are not expected to further alter the charge states of the seed population (Hovestadt et al. 1984; Luhn & Hovestadt 1987).

Recent measurements (Oetliker et al. 1997; Mazur et al. 1999; Möbius et al. 1999) of the charge states of a number of solar energetic ions (up to Fe) over a wide energy range (up to a few tens of MeV nucleon^{-1}) show an energy dependence of the inferred mean charge $\langle q \rangle$. Higher energy ions, especially iron, seem to have a higher $\langle q \rangle$. This feature, seen in both the 1992 and 1997 events, was not reported in earlier observations from 1978–1979 (Luhn et al. 1984), which covered a smaller energy interval. Oetliker et al. (1997) suggested that either different (frozen-in) seed populations (i.e., coronal vs. solar wind) with different electron temperatures or a rigidity-dependent acceleration mechanism that tends to favor lower rigidity (higher $\langle q \rangle$) ions could be responsible for the energy-dependent feature seen in the 1992 events.

Luhn & Hovestadt (1985) examined the effects of non-Maxwellian electron-velocity distribution and plasma heating on the observed charge states and the inferred temperature. They concluded that neither effect was strong enough to explain the variation in the temperature inferred for different ion species. (It was in this work that the first qualitative suggestion regarding the possible role of equilibration and acceleration timescales in the observed charge states was made.)

In a steady state model, Kurganov & Ostryakov (1991) explored the effects of shock-induced acceleration and charge-changing processes on the energy spectra and charge distributions of light solar energetic ions. Their analytic model is similar to the one presented here. Their treatment of the charge-changing processes, however, is oversimplified.

On the basis of a nonequilibrium calculation without acceleration, Ruffolo (1997) examined the roles of shock heating and further charge stripping. He found neither effect able to explain the different apparent temperatures for different ion species. Using measured charge states, he inferred an upper limit ($\lesssim 3 \times 10^9$ cm^{-3} s) on the electron density × acceleration time product. The order of this (upper) limit suggests that the acceleration could not take place in coronal loops, with assumed densities of $\sim 10^{11}$–10^{12} cm^{-3}, but must occur on open magnetic field lines, where the density begins to drop rapidly—reaching $\sim 10^5$–10^6 cm^{-3} at heights $\sim r_\odot$.

In this Letter, we present a quantitative explanation for the observed energy dependence of $\langle q \rangle$ based on the timescales for the charge-changing (due to ionization-recombination) and shock-induced acceleration. When the timescales are comparable, the two processes tend to equilibrium in a dynamic, time-dependent fashion, and the interplay between the two cannot be ignored. This equilibrium will then reflect an accelerated seed population with not only its own characteristic temperature but also with its nonthermal energy spectrum. This interplay between charge-changing and acceleration processes can give rise to an energy-dependent $\langle q \rangle$. We will show that it is only when the timescales are very different that the concept of a frozen-in seed population that is *subsequently* accelerated is applicable to the seed population. In this latter case, $\langle q \rangle$ will largely reflect the preaccelerated as well as the accelerated equilibrated population.

We concentrate here on solar energetic iron for two different reasons. First, as discussed by Oetliker et al. (1997), Aellig et

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al. (1998), and Gloeckler et al. (1999), iron charge states are expected to be quite sensitive to temperature variations, and there is evidence for an energy dependence in these measurements to date. Secondly, the set of atomic physics parameters needed to calculate all the charge states of iron, as compiled by Arnaud & Raymond (1992), is relatively complete and data-tested. Further charge stripping due to photoionization (e.g., Mullan & Waldron 1986) is ignored.

2. CHARGE-STATE BALANCE EQUATION WITH ACCELERATION

The nonequilibrium balance equation of our model includes stochastic (preacceleration) and shock-induced acceleration in addition to source and sink coupling terms. In a spatially homogeneous acceleration region, the shock is idealized as a finite, plane shock with its normal parallel to the flow of the coronal plasma (e.g., Schlickeiser 1985). The balance equation is

$$\frac{\partial f^q}{\partial t} = \frac{1}{p^2}\frac{\partial}{\partial p}\left[p^2 D_{pp}(p)\frac{\partial f^q}{\partial p}\right] - \frac{1}{p^2}\frac{\partial}{\partial p}\left[p^2\frac{dp}{dt}f^q\right] - \frac{f^q}{t_{\rm esc}} + Q^q_{\rm ir}(p,t), \tag{1}$$

where $f^q(p, t)$ is the phase-space density function of ion with charge q as a function of momentum p at time t. Because propagation effects are assumed negligible, explicit spatial variations of $f^q(p, t)$ are ignored.

