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ROBERT A. SMITH

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MODELS OF JOVIAN DECAMETRIC RADIATION*

By

Robert A. Smith† **

Theoretical Studies Group
NASA Goddard Space Flight Center
Greenbelt, Maryland 20771

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†NAS-NRC Postdoctoral Resident Research Associate.

**Present address: Observatoire de Paris, 92190 Meudon, France

ABSTRACT

We present a critical review of theoretical models of Jovian decametric radiation, with particular emphasis on the Io-modulated emission. The problem is divided into three broad aspects: the mechanism coupling Io's orbital motion to the inner exosphere, the consequent instability mechanism by which electromagnetic waves are amplified, and the subsequent propagation of the waves in the source region and the Jovian plasmasphere. At present there exists no comprehensive theory that treats all of these aspects quantitatively within a single framework. Acceleration of particles by plasma sheaths near Io appears to be a promising explanation for the coupling mechanism, while most of the properties of the emission may be explained in the context of cyclotron instability of a highly anisotropic distribution of streaming particles. The present state of the theory is evaluated, and some suggested approaches for future work are discussed.

I. INTRODUCTION

The nature and origin of the Jovian decametric radiation (DAM) have prompted much theoretical effort, especially since the remarkable modulating effect of Io was demonstrated by Bigg [1964]. The problem occupies an unusual niche in the gallery of radio-astronomical phenomena. As is true of emission from quasars, pulsars, and the sun, DAM presents a rich phenomenology. Unlike the cases of quasars and the sun, however, which exhibit detailed line spectra at optical wavelengths, radio observations [of both DAM and the weaker decimetric emission (DIM)] provided the only information about the plasma environment of Jupiter before the recent probes of the Jovian magnetosphere by Pioneers 10 and 11. Thus, for nearly twenty years the physical characteristics of the Jovian magnetosphere were inferred solely from radio astronomy. In this respect the problem is similar to that of pulsars. Like the sun, however, Jupiter is close enough that a detailed observational morphology of its radio emissions may be constructed. This unusual combination of circumstances, together with the extensive development of understanding of terrestrial magnetospheric processes, has allowed for wide variability in the mixture of inductive and deductive arguments in various individual attempts to explain DAM.

Despite the great amount of theoretical effort that has so far been expended, we have not yet arrived at a comprehensive, convincing understanding of DAM. Nevertheless, a review of the theoretical work to date is both pertinent and timely. The results from the Pioneer 10 and 11 missions have excited widespread interest in Jovian physics in general, and much of the conceptual framework in which we evaluate these results has been developed in the course of attempts to elucidate DAM. We anticipate that the relations between the problems of understanding DAM in particular and the Jovian magnetosphere in general will be more widely appreciated in the future. Furthermore, although we do not view any of the attempts heretofore to explain various facets of DAM as wholly successful, we believe that in the paradigmatic nature of science a successful theory will emerge from the gradual modification and consolidation of previous work, rather than from radically new hypotl ses. In this spirit, we attempt in this review to present a critical framework for the evaluation of the theoretical efforts to date, in the hope that such a framework will prove useful in both motivating and evaluating future work. The bulk of our discussion concerns the Io-modulated DAM; we discuss the Io-independent emission briefly near the end.

A current review of DAM observations appears in this volume [Carr and Desch, 1975]. Previous reviews by Warwick [1964, 1967, 1970] and by Carr and Gulkis [1969] deal with observations of both DAM and DIM and the inferences about the Jovian magnetosphere drawn from these observations. In addition, Warwick [1967] presents a critical review of some of the theoretical work up to 1967.

Our approach in this review is to separate the theoretical problem into various more-or-less distinct facets, and to consider these facets both in themselves and in their structural relationships to each other. Most of the theoretical works to date deal principally with only one or two of the different elements of the problem, and so we discuss each work in the most appropriate contexts. We assume that the reader is familiar with the observations as they are described in the above-mentioned reviews, and so our references to observational results will usually be superficial and will seldom be referenced; the article by Carr and Desch in this volume will generally provide adequate observational background for our discussion.

