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**FREQUENCY-SWEEPING: A NEW TECHNIQUE FOR
ENERGY-SELECTIVE TRANSPORT**

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Frequency-Sweeping: A New Technique for Energy-Selective Transport

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Abstract

A new method is described for inducing energy-selective transport by 'sweeping' the frequency of applied low- n magnetic perturbations. The mechanism, formally analogous to the 'rising buckets' concept in accelerator physics, can move particles with a selected velocity in a nondiffusive fashion from one specified radius to another. The technique is considered principally as a means for removal of Helium ash. Other likely applications are as a method for burn control, profile control, as a diagnostic, and perhaps as a nonstochastic means of effecting the direct coupling of alpha power recently discussed by Fisch and Rax.

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In this work, we introduce a new technique for inducing energy-selective transport in a tokamak, achieved by ‘sweeping’ the frequency of applied low- n magnetic perturbations. The mechanism, formally analogous to the ‘rising buckets’ concept[1] for accelerating particles, can nondiffusively move particles with a selected velocity from one specified radius to another, a distance comparable to the tokamak minor radius $r = a$, in on the order of a hundred toroidal transit times $\tau_\zeta \simeq v_{||}/(qR)$. Here we focus on the use of the technique as a means of ash removal, but the same method could be employed for burn control, profile control, as a diagnostic, and perhaps as a nonstochastic means of effecting the direct coupling of alpha power recently proposed by Fisch and Rax.[2]

Earlier, we have proposed[3, 4] another possible method for ash removal, based on the energy-selective properties of a *stochastic* transport mechanism. The technique considered here shares with the earlier stochastic method the energy-selective property, and the virtue of being energetically inexpensive, and may be optimized by making use of the energy scalings underlying the earlier mechanism. It has the additional strength of requiring substantially smaller applied perturbation amplitudes, and in addition probably requires less careful tailoring than the stochastic mechanism of the spatial structure of the needed perturbation.

The basic idea of the mechanism is simple. For fixed frequency ω , a single-harmonic perturbation $\tilde{b} \equiv \delta B_r/B_0 = b_{mn}(r) \sin(n\zeta - m\theta - \omega t)$ induces an island in the particle drift surfaces having the same (m, n) symmetry (plus sideband islands,[3] subdominant for current considerations), whose resonant surface r_{res} depends upon ω , as given by the resonance condition $\omega = n\Omega_\zeta - m\Omega_b$. (Here, Ω_b is the poloidal transit/bounce frequency and $\Omega_\zeta \equiv 2\pi/\tau_\zeta$ is the toroidal transit/drift frequency.) For passing particles, one has $\Omega_\zeta \simeq \bar{v}_{||}/R, \Omega_b \simeq \Omega_\zeta/q(r)$ ($\bar{v}_{||}$ is the transit-averaged parallel

velocity), and thus

$$q_{res}(\omega) \equiv q[r_{res}(\omega)] \simeq q_{mn}/[1 - (\omega/n\Omega_\zeta)], \quad (1)$$

where $q_{mn} \equiv m/n$. This is illustrated in Figure 1, showing three Poincare plots of the drift surfaces for particles at one tenth the alpha birth energy $E_0 = 3.5$ MeV and with pitch $\lambda \equiv v_{||}/v = 1$, using TFTR parameters [$R = 2.62$ m, $a = .96$ m, $B = 5$ T, and taking $q_0 \equiv q(r = 0) = 1, q_a \equiv q(a) = 4$]. For each, the same $(m, n, b_{mn}^{max}) = (2, 1, .8 \times 10^{-3})$ perturbation is used, but with the frequency assuming three values. One notes that the primary (2,1) island moves outward as the frequency changes, from near $r = 0$ (a) to near the machine edge (c).