The source and sink function $Q^q_{\rm ir}(p,t)$, due to ionization and recombination, at a given temperature T is given by

$$Q^q_{\rm ir}(p,t) = n_e\{R_i^{q-1}(T)f^{q-1} + R_r^{q+1}(T)f^{q+1} - [R_i^q(T)f^q + R_r^q(T)f^q]\}, \tag{2}$$

where n_e is the mean electron density and $R_i(T)$ and $R_r(T)$ are the temperature-dependent ionization and recombination rates. To estimate these rates, we use the formulation[2] and various atomic physics parameters of Arnaud & Raymond (1992). For the ionization rates, their formulation includes direct ionization (due to collisions with electrons) and autoionization (as noncollisional processes). The recombination rates include radiative and dielectronic contributions.

In equation (1), the term $f^q/t_{\rm esc}$ describes the rate of escape from the shock acceleration region, where $\mathcal{P}_{\rm esc} = t_s/t_{\rm esc} = 4v_s/\eta v$ is the escape probability per shock crossing (e.g., Forman & Webb 1985; Drury 1983); v_s is the shock's propagating speed, v is the ion's speed, and η is the shock's compression ratio. In the rest frame of the plasma, $v_s = (3k_{\rm B}T_e/m_0)^{1/2}c$, where T_e is the electron temperature, $k_{\rm B}$ is Boltzmann's constant, m_0 is the nucleon's rest mass, and c is the speed of light. The cycle time t_s, or time to complete a crossing of the shock front, is given by $t_s = 4(1+\eta^{-2})D_\parallel/v_s v$. The quasi-linear diffusion coefficient $D_\parallel(p)$ along the magnetic field line is expressed as (Jokipii 1971)

$$D_\parallel(p) = D_0 P^{2-\alpha}\beta\frac{B_0}{B_s}, \tag{3}$$

where D_0 is a constant expressing the strength of the diffusion, P is the ion's rigidity, α is the spectral index of the magnetic turbulence, and $\beta = v/c$. B_0/B_s is the ratio of the strength of the magnetic field at 1 AU to that at the acceleration site, typically included to ensure that the particle remains tied to the field line (Jokipii & Morfill 1987).

The rate of momentum gain per shock crossing is written as

$$\frac{dp}{dt} = \frac{\zeta v_{\rm sw}^2}{4D_\parallel}\,p, \tag{4}$$

where $\zeta = (\eta-1)/6\eta$ and $v_{\rm sw}$ is the solar wind speed. In the hard-sphere approximation, the diffusion coefficient in momentum space $D_{pp}(p)$ is related to $D_\parallel$ as

$$D_{pp}(p) = \frac{v_{\rm sw}^2}{9D_\parallel}\,p^2. \tag{5}$$

For the solution of equation (1), the initial $f^q(p, t=0) = f^q_0(p)$ is assumed Maxwellian in p characterized by the electron temperature T_e. The initial distribution in q is taken to be that for an equilibrated distribution corresponding to some $T < T_e$. Due to the diffusive nature of the acceleration and the diffusion-like behavior of the coupling terms (Ivanov, Kukushkin, & Lisitsa 1987), however, the nonthermal, steady state energy spectrum of $\langle q\rangle$ is insensitive to the exact form of either distribution of the seed population. The boundary conditions in p are $f^q(p=p_0, t) = f^q_0(p=p_0)$ and $\partial f^q(p,t)/\partial p|_{p=p_0} = \partial f^q_0(p)/\partial p|_{p=p_0}$, where p_0 is the lower limit for p. Shock acceleration is assumed to affect only those ions preaccelerated by diffusing in momentum space and attaining a momentum $p > p_{\rm A}$, where $p_{\rm A}$ is the Alfvén momentum given by $p_{\rm A} = m_0 B_s/(\mu n_e)^{1/2}$, μ being the magnetic permeability constant.