Some preliminary semantic definitions may be in order. As we have already done above, we shall sometimes distinguish between the phenomenology and morphology of DAM. By "phenomenology" we mean those aspects that are directly observed, such as intensity, time scales, bandwidth, polarization, and such features of the dynamic spectra as frequency drifts and modulation lanes. In contrast, we mean by "morphology" the organization of the data into such features as rotational profiles of occurrence probability and polarization, the similarity of the dynamic spectra of different events observed under the same geometrical conditions, the declination effect, the induction by Bigg of Io-modulation, the concomitant recognition of an

Io-independent component, and so forth. Moreover, we do not distinguish between the familiar decametric emissions and the hectometric emissions observed from satellites by Brown [1974a, b] and by Desch and Carr [1974]. Except for the component of hectometric emissions that was designated by Brown [1974b] as "mid-frequency" (MF) emission, the hectometric emission appears to be merely a continuation of the decametric spectrum to longer wavelengths, and we shall include the hectometric range in the traditional designation DAM. We note, however, that at this writing (mid-1975) the observational statistics are insufficient to determine whether either the "normal" or the MF components of hectometric emission are modulated by any of the Jovian satellites.

Before plunging into the literature, we discuss some theoretical preliminaries.

The high intensities, limited bandwidths, short time scales, and sporadic nature of individual DAM bursts indicate that they originate in stimulated emission from one or more collective micro-instabilities. The remarkable repeatability of the dynamic spectra observed under similar geometrical conditions, and the complex beaming pattern manifested in the source morphology, imply that the structure of the ambient medium is stable in the long-term sense and determines both the local conditions for plasma instability and the subsequent propagation of the radiation. Although this statement seems indisputable in the general sense, it must nevertheless be taken with some qualification. From detailed study of certain repeatable features of dynamic spectra, Warwick [1963] showed that these features always appeared for particular corresponding central meridian longitudes (CML), within a range of $\pm 9°$. Warwick interpreted these results as implying that the emission was beamed into a cone of $9°$. Alexander [1975], however, compared several pairs of identical-looking dynamic spectra observed one Jovian year (≈ 11.9 y) apart. Although the members of each pair of spectra commonly shared quite detailed features, Alexander concluded from a study of the rotation rates inferred from the different pairs that the beaming pattern might be subject to longitude shifts over a range of up to $\pm 10°$. Another point is so obvious that it is not often made explicitly, but we feel that perhaps for that very reason its implications are sometimes lost sight of in the literature. This point is that the occurrence probability distribution in the (CML, γ_{Io}) plane [γ_{Io} is the phase of Io from superior geocentric conjunction (SGC)] is conventionally interpreted as indicative of a beaming pattern. The occurrence probability over most of the plane, however, is substantially less than unity. This fact alone indicates that there must be some variable influence on the modulation mechanism or the waveguide structure of the plasma medium, or both.

Plasma instabilities require a source of free energy. Large-amplitude magnetohydrodynamic (MHD) waves or shocks may provide free energy under some circumstances. In micro-instabilities (by which we mean instabilities that are predicted only in kinetic-theoretical analyses and cannot be predicted from fluid theory) the free energy resides in some nonthermal feature of the particle distribution function $F(\vec{r}, \vec{v}, t)$. Such a nonthermal feature may be manifested in either the configuration-space or velocity-space dependence of F. For example, density gradients in a magnetized plasma create drift currents, which may drive a variety of instabilities. Velocity-space features such as beams of high-velocity particles, gaps in the distribution in some velocity range, or anisotropy in either pitch angles or Larmor phases, may result in instability.

The repeatability of DAM spectra suggests that the instability occurs for waves of frequency at or related to one of the characteristic frequencies of the plasma, as is commonly the case. In the high-frequency regime, these are the electron plasma frequency f_{pe}, the electron gyrofrequency f_{ce}, and the upper hybrid frequency $f_{UH} = \left(f_{pe}^2 + f_{ce}^2\right)^{1/2}$. At Jupiter, the DAM frequency range may be identified with these characteristic frequencies only near the planet: roughly speaking, at altitudes of $\leqslant 1\ R_J$.

Thus, the central problem in the theory of DAM is to elucidate the <u>coupling mechanism</u> whereby the orbital motion of Io provides a free-energy source in the inner plasmasphere. The existence and nature of an instability are determined by the form in which the free energy appears. The coupling mechanism and the subsequent instability mechanism by which the radio emission is amplified are, however, viewed here as two logically distinct elements of the problem. Roughly speaking, we might expect that the coupling mechanism primarily affects the DAM morphology, while the nature of the instability is reflected in the phenomenology. This statement is intended simply as a general guide, however, and not as an organizing postulate.