If ω is varied slowly enough, a particle starting inside an island (or ‘bucket’) at $r_{res}(\omega)$ will move adiabatically with that island, at radial velocity $\dot{r}_{res} = \dot{q}_{res}/q'_{res}$, with $\dot{q}_{res} = q_{res}\dot{\omega}/(n\Omega_\zeta - \omega)$, where $\dot{\omega} \equiv d\omega/dt$ is the rate at which the frequency is swept. In this way, particles at a specified phase point $(r, \mathbf{v})_1$ can be picked up in a bucket at frequency ω_1 , and then walked nonstochastically to another radius r_2 , the bucket radius for frequency ω_2 . Particles initially outside the island remain outside the island, receiving only a small radial kick as the island sweeps over them.

As will be shown, the limit on the rapidity $\dot{\omega}$ of the frequency sweep is that the nonlinear trapping frequency ω_{tr} of a particle in the island be greater than the characteristic frequency for the island moving radially an island halfwidth δr_d . Including the numerical factors found later,

$$\omega_{tr} \gtrsim 2\dot{r}_{res}/\delta r_d. \quad (2)$$

The characteristics just summarized can be used to achieve the energy-selectivity needed for the various applications of the technique mentioned earlier. For ash removal, one wants to greatly enhance the transport of the alphas at an energy E_1 small compared with E_0 , while not greatly enhancing the transport of both energetic alphas, with $E \sim E_0$, or of particles in the

thermal bulk $E \sim T$. Since one has $E_0/T \sim 10^2$, we choose the target ash energy E_1 equal to $E_0/10$. Thus, particles at E_1 have already given most of their original energy to the plasma, but may be still be readily distinguished in their transport characteristics from thermal ions. In the earlier stochastic mechanism, energy selectivity was provided by orbit-averaging effects. Frequency-sweeping provides three additional parameters to achieve this, *viz.*, ω_1, ω_2 , and $\dot{\omega}$. As indicated above, for a typical particle of energy E_1 , with $\Omega_\zeta = \Omega_{\zeta 1}$, one can set ω_1 so that the induced bucket lies deep inside the plasma [$r_1(\Omega_{\zeta 1}) \simeq 0$], and ω_2 so that the bucket has moved the target particle nearly or fully outside the machine [$r_2(\Omega_{\zeta 1}) \simeq a$]. From Eq.(1), for the (2,1) perturbation, particles with $E \sim E_0$, having a much larger typical Ω_ζ than those with $E \sim E_1$, will have an r_1 much further out (near the $q = 2$ surface), and an r_2 much further in. Thus, many fewer birth alphas will be swept, both because the bucket picks up particles where not many birth alphas exist, and because the distance $r_2 - r_1$ swept is much smaller. Similarly, thermal particles, having Ω_ζ small compared with those with $E = E_1$ will have a $q_1 \equiv q(r_1) < q_0$, *i.e.*, the bucket picks up no thermals, because the initial bucket position does not exist within the plasma. Moreover, one can choose $\dot{\omega}$ so that criterion (2) holds for particles with $E \gtrsim E_1$, but is violated for those with $E \sim T$, further diminishing the effect of the swept buckets. Finally, further energy selectivity can be provided, if needed, by making use of the orbit-averaging mechanism on which the stochastic scheme is based.

The energy selectivity is illustrated by the orbits shown in Figure 2, for $E/E_0 =$ (a).01, (b).1, and (c)1.0, and for each particle having pitch $\lambda = .8$. The same (2,1) perturbation is used as that in Figure 1, but with frequency swept from $\omega/\Omega_g = -5 \times 10^{-3}$ to 2.6×10^{-3} in a time $200\tau_{\zeta 1}$, where $\tau_{\zeta 1} \equiv \tau_\zeta(E = E_1, \lambda = 1)$. (Ω_g is the alpha gyrofrequency, equal to $2.38 \times 10^8 s^{-1}$ for $B = 5T$.)

