PPPL-2457

UC20-G

120
10-1-87

37 (2) (5)

I-32032

OR 0314-X
PPPL-2457

PPPL--2457

DE88 000117

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JUNE 1987

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PREPARED FOR THE U.S. DEPARTMENT OF ENERGY,
UNDER CONTRACT DE-AC02-76-CEO-3073.

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Printed in the United States of America

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Springfield, Virginia 22161

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Inverse Problem for Incremental Synchrotron Radiation

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Abstract

Significantly more information is available from synchrotron emission from a plasma when the plasma is purposefully disturbed. An inverse problem, to deduce properties of the disturbance given time-dependent radiation data, is proposed. The fast time response of radiation detectors is fully exploited by this approach. A special case of interest, perpendicular observation of a steady-state plasma, lends itself to an analytic inversion.

MASTER

A magnetically confined plasma is a copious emitter of synchrotron radiation. In fact, the radiated power is a matter of concern in the design of high temperature nuclear fusion reactors. Here, however, we concern ourselves with the possibility of diagnosing the plasma by means of this radiation.

A method is proposed by which plasma synchrotron emission from tokamaks may be used to infer the two-dimensional electron distribution function $f(p_{\parallel}, p_{\perp})$, where $p_{\parallel}$ and $p_{\perp}$ refer to the electron momentum, respectively, parallel and perpendicular to the magnetic field. The straightforward approach would be to deduce $f(p_{\parallel}, p_{\perp})$ from the 2-D pattern $R(\omega, \theta)$ of radiation emitted at frequency ω into angle θ , where θ measures the angular deviation from purely perpendicular observation of the magnetic field. (We observe the tokamak in the vertical plane that includes the tangent to the magnetic field $\mathbf{B}$, so the strength of $\mathbf{B}$ may be assumed constant and known.) This approach, however, is impractical because it requires the deployment of an array of microwave detectors to resolve the θ -dimension. In practice, possibly only one detector is available, and, while it has been used to place useful constraints¹⁻⁵ on $f(p_{\parallel}, p_{\perp})$, or to deduce elegantly⁶ a 1-D f , a full inversion is not possible.

We solve instead a simpler, but, as we show, very useful, problem. Suppose that rf or other power induces in the plasma a momentum-space flux $\Gamma(\mathbf{p}, t)$, and we measure the *incremental* radiation emitted by the perturbed distribution. We are at liberty, however, to modulate or otherwise control the time-dependence here; so suppose we impose an impulse $\Gamma = \mathbf{S}(\mathbf{p})\delta(t)$. The incremental radiation now decays in time as the electrons suffer collisions, obeying laws we think we know, so that this time decay is revealing. For example, radiation associated with fast electrons would decay slowly; for nonrelativistic electrons the decay time goes as $1/p^3$. By analyzing the 2-D impulse response $R(\omega, t)$, we hope to deduce the 2-D wave-induced flux $\mathbf{S}(\mathbf{p})$. Only one observation angle is necessary here. Of course, if more detectors were available, the multiplication in information would be the same, and the additional information might be used to uncover spatial dependencies. (In this paper, we take f to be homogeneous in the region where the incremental radiation is produced.)

Note that while frequency and time are formally conjugate variables, a huge separation in time scales permits them to be treated as independent variables; hot electrons in a reactor slow down in hundreds of milliseconds, while the frequencies to be measured are hundreds of gigahertz. Radiation detectors with microsecond response time are readily available, so that $\sim 10^5$ time points are available, in principle, in a decay time. Thus, the potential exists for a huge multiplication in the amount of processible information.

Together with theories of wave-damping and quasilinear diffusion, $\mathbf{S}(\mathbf{p})$ may be used to infer $f(\mathbf{p})$, at least where $\mathbf{S}(\mathbf{p})$ is finite. However, details of $\mathbf{S}(\mathbf{p})$ are often of immediate interest. For example, the efficiency of current drive by lower hybrid waves depends critically on the speed of the resonant electrons.⁷ Suppose the plasma current were maintained by these waves. By delivering through the waveguides an incremental impulse of power, under otherwise steady conditions, the resonant electron velocity distribution may be available from the incremental radiation, possibly providing useful feedback for waveguide phasing. Moreover, in such a manner one might resolve the outstanding question of the so-called *spectral gap* observed in recent experiments.