The third element of the problem is the propagation of the radiation in the Jovian plasmasphere. Propagation effects are expected to be manifested in both the phenomenology and morphology. As we shall see below, for example, certain instability mechanisms that have been proposed actually amplify waves that are trapped in the Jovian plasmasphere. Therefore, some mode-conversion process is necessary in order for the radiation to propagate into free space. In this case the polarization of the escaping wave, which is generally the datum by which observers identify the base mode, may be determined by the escape mechanism rather than by the instability. The escape mechanism may also determine the initial wave-vector of the escaping ray, thus determining the ray path in the plasma medium and the beaming pattern.

Provided the boundaries between the different aspects of the problem are not regarded as impenetrable, our formulation is a useful one which enables us to construct a general conceptual framework within which to consider particular theories.

II. THE COUPLING MECHANISM

The question of the coupling mechanism has been addressed, to varying degrees, by many authors. Smith and Wu [1974] divided the majority of the proposed mechanisms into three broad categories:

(i) Wave models, in which the motion of a conducting Io through the magnetoplasma may generate whistlers or MHD waves that propagate to the Jovian ionosphere;

(ii) Acceleration models, in which the electric field induced in the frame of Io by the satellite's motion through the ambient magnetic field is invoked to accelerate particles to energies from tens to hundreds of keV. These particles then precipitate along the Io-threaded flux tube (IFT);

(iii) The sweeping model, in which energetic particles trapped on the IFT are removed from the trapping volume by impingement on Io.

To our knowledge, the above categories embrace all the cases in the literature in which a coupling mechanism has been examined quantitatively. This does not mean that their corresponding instability mechanisms have been considered quantitatively; as we shall see in section III, no quantitative model of an instability mechanism that might result from MHD-wave coupling has yet been presented in the literature. Whether or not such mechanisms may exist, the generation of MHD-wave disturbances by Io may have important consequences which are discussed below.

In addition to the coupling models described above, other more speculative hypotheses have been advanced. Duncan [1970] suggested that the source morphology was determined by a combination of the geocentric beaming pattern and heliocentric effects. He constructed a model of the Jovian magnetosphere in which the subsolar point of the magnetopause was at 5 R_J, and conjectured that the strong confinement of Io modulation to orbital phases near $\gamma_{Io} \simeq 90°$ SGC and 240° SGC implied that Io crossed the magnetopause at about these phases. Duncan further speculated that these magnetopause crossings would precipitate trapped particles into the Jovian ionosphere, although he did not indicate why or how this might occur. To evaluate Duncan's hypothesis, however, it suffices to note that the Pioneer 10 and 11 probes have revealed the Jovian magnetosphere to be extended to characteristic dimensions of 50-100 R_J in the dayside hemisphere.

Another phenomenological suggestion in a spirit similar to that of Duncan's was made by Conseil et al. [1971, 1972]. These authors inferred a correlation between the velocity of the solar wind at Jupiter (extrapolated from satellite measurements at 1 AU) and the phase of Io at the onset of DAM storms. They suggested that the correlation might be attributed to a plasmapause bulge, similar to that of the terrestrial plasmapause, the location of which varied in response to solar wind ram pressure. They speculated that Io might trigger DAM storms upon passing through the plasmapause boundary.

The statistical significance of the correlation asserted by Conceil *et al.* is unclear; it has been questioned by Carr and Desch [1975]. Furthermore, a Jovian plasmapause would be unlikely to contain a bulge analogous to that in the terrestrial case. It is believed that the terrestrial plasmapause occurs at the boundary between closed and open streamlines arising from the combined flow fields of corotation and convection. In the Jovian magnetosphere, corotation would be expected to dominate convection out to large distances, because of the rapid rotation rate of Jupiter. A comparison of the flow pattern to be expected at Jupiter and that in the terrestrial magnetosphere is given by Brice and Ioannidis [1970]. Although this picture may be subject to considerable modification in light of the unexpectedly large extension of the Jovian magnetosphere and the dominance of its outer structure by plasma outflow, it should not be greatly changed in the inner region. A true plasmapause at Jupiter would probably be determined by the breakdown of corotation.