3. CALCULATION FOR SOLAR ENERGETIC IRON

To arrive at an energy spectrum for $\langle q\rangle$, the 27 coupled balance equations (eq. [1]) are numerically integrated starting from a thermal seed population as described above. The steady state spectrum $\langle q\rangle(p)$ is then calculated from the steady state solution $f^q(p)$ as $\langle q\rangle(p) = \sum_q qf^q(p)/\sum_q f^q(p)$ and similarly for the second moment of q.

Before presenting the calculation for Fe, we address the question of comparable timescales in the equilibration and acceleration processes. In Figure 1a, the equilibration curve depicts the solution to the balance equations with no acceleration, i.e., $\partial f^q/\partial t = Q^q_{\rm ir}(p,t)$, with $Q^q_{\rm ir}(p,t)$ given by equation (2). The acceleration curves 1–4 in Figure 1b are arrived at as follows. The characteristic time t_c it takes an ion to be accelerated from momentum $p_i \approx p_{\rm A}$ to a momentum p_c via the mechanism of diffusive shock acceleration can be estimated from $t_c \sim \int_{p_i}^{p_c}(dp/dt)^{-1}dp$, with dp/dt from equation (4). In kinetic energy and for the parameters discussed below, for $p_c = 10p_{\rm A}$ this range in p corresponds to ≈0.25–25 MeV nucleon^{-1}.

Curves 1–4 are arrived at with the following set of physical parameters assumed to be associated with an active coronal region in a typical large solar event: the characteristic electron

[2] In their expressions for the rates, Arnaud & Raymond (1992) as well as Arnaud & Rothenflug (1985) use analytic approximations for the exponential integral function $f_1(x)\exp(-x) = E_1(x) \equiv \int_1^\infty \exp(-tx)/t\,dt$ and the related function $f_2(x)\exp(-x) = \int_1^\infty \exp(-tx)\ln(t)/t\,dt$. The approximation for $E_1(x)$ given in Abramowitz & Stegun (1972) is found to be more accurate. A reasonably accurate approximation for $f_2(x)$ is given by $f_2(x)\exp(-x) \approx 1/2[E_1(x)]^2$.

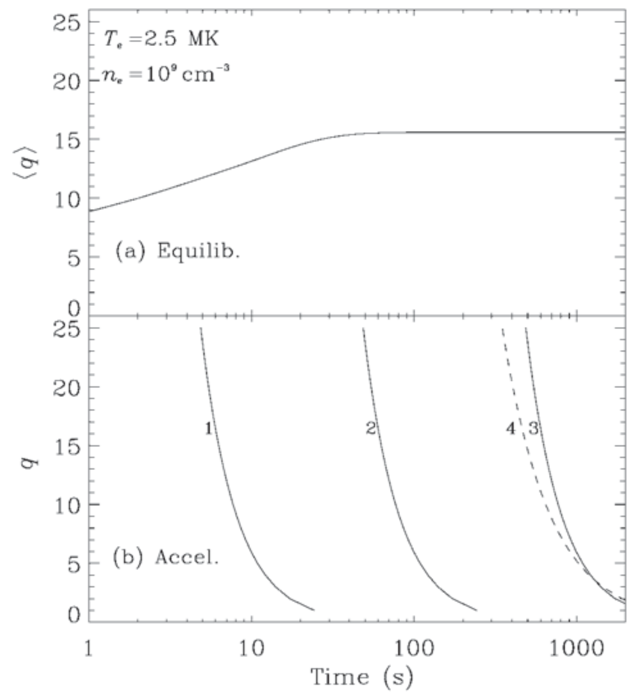

FIG. 1.—(*a*) Equilibration timescale t_{eq}, where the equilibration curve for $\langle q \rangle$ is due to the charge-changing processes only. The initial distribution in q corresponds to a temperature of 1 MK. (*b*) Acceleration timescale t_c, where the acceleration curves are due to the momentum-changing processes only. Curve 1 depicts a weaker diffusive process than curves 2 and 3, respectively. Dashed curve 4 is similar to curve 2 except it depicts a stronger rigidity-dependent acceleration process than curve 2. See text for further discussion of the set of parameters used in the calculation.