•Theory

We analyse the action of frequency sweeping analytically. The radial motion of the particle due to the perturbation is given by considering $\bar{r}$, the transit/bounce-averaged minor radius, which evolves according to [3, 4]

$$\dot{\bar{r}} = \dot{\bar{r}}_A + \dot{\bar{r}}_B = \sum_{l_b, n} v_{l_b} \sin \eta_{l_b}, \quad (3)$$

where $\eta_{l_b} \equiv n\bar{\zeta} + l_b\theta_b + \phi_b$, $\bar{\zeta}$ is the transit/bounce-averaged toroidal azimuth ζ , θ_b is the bounce phase, and $l_b = l - \sigma m$ is the bounce harmonic, with $l = 0, \pm 1, \pm 2, \dots$, and $l = 0$ corresponding to the primary island, and with trapping-state index σ , equal to 1(0) for passing (trapped) particles. ϕ_b is the perturbation phase, with $\dot{\phi}_b = -\omega$. In the absence of a perturbation, one has $\dot{\theta}_b = \Omega_b$ and $\dot{\bar{\zeta}} = \Omega_\zeta$, so that

$$\dot{\eta}_{l_b} \simeq n\Omega_\zeta + l_b\Omega_b - \omega. \quad (4)$$

The location of the l_b^{th} resonant surface is thus given by $\dot{\eta}_{l_b} = 0$. As discussed in detail in Refs. 3 and 4, $\dot{\bar{r}}_A \sim \tilde{b}$ is the portion of $\dot{\bar{r}}$ due to the perturbing radial field δB_r , which induces ‘magnetic braiding,’ while $\dot{\bar{r}}_B \sim \delta B/B_0$ is the contribution due to ‘grad- B ’ drifts from ripple perturbations δB in the magnitude B of the field. For each l_b and n , the amplitudes v_{l_b} contain both the braiding and ripple contributions, and a sum over the contributions from all magnetic harmonics m having that n . For simplicity, here we drop the ripple contribution, take a single n , and consider a single island l_b . Transforming $\bar{r}$ to the frame $\bar{r}' \equiv \bar{r} - r_{res}$ moving with the swept island and expanding $\dot{\eta}_{l_b}$ in r about the resonant surface, one has

$$\dot{\bar{r}}' \simeq v_{l_b} \sin \eta_{l_b} - \dot{r}_{res}, \dot{\eta}_{l_b} \simeq \bar{r}'/M_{l_b}, \quad (5)$$

formally the equations for one dimensional motion of a particle with mass (for passing particles) $M_{l_b} \equiv -q/(l_b\Omega_b q')$, with canonically conjugate coordinate η_{l_b} and momentum $\bar{r}'$, moving according to Hamilton’s equations $\dot{\eta}_{l_b} =$

$\partial_{\tilde{r}'} H$ and $\dot{\tilde{r}'} = -\partial_{\eta_{l_b}} H$, with Hamiltonian $H(\eta_{l_b}, \tilde{r}') = \tilde{r}'^2/(2M_{l_b}) + V(\eta_{l_b})$, and with potential $V(\eta_{l_b}) \equiv v_{l_b} \cos \eta_{l_b} + \eta_{l_b} \dot{r}_{res}$. For nonzero $\dot{r}_{res} \propto \dot{\omega}$, the sinusoidal wells in V lie on a slope. Thus, as well known from the formally similar problem of particle motion in stellarators or rippled tokamaks,[5] the separatrices (buckets) in phase space have single x-points, with orbits outside the separatrices passing through the $\tau' = 0$ resonant surface, and a reduced well-depth $v_{ef} \equiv [V(\eta_{l_b}^{max}) - V(\eta_{l_b}^{min})]/2 \simeq v_{l_b} \{[1 - (\dot{r}_{res}/v_{l_b})^2]^{1/2} - (\dot{r}_{res}/v_{l_b})[\pi/2 - \arcsin(\dot{r}_{res}/v_{l_b})]\}$. The island halfwidth and trapping frequency are given by $\delta r_d = (4v_{ef}M_{-m})^{1/2}$, and $\omega_{tr} \simeq (v_{ef}/M_{-m})^{1/2}$. In the limit $\dot{r}_{res}/v_{l_b} \rightarrow 1$, $v_{ef} \rightarrow 0$, and the separatrix width vanishes, yielding criterion (2) limiting the size of $\dot{\omega}$.