Frank *et al.* [1975], however, inferred the existence of a plasmapause at the orbit of Io, where the density of protons of energies 100 eV - 4.80 keV was observed to increase by a factor of about 3 to 4 on the inbound pass of Pioneer 10. As may be seen in Figure 1 [Frank *et al.*, 1975], however, this density decreases again inside Io's orbit, before resuming a general increasing trend as the radial distance decreases. These observations could also be explained by injection of these particles by Io, rather than as a plasmapause structure.

Whether or not the structure of the thermal-plasma density near Io is indicative of a plasmapause, it is clearly complicated. The basic idea behind the suggestion of Conseil et al. was that the stimulation of DAM by Io might depend on the local plasma density. This density may, in fact, be an important parameter in the MHD-wave model of Goertz and Deift [1973; c.f. also Deift and Goertz, 1973; Goertz, 1973a, b] and the recent acceleration models of Shawhan et al. [1974] and of Smith and Goertz [1975].

We now consider the quantitative categories (i) - (iii). As was mentioned above, in most of the works discussed in this review the major focus is on only one aspect of the total problem; in the majority of the papers we consider in this section it is the coupling mechanism. In most of these papers, however, the authors had some hypothesis in mind concerning the subsequent generation or propagation of the electromagnetic waves. In some other works, such as that of Wu [1973], two or more aspects of the problem are treated integrally. Therefore, in order not to fragment our exposition unduly, we shall occasionally anticipate some points of the later discussion in order to place the immediate topic in context.

A. Wave Models

Ellis [1965] first speculated that the coupling might be due to either MHD or whistler (which he called "electromagnetic") waves generated by Io. He did not investigate how such waves might be excited, except to note that the relative velocity of Io normal to the corotating magnetic field was greater than the phase velocities of these modes, all of which have resonances (wave number $k \to \infty$) at angles $\theta < \pi/2$ with respect to the field. Ellis, therefore, assumed that these low-frequency modes might be generated by Cerenkov radiation from Io. This is by no means obvious, however, because in order to be a source for the waves, Io must either have an appropriate current distribution on its surface or cause such a current distribution in the ambient plasma by perturbing its flow.

Ellis hypothesized that the Io-generated waves would interact with pre-existing particle streams or "bunches" to stimulate the electromagnetic radiation observed as DAM, in analogy to the well-known phenomenon of stimulated VLF emissions in the terrestrial magnetosphere. Thus, the whistler waves are the important ones in this context; this mechanism was also considered by Chang [1963], before Io-modulation had been recognized. It is implicit in this scheme that the resulting DAM waves travel along field lines, and that the source region would therefore be the CML.

McCulloch [1967] developed Ellis's ideas for the coupling more quantitatively. He too assumed whistlers to be generated by Io, and did ray tracing in a model magnetosphere. By varying the parameters of the model, and by assuming that the coupling region where the whistlers interacted with particle streams was at magnetic latitudes between 75° and 80°, McCulloch obtained a distribution of coupling regions in the (CML, γ_{Io}) plane, for which the relative intensities of the propagated whistlers corresponded well with the distribution of occurrence probability of DAM (c.f. in particular his Figures 7, 9, and 10). In particular, McCulloch's results indicated that the time required for the propagation of the whistlers from Io to the Jovian ionosphere was of the order of 33 min. This is the time required for Io to move 15° in orbital phase relative to the corotating medium; in this way, McCulloch accounted for the asymmetry of the Io phase, for peak occurrence probability, about the Earth-Jupiter line (Figure 2).

The parameters which McCulloch found necessary to describe his model magnetosphere, in order to obtain good agreement between the coupling mechanism and observational morphology, are not in accord with the Pioneer observations. For example, McCulloch used an ionospheric electron temperature T_e of 2000 K, whereas Kliore et al. [1975] have inferred $T_e \simeq 750$ K. Furthermore, McCulloch states