temperature T_e is taken to be 2.5 MK (T for the initial distribution in q is 1 MK), and the density n_e is 10^9 cm^{-3} (i.e., the seed population is taken to be that of a flare-heated coronal material accelerated at the flare site on open magnetic field lines at an altitude of $\lesssim 0.1 r_\odot$). The strength of the local magnetic field at the acceleration site B_s is 10^2 G. The benchmark strength of the local diffusion coefficient D_0 is 10^{21} cm^2 s^{-1}, and the benchmark spectral index of the magnetic turbulence α is 3/2 (i.e., a Kraichnan spectrum). (Note that outside the site's assumed physical parameters [i.e., T_e, n_e, and B_s], D_0 and α are the only two adjustable [physical] parameters in this model.)

Acceleration curve 2 in Figure 1*b* depicts the benchmark timescale using the above cited parameters. It is seen that the acceleration timescale t_c for $q > 10$ is $\lesssim 10^2$ s, which is comparable to the timescale for equilibration t_{eq}. The timescale is only a weak function of q (or rigidity) for $q > 10$. Curve 3 is arrived at using the same parameters as in curve 2 except that D_0 is now increased by a factor of 10, simulating a strongly diffusive acceleration. Similarly, curve 1 has the same parameters as curve 2 except for D_0, which is decreased by a factor of 10 from curve 2. Curve 4 shares the same set of parameters as curve 2, including D_0, except now α is decreased to 4/3, simulating a stronger rigidity-dependent acceleration. Clearly from Figure 1 the acceleration timescales in curves 3 and 4 (strongly diffusive and strong rigidity dependence, respectively) are much greater than the equilibration time. For curves 1 and 2 (weakly to moderately diffusive), the acceleration timescale is comparable to or less than the equilibration time.

In light of Figure 1 and the above discussion, Figure 2 shows

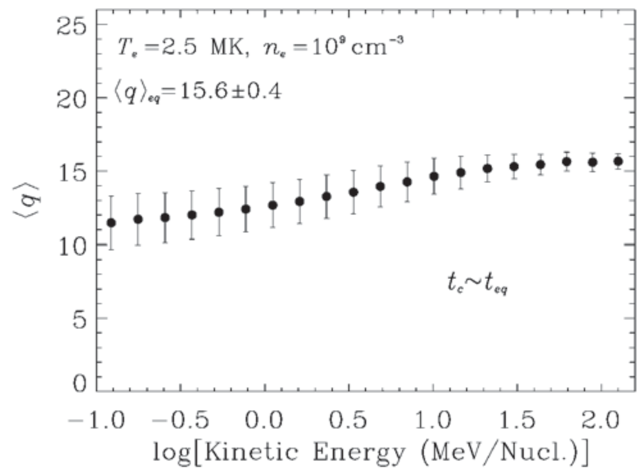

FIG. 2.—Steady state solution for $\langle q \rangle$, as a function of kinetic energy, from the coupled set of eq. (1). The equilibrium value $\langle q \rangle_{eq}$ is due only to ionization and recombination processes. Time to reach steady state t_∞ is typically greater than $t_c > t_{esc}$. This energy spectrum of $\langle q \rangle$ corresponds to curve 2 of Fig. 1*b*.

the steady state $\langle q \rangle \pm (\langle q^2 \rangle)^{1/2}$ as a function of kinetic energy for the benchmark timescale for acceleration, i.e., curve 2. The equilibrium value for $\langle q \rangle$, i.e., without acceleration, is denoted by $\langle q \rangle_{eq}$. Time to reach steady state t_∞, while both q and p dependent, is typically greater than $t_c > t_{esc}$. The energy-dependent feature is quite noticeable and is similar to that reported by Oetliker et al. (1997) for the 1992 events. Two significant characteristics of this energy dependence can also be seen here. First, $\langle q \rangle$ seems to approach $\langle q \rangle_{eq}$ as the energy increases but does not overtake it, i.e., for a given temperature $\langle q \rangle_{eq}$ is an upper limit for $\langle q \rangle$. Second, $(\langle q^2 \rangle)^{1/2}$ seems to decrease appreciably with increasing energy, approaching $(\langle q^2 \rangle_{eq})^{1/2}$. The energy spectrum of $\langle q \rangle$ corresponding to curve 1 is qualitatively similar to that for curve 2 except that it is lower by a few charge states.