Consider the incremental radiation R due to an externally imposed momentum-space flux $\Gamma(\mathbf{p}, t)$. The distribution function f can be written as $f = f_M(1 + \phi_B + \phi)$, where f_M is a Maxwellian distribution, ϕ_B describes the deviation from Maxwellian of the background distribution, and ϕ describes the distribution specifically associated with the source Γ . Assume that f obeys the linearized Fokker-Planck equation, which for the term we are interested in may be written as

$$f_M \partial \phi / \partial t - C(\phi) = -\nabla_{\mathbf{p}} \cdot \Gamma(\mathbf{p}, t) \equiv -\delta(t) \nabla_{\mathbf{p}} \cdot \mathbf{S}(\mathbf{p}), \quad (1)$$

where C is a collision term and steady state (no dc electric field) has been assumed. Then the incremental radiation due to Γ may be written for an optically thin plasma as

$$R(\omega, \tau) = \int d^3 p f_M \phi(\mathbf{p}, \tau) I(\omega, \theta, \mathbf{u}) = \int d^3 u \psi(\omega, \mathbf{u}, \tau) Q(\mathbf{u}), \quad (2)$$

where we employ normalized momentum, $\mathbf{u} = \mathbf{p}/mc$, and normalized time, $\tau = \nu_c t$, with collision frequency $\nu_c = n q^4 \log \Lambda / 4 \pi m \epsilon_0^2 c^3$, and we define $Q(\mathbf{u}) \equiv -(mc)^3 \nabla_{\mathbf{p}} \cdot \mathbf{S}(\mathbf{p})$. Here I is the radiated power into angle θ and frequency ω of a single electron at momentum $\mathbf{p}$, and ψ is the Green's function for the radiation response, i.e., ψ solves the relativistic Fokker-Planck adjoint equation, written for superthermal excitation in the high-velocity limit⁷ as

$$\frac{\partial \psi}{\partial \tau} + \frac{1}{u^3} \left(\gamma^2 u \frac{\partial \psi}{\partial u} - \gamma \frac{1+Z}{2} \frac{\partial}{\partial \mu} (1-\mu^2) \frac{\partial \psi}{\partial \mu} \right) = 0, \quad (3)$$

with the initial condition $\psi(\omega, \theta, \mathbf{u}, \tau = 0) = I(\omega, \theta, \mathbf{u})$, where Z is the ion charge state and we defined $\gamma^2(u) \equiv 1 + u^2$ and $\mu \equiv p_{\parallel}/p_{\perp}$. Separate ψ and the initial conditions into Legendre harmonics, ψ_k and I_k , and the resulting equations for the ψ_k may be integrated along characteristics to obtain

$$\psi_k(\omega, u, \tau) = I_k(\omega, \rho) \left[\frac{\rho}{u} \left(\frac{\gamma(u)+1}{\gamma(\rho)+1} \right) \right]^{\alpha_k}, \quad (4)$$

where the characteristic $\rho(u, \tau)$ solves $G(\rho) = G(u) - \tau$, where $G(u) = \int_0^u (x^2/(1+x^2)) dx = u - \tan^{-1} u$, and where $\alpha_k \equiv (1+Z)k(k+1)/2$. For u nonrelativistic, $\rho \rightarrow (u^3 - 3\tau)^{1/3}$, indicating that after a time $\tau = u^3/3$, electrons initially with speed u have slowed down to $u = 0$, at which point they no longer radiate. Although Eq. (3) is derived rigorously for superthermal u only, it may be applied here universally, since radiation from slowed-down, less energetic electrons is small compared to initially, and the terms neglected in Eq. (3), energy diffusion, merely act to thermalize, so that, in any event, the incremental radiation vanishes.

Denoting the Legendre components of $Q(\mathbf{u})$ by $Q_k(u)$, we can write Eq. (2) in the form

$$R(\omega, \tau) = \sum_k \frac{4\pi}{2k+1} \int_0^\infty u^2 du \psi_k(\omega, u, \tau) Q_k(u). \quad (5)$$

The task here is to invert Eq. (5), i.e., solve for the Q_k given R . Note that Eq. (5) is a Fredholm equation of the first kind, with a 4-D kernel. Neither the existence nor the

uniqueness of the solution is guaranteed. In fact, of particular interest is the null space and range of the kernel, which maps Q into R . The null space, spanned by those $Q_k(u)$ which are mapped into $R = 0$, represents the irreducible ambiguity in deducing Q from R . The range of ψ is spanned by those $R(\omega, \tau)$ for which solutions Q exist. Experimental data falling outside the range cast doubt either on the validity of the data or on the applicability of the assumptions that form the basis of the theory.