•Multiple Particles, Multiple Sweeps

A single island sweep will not capture all the particles in the target distribution, both because the initial particle phase will place particles outside the bucket, and because the ash distribution exists over a range of r and λ . Since a typical sweep time $\tau_{sw} \simeq 100\tau_{\zeta 1} \simeq 200\mu\text{sec}$ is very small compared to a typical slowing down time $\tau_{sl} \sim 200 \text{ msec}$, many sweeps may be used, with frequencies $\omega_{1,2}$ tuned to access $E \sim E_1$ particles at different values of (r, λ) . In Figure 3a we show the effect of a series of identical sweeps, each of duration $200\tau_{\zeta 1}$ for an ensemble of co-going particles, with initial radii in the range $r/R = .01$ to $.10$, and pitch in the range $\lambda = 0.6$ to 0.9 . The sweeps are from $\omega_1/\Omega_g = -4 \times 10^{-3}$ to $\omega_2 = -\omega_1/2$. We plot the ensemble-averaged radius $\langle r \rangle$ of the particles versus time, in units of the number of sweeps. One notes, as in Figure 2, that the $E = E_1$ particles are selectively swept. One also notes a diminution with sweep number of the rate of change of $\langle r \rangle$, as the region of phase space swept by that particular (ω_1, ω_2) bucket is cleared out, indicating time to move to a new (ω_1, ω_2) -setting.

The earlier stochastic mechanism required that the perturbation strength $\tilde{b}$ exceed an overlap threshold in order to be operative. For the (2,1) pertur-

bation, a threshold of $b_{mn} \simeq 4 \times 10^{-3}$ was found,[3] a very large perturbation. In contrast, the present mechanism is operative at arbitrarily low b_{mn} , thus avoiding potential deleterious effects on other aspects of the machine operation. The Figures shown in this work have all used $b_{mn}^{max} = .8 \times 10^{-3}$, only a fifth of the stochastic threshold. This is still a rather large perturbation, used to provide easy visibility of the sweeping. In a real implementation $\tilde{b}$ could readily be taken far below this, limited only by the need to compensate for the smaller bucket width $\delta r_d \propto v_{cf}^{1/2} \propto \tilde{b}^{1/2}$ by an increased number of sweeps to access the full target distribution.

•Other Considerations

A virtue of the earlier stochastic mechanism was that, because it can operate at zero or low frequencies, it does little work in moving the ash to the outside. The frequency $\omega/\Omega_g = -6 \times 10^{-3}$ used in Figure 1a to place the bucket near $r = 0$ translates to $\omega/(2\pi) = -230$ kHz, no longer a low frequency. As a result, there can be significant energy exchange $\dot{E}$ between the particles and the perturbation. However, as can be understood by averaging the Hamilton's equation $\dot{E} = -\partial_t h$ over a trapping period $\tau_{tr} \equiv 2\pi/\omega_{tr}$, the sign of $\langle \dot{E} \rangle_{\tau_{tr}}$ reverses when the sign of ω does, positive for $\omega < 0$, and negative for $\omega > 0$ for co-going particles. Because a sweep with a (2,1) perturbation requires both $\omega < 0$ and $\omega > 0$, the net energy change over a sweep is again small. One sees this in Figure 3b, showing the ensemble-average of the energy at each sweep versus time, for the same ensemble as in Figure 3a. It is thus possible to make the ash-removal mechanism be energetically nearly neutral, by such a choice of q_{mn} . However, we note that it is also possible to cause *negative* net work to be done on the particles by choosing q_{mn} near q_0 . Tuned to resonate with alphas with $E \sim E_0$ rather than E_1 , such a perturbation could draw net energy from the newborn alphas, perhaps permitting a nonstochastic implementation of the scheme of Fisch and Rax[2] or an approach for direct conversion.