[1967, p. 1302] that he uses a density model for the thermal plasma of the form derived by Melrose [1967]. But he also states that "(the) most satisfactory models, however, were those for which the electrons were concentrated in a disc about the magnetic equator, and the plasma frequency and gyrofrequency at Io were nearly equal" [McCulloch, 1967, p. 1303]. The latter condition, $f_{ce} \simeq f_{pe}$, is actually in accord with observations: for a dipole moment of 4 $G\text{-}R_J^3$ [Smith et al., 1975], the field strength $B_o \simeq 0.02$ G at Io and $f_{ce} = 56$ kHz, while for a plasma density at Io of $N_o \simeq 30$ cm^{-3} [Frank et al., 1975], $f_{pe} \simeq 55$ kHz. The feature of a plasma disk in the equatorial plane, however, is not contained in Melrose's density model, in which the density decreases monotonically with latitude along a field line.* Therefore, the possibility exists that McCulloch's work may be internally inconsistent. Finally, McCulloch's ray tracing calculations assumed a centered, tilted dipole field of equatorial field strength 15 G; it is not clear in context (c.f. p. 1306) whether *ad hoc* anomolies, such as were postulated by Ellis and McCulloch [1963] were incorporated or not. As has been shown by Acuña and Ness [1975] and Smith *et al.* [1975], however, the field near the surface is strongly influenced by higher-order moments.

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*The plasma disk *is* contained in density models developed by Gledhill [1967] and Goertz [1973a, 1975].

Apart from questions about the associated instability mechanism and consequent propagation effects, these considerations indicate that, at the very least, coupling by whistlers must be re-examined in the context of the plasma environment deduced from the Pioneer observations. This comment does not apply uniquely to the work of McCulloch, however.

Other authors have concentrated on low-frequency ($f \ll f_{pi}$) MHD waves to couple Io to the DAM source regions. Marshall and Libby [1967] suggested that such waves might induce spin transitions in free radicals in the Jovian ionosphere, presumably because of the presence of off-diagonal terms in the pressure tensor, in analogy to the interaction of phonons with free radicals in crystals [Warwick, 1967].

The first quantitative estimates of MHD-wave coupling were made by Warwick [1967], who considered Io to be perfectly conducting and, therefore, argued that it would push aside the ambient field. He estimated that for "resonant" waves of phase velocity $V_{ph} \simeq V_{Io}$ (where $V_{Io} = 56$ km/sec is the speed of Io relative to the corotating plasma) and wavelength 2 R_{Io}, waves of amplitude $B_w = (V_A/\pi V_{Io})\ B_o$ might be produced, where V_A is the Alfvén velocity near Io. At the current estimate of $V_A \simeq 700$ km/sec, this would give $B_w \simeq 4\ B_o$. From other considerations, involving large but finite conductivity for Io, Warwick [1970] estimated $B_w \simeq 2.6\ B_o$. Waves of this magnitude are inherently nonlinear, and would be expected to steepen into shocks as they propagated toward Jupiter. Warwick [1970] estimated that about 4×10^{16} watts would thus be radiated from Io. He did not attempt to partition this energy among the various MHD modes (shear Alfvén, compressional Alfvén, and magnetosonic). Instead, he noted that if the MHD radiation from Io were isotropic, about 3×10^9 W would be intercepted by an area of 1.6×10^5 km^2 (the presumed size of the DAM source region) at the surface of Jupiter.

The estimates by Warwick [1970] were made partly in the context of a critique of the paper by Goldreich and Lynden-Bell [1969], which is discussed below. These latter authors assumed that the electric field induced across Io by its motion relative to the Jovian magnetic field would drive a current system that closed in the planetary ionosphere. In contrast, Warwick asserted that the current system would be closed within Io itself. This idea was also adopted by Schmahl [1970], who considered Io to have a permanent dipole moment $\vec{M}_{Io}$ induced by the DC average of its ambient magnetic field, and thus to be a point current source

$$\vec{J}_s(\vec{r}, t) = c\, \vec{M}_{Io} \times \nabla \delta (\vec{r} - \vec{V}_{Io}\, t) \quad .$$

Schmahl calculated that 8×10^9 W would be delivered to the Jovian ionosphere in the shear-Alfvén mode, which propagates essentially along the IFT.

In a contrasting approach, Goertz and Deift [1973] assumed that currents leak from Io into the ambient plasma. They attempted to calculate the steady flow pattern of the ambient medium around Io. The resulting distortions of the ambient magnetic field are to be understood in the sense of MHD wave fields. Some particularly striking features of their model are: (i) the perturbations of the field ahead of Io propagate to the Jovian ionosphere as solitons and are reflected as solitons, leading to a stretching of the field roughly analogous to the stretching of