The numerical inversion of Eq. (5) is simplified by exploiting the properties of the kernel ψ ; in particular, note that Eq. (5) is of the form

$$R(\omega, \tau) = \sum_k \int_0^\infty d\xi \Psi_k(\omega, \xi - \tau) q_k(\xi), \quad (6)$$

where $\xi = G(u)$. Let $\tilde{\Psi}_k(\omega, s) \leftrightarrow \Psi_k(\omega, \xi)$, $\tilde{R}(\omega, s) \leftrightarrow R(\omega, \tau)$, and $\tilde{q}_k(s) \leftrightarrow q_k(\xi)$ be Fourier transform pairs. Then Eq. (6) can be put in the form

$$\tilde{R}(\omega, s) = \sum_k \tilde{\Psi}_k(\omega, -s) \tilde{q}_k(s), \quad (7)$$

which simplifies Eq. (6) to a set, parameterized by s , of Fredholm summations of the first kind, but with 2-D, rather than 4-D, kernels.

An important special case is purely perpendicular observation ($\theta = 0$), i.e., observation along a line of sight parallel to the tokamak major axis. The radiation intensity, say, for ordinary polarization (i.e., with $\mathbf{E}$ vector parallel to the magnetic field) may then be written as

$$I(\omega, \theta = 0, \mathbf{u}) = \sum_n \frac{e^2 \omega^2}{2\pi c} \frac{u^2 \mu^2}{\gamma^2} J_n^2 \left(n \frac{u}{\gamma} (1 - \mu^2)^{1/2} \right) \delta(\omega - n\omega_c/\gamma), \quad (8)$$

where n is the cyclotron harmonic, J_n is a Bessel function of the first kind, and $\omega_c = eB/mc$ is the cyclotron frequency of nonrelativistic electrons. Note that the sign of the parallel electron velocity is not resolvable by this measurement, since the Doppler frequency shift is absent when observing perpendicularly. Compensating for this drawback is a fortuitous circumstance: electrons initially at the same energy all radiate at the same frequency, *both initially and subsequently*.

Consider excitations over a finite region of momentum space, where the perturbed electrons contribute to the incremental radiation by radiating, at least initially, through at most one harmonic, say the n th, into the range of detected frequencies. Suppose that we wish to deduce the distribution of electrons initially at momentum u_0 , i.e., to deduce the $Q_k(u_0)$. At time $\tau = \tau_0$, only the measurement of the frequency $\omega = n\omega_c/\gamma(u_0)$ constrains the $Q_k(u_0)$. In the absence of a dc electric field, superthermal electrons slow down in energy, but diffuse only in pitch angle. Hence, at some later time, say $\tau = \tau_1$, electrons initially at momentum u_0 will all be found on the same energy surface, say $u = u_1$, and radiating at the same frequency $\omega = n\omega_c/\gamma(u_1)$. Consequently, to deduce the $Q_k(u_0)$, all available information is found in the projection $R(\omega, \tau) \rightarrow r(\tau) \equiv R(\omega(\tau), \tau)$, where, $\omega(\tau) = n\omega_c/\gamma(\rho)$ tracks electrons slowing down on the trajectory $\rho(u_0, \tau)$.

Thus, rewriting $\delta(\omega - n\omega_c/\gamma(\rho)) = \delta(u - u(\omega, \tau))\rho\gamma(u)/(\omega u^2)$, we substitute Eqs. (4) and (8) into Eq. (5). We take the projection $\omega \rightarrow \omega(u, \tau)$ and integrate over the δ -function, which now selects only one u , which enters as a parameter. We obtain

$$r(\tau) \equiv R(\omega(\tau), \tau) = \frac{e^2 \omega}{c} \sum_k \gamma(u) Q_k(u) \left(\frac{\gamma(u) + 1}{u} \right)^{\alpha_k} \left(\frac{\rho}{\gamma(\rho) + 1} \right)^{\alpha_k} \frac{\rho^3}{\gamma^2(\rho)} H_{nk}(\rho), \quad (9)$$

where $\omega = \omega(u, \tau)$ and $\rho = \rho(u, \tau)$, as discussed, and

$$H_{nk}(\rho) = \int_{-1}^1 P_k(\mu) \mu^2 J_n^2 \left(n \frac{\rho}{\gamma(\rho)} (1 - \mu^2)^{1/2} \right) d\mu, \quad (10)$$

where P_k is the k th Legendre harmonic and $H_{nk} = 0$ for k odd.