In Figure 3 the energy spectrum of $\langle q \rangle$ that corresponds to curve 3 is shown (i.e., strongly diffusive case), where $\langle q \rangle$ and $(\langle q^2 \rangle)^{1/2}$ are now weak functions of energy. The energy spectrum of $\langle q \rangle$ corresponding to curve 4 is identical to that for curve 3. (Simulated charge distributions and energy spectra are reported in Barghouty & Mewaldt 1999.)

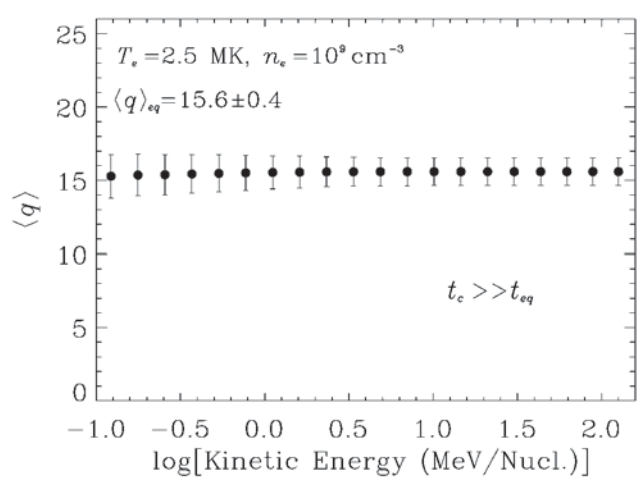

FIG. 3.—Similar to Fig. 2, except the energy spectrum of $\langle q \rangle$ corresponds to curve 3 of Fig. 1*b*.

4. DISCUSSION

The calculation in Figures 1–3 above suggests that when the timescales for equilibration t_{eq} and acceleration t_c tend to be comparable, $\langle q \rangle$ tends to assume an energy-dependent profile (curves 1–2 and Fig. 2). In this case, the interplay between equilibration and acceleration is inconsistent with the concept of a frozen-in, then accelerated seed population. Only when the timescales are very different (curves 3–4 and Fig. 3) and $\langle q \rangle$ assumes the preaccelerated, equilibrium value does the frozen-in concept become applicable.

The calculation has shown that a rigidity-dependent acceleration process with $2 - \alpha \leq 0.5$ (curve 2) is consistent with an energy-dependent $\langle q \rangle$. However, when the rigidity dependence is too strong, i.e., $2 - \alpha > 0.5$, it tends to increase the characteristic acceleration time relative to the equilibration time (curve 4), resulting in a seed population that is only subsequently accelerated. This suggests that the energy spectrum of $\langle q \rangle$ may be a more sensitive measure of the characteristic acceleration time than the rigidity dependence of the acceleration process.

The behavior of the second moment of q with energy may also be as significant. While this behavior is related to the ionic shell structure of Fe, the diffusive nature of the acceleration process may also play a part. For Fe, with its relatively large number of charge states, one can numerically treat q as a continuous variable on par with p, so that equation (1) describes the dynamic evolution (diffusion-like) of the density function f in q and p space (Ivanov et al. 1987). Since f is conserved in this two-dimensional space, the distribution will be wider along q for small p, and vice versa. This suggests that the observed energy spectrum of $\langle q^2 \rangle$ can serve as a signature of the diffusive nature of the acceleration process.

According to Ruffolo (1997), an energy-dependent $\langle q \rangle$ places more severe constraints on his inferred upper limit for the product $n_e \times t_c$. Using a characteristic t_c of $\lesssim 10^2$ s and density of 10^9 cm^{-3}, our calculation suggests a more relaxed upper limit of $\sim 10^{11}$ cm^{-3} s (compared to 3×10^9 cm^{-3} s). However, even the more relaxed upper limit for $n_e \times t_c$ still does not lead to a t_c of $\sim 10^{-2}$ s, which seems to be required (Ruffolo 1997) for acceleration in coronal loops in which the density is $\sim 10^{11}$–10^{12} cm^{-3}.