However, q_{mn} is also constrained from below. One notes from Eq.(1) that reversing Ω_ζ or $\bar{v}_\parallel$ requires reversing $\omega_{1,2}$ to sweep particles across the same $r_1 \rightarrow r_2$ interval. If $|\omega_1/\omega_2| \leq 1$, therefore, sweeps tuned to sweep co-going particles from $r = 0$ outward will also sweep counter-going particles from inside $r = a$ inward, tending to cancel the effect of sweeps tuned to move counter-going particles out. From Eq.(1), one finds $\omega_1/\omega_2 = (q_a/q_0)(q_0 - q_{mn})/(q_a - q_{mn})$, negative and with magnitude greater than unity only for $q_{mn} > 2q_0q_a/(q_0 + q_a)$. For the q_0, q_a used here, this requires $q_{mn} > 8/5$, satisfied by the (2,1) perturbation, for which $\omega_1/\omega_2 = -2$.

The energy dependence in Eq.(1) arises through $\bar{v}_\parallel$ or Ω_ζ . While this equation is only approximate, neglecting the additional E -dependence in the particle drifts, it raises the difficulty that particles with $E \sim E_0$, but with small λ (*viz.*, trapped or nearly trapped particles) can have values of Ω_ζ the same as that typical of the $E \sim E_1$ ash, and so might be swept as well. To suppress the transport of such energetic particles, as noted earlier, one can employ the further energy selectivity arising from orbit-averaging.[3, 4] For example, for a single- m braiding perturbation, one has $v_{l_b} \simeq \bar{v}_\parallel b_{mn} J_l(\eta_1)$, with perturbing amplitude b_{mn} , orbit-avging parameter $\eta_1 \simeq k_\perp \rho_b$, and $\rho_b \propto E^{1/2}$ the particle drift-orbit excursion. As shown in Refs. 3, 4, taking $n \simeq 3$ is large enough to put J_0 at its first zero for particles with $(E, \lambda) = (E_0, 1)$, thereby causing the primary bucket width to vanish. Energetic trapped particles have still larger η_1 , and are therefore even more susceptible to orbit-averaging effects.

Summarizing, we have proposed a new technique for selectively transporting particles from one minor radius to another, and from a limited region of velocity space. A number of potential applications of this technique needed for a tokamak reactor have been noted, perhaps the most demanding of which is ash removal, the principal focus of this work. A theory of the frequency-sweep mechanism has been developed, most fully for pass-

ing particles, numerical evidence has been presented supporting the efficacy of the proposed ash removal scheme, and some of the constraints and potential difficulties of the technique have been indicated. The precise type of perturbations needed to implement the scheme and the means by which the perturbations can be induced have not been specified here. However, the various constraints considered do serve to narrow the field of possibility. Required are low- n magnetic perturbations, with frequencies up to a few hundred kHz, so they can resonate with energetic alphas. An antenna system with a tunable frequency, similar to the saddle-coil array proposed[6] to excite TAE and EAE modes in JET seems appropriate to test the present scheme, particularly as an acceptable test could be done with a perturbation strength as small ($\tilde{b} \sim 10^{-5}$) as that for which that system is designed.

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Figures

Fig. 1. Poincare plots of drift surfaces of particles with $E = E_0/10, \lambda = 1$ in TFTR, in the presence of the same (2,1) perturbation, but where the frequency assumes three values $\omega/\Omega_g =$ (a) -6×10^{-3} , (b) 0, and (c) 2×10^{-3} .

Fig. 2. Poloidal projections of passing orbits with $\lambda = .8$, and with energies $E/E_0 =$ (a).01, (b).1, and (c) 1.0, moving under the influence of a single sweep of a (2,1) perturbation.

Fig. 3. (a) Plot of ensemble-averaged radius of particles versus time, in units of the number of sweeps applied, for an ensemble of co-going particles, with energies $E/E_0 = .01, .1$, and 1.0. (b) Ensemble-averaged energy versus time, for the same simulation as in (a).

Fig. 1

Fig. 2

Fig. 3

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