Now Eq. (9) is a much easier inverse problem; it is 1-D with u entering as a parameter. Moreover, there is a trick available here to invert this equation analytically. We deduce the Q_k from the behavior of r as $\rho \rightarrow 0$, or, equivalently, $\tau \rightarrow G(u)$. The exploitable property of Eq. (9) is that the coefficients of the higher order Q_k have higher order zeroes. Since we are interested in the behavior particularly as $\rho \rightarrow 0$, we can Taylor expand the Bessel function to put Eq. (9) into the form

$$\Phi(\rho) = \sum_k \hat{Q}_k F_k(\rho), \quad (11)$$

where, we rewrote $\tau = G(u) - G(\rho)$, and we defined

$$\begin{aligned} \Phi(\rho) &= r(\tau) / \left[(e^2 \omega \rho / c) (\rho / \gamma(\rho))^{2n+2} \right], \\ \hat{Q}_k &= \gamma(u) \left((\gamma(u) + 1) / u \right)^{\alpha_k} Q_k(u), \\ F_k(\rho) &= \left(\frac{\rho}{\gamma(\rho) + 1} \right)^{\alpha_k} \left(a_{nk}^{(0)} - \frac{n^2}{2(n+1)} \left(\frac{\rho}{\gamma(\rho)} \right)^2 a_{nk}^{(1)} + \dots \right), \end{aligned} \quad (12)$$

where the $a_{nk}^{(i)}$ are constants that arise from the small argument expansion of the Bessel function, e.g.,

$$a_{nk}^{(0)} = (n^n / 2^n n!)^2 \int_{-1}^1 P_k(\mu) \mu^2 (1 - \mu^2)^n d\mu. \quad (13)$$

Note that as $\rho \rightarrow 0$, $F_k(\rho) \sim \rho^{\alpha_k}$. Now suppose that $\Phi(\rho)$ is analytic — i.e., assure this by replacing the experimental projection $r(\tau)$ by an appropriately smoothed fit. Then, if we differentiate Eq. (11) α_k times and evaluate at $\rho = 0$, terms higher than k vanish, and we find (for k even)

$$\hat{Q}_k = \frac{2^{\alpha_k}}{\alpha_k! a_{nk}^{(0)}} \frac{\partial^{\alpha_k}}{\partial \rho^{\alpha_k}} \left[\Phi(\rho) - \sum_{j=0}^{k-1} \hat{Q}_j F_j(\rho) \right]_{\rho=0}, \quad (14)$$

which solves for the $\hat{Q}_k$ (and hence for the Q_k) by recursion. In writing Eq. (14), we assumed that α_k is an integer — but, in fact, more general α_k may be handled similarly.

It is possible that the excitation $S(\mathbf{p})$ is known to be unidirectional, i.e., $Q(\mu < 0) = 0$. This auxiliary knowledge may be exploited to particular advantage. Separate $Q(\mu)$ into $Q_e + Q_o$ which are, respectively, its even and odd components. Now $Q(\mu < 0) = 0$ implies that $Q_e = Q_o$, and, hence, for $\mu > 0$, we have $Q = 2Q_e$. Thus, our solution for the even Legendre harmonics, which gives Q_e , represents, in fact, a complete and unambiguous inversion when unidirectional excitation can be assumed. Without this additional constraint, the null space of ψ consists of all odd harmonics; with the unidirectional assumption, the null space vanishes. The range of ψ , however, may be shown to be a notably small subset of R -space.

Although we have solved the inverse problem given emissions of the ordinary wave polarization, an entirely analogous procedure may be used to deduce $Q(\mathbf{p})$ from emissions of the extraordinary wave polarization. This gives rise to several interesting questions, including the puzzle of how to approach the formally overspecified inverse problem when, happily, emission data of both polarizations are available.

In summary, we have proposed an invasive method for multiplying by orders of magnitude the information obtainable from synchrotron radiation. The key has been to use knowledge of the dynamics of fast electrons. These are the electrons that emit the most radiation and about which we are generally most curious. The inverse problem for an interesting special case has been solved exactly. A more general case may be approached numerically, and further generalization is available by considering more complicated plasma dynamics. Perturbative solutions about our analytical result, say for small dc electric field or observation slightly off from $\theta = 0$, might be useful extensions. If the dynamics of fast electrons is not known precisely, then the methods here might be employed to pick between competing models of their dynamics. We speculate, too, that on the basis of the very detailed information obtained here, inferences might be drawn concerning macroscopic plasma parameters, such as the current profile, which are difficult to deduce through other means.

It is a pleasure to acknowledge stimulating discussions with P. Efthimion, T. Luce, and H. Segur. This work, initiated under a grant from the John Simon Guggenheim Memorial Foundation, was supported by United States Department of Energy under contract number DE-AC02-76-CHO3073.

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