Luhn & Hovestadt (1985) and Ruffolo (1997) concluded, even though acceleration was not explicitly taken into account in their respective nonequilibrium calculations, that plasma heating had little effect on the charge states. Measurements from the 1997–1998 events at widely separated energies ($\sim$1 MeV nucleon^{-1} by Möbius et al. 1999 and $\sim$50 MeV nucleon^{-1} by Mazur et al. 1999) suggest a mean charge for Fe higher than the characteristic $\langle q \rangle_{eq}$. A higher temperature can result in a higher $\langle q \rangle_{eq}$. Our finding here that $\langle q \rangle \leq \langle q \rangle_{eq}$ at a given temperature, even when the energy dependence of $\langle q \rangle$ is accounted for, suggests that photoionization may also be responsible for further charge stripping. The effect of photoionization on the energy spectrum of $\langle q \rangle$ in the contexts of both equilibrium and nonequilibrium calculations is currently being tested.

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REFERENCES

Abramowitz, M., & Stegun, I. A., eds. 1972, Handbook of Mathematical Functions (New York: Dover)
Aellig, M. R., et al. 1998, J. Geophys. Res., 103, 17,215
Arnaud, M., & Raymond, J. 1992, ApJ, 398, 394
Arnaud, M., & Rothenflug, R. 1985, A&AS, 60, 425
Barghouty, A. F., & Mewaldt, R. A. 1999, Proc. 26th Int. Cosmic-Ray Conf. (Salt Lake City), in press
Boberg, P. R., Tylka, A. J., & Adams, J. H., Jr. 1996, ApJ, 471, L65
Drury, L. O'C. 1983, Rep. Prog. Phys., 46, 973
Forman, M. A., & Webb, G. M. 1985, in Collisionless Shocks in the Heliosphere: A Tutorial Review, ed. Stone, R. G. & Tsurutani, B. T. (Geophys. Monog. 34; Washington DC: American Geophysical Union), 91
Gloeckler, G., et al. 1999, Geophys. Res. Lett., 26, 157
Hovestadt, D., Gloeckler, G., Klecker, B., & Scholer, M. 1984, ApJ, 281, 463
Ivanov, V. V., Kukushkin, A. B., & Lisitsa, V. S. 1987, Soviet J. Plasma Phys., 13, 774
Jokipii, J. R. 1971, Rev. Geophys. Space Phys., 9, 27
Jokipii, J. R., & Morfill, G. 1987, ApJ, 312, 170
Kurganov, I. G., & Ostryakov, V. M. 1991, Soviet Astron. Lett., 17, 77
Leske, R. A., Cummings, J. R., Mewaldt, R. A., Stone, E. C., & von Rosenvinge, T. T. 1995, ApJ, 452, L149
Lin, R. P. 1987, Rev. Geophys., 25, 676
Luhn, A., & Hovestadt, D. 1985, Proc. 19th Int. Cosmic-Ray Conf. (La Jolla), 4, 245
———. 1987, ApJ, 317, 852
Luhn, A., et al. 1984, Adv. Space Res., 4, 161
———. 1985, Proc. 19th Int. Cosmic-Ray Conf. (La Jolla), 4, 241
Mason, G. M., Mazur, J. E., Looper, M. D., & Mewaldt, R. A. 1995, ApJ, 452, 901
Mazur, J. E., Mason, G. M., Looper, M. D., Leske, R. A., & Mewaldt, R. A. 1999, Geophys. Res. Lett., 26, 173
Möbius, E. M., et al. 1999, Geophys. Res. Lett., 26, 145
Mullan, D. J., & Waldron, W. L. 1986, ApJ, 308, L21
Oetliker, M., et al. 1997, ApJ, 477, 495
Reames, D. V. 1990, ApJ, 358, L63
Ruffolo, D. 1997, ApJ, 481, L119
Schlickeiser, R. 1985, A&A, 143, 431
Tylka, A. J., Boberg, P. R., Adams, J. H., Jr., Beahm, L. P., Dietrich, W. F., & Kleis, T. 1995, ApJ, 444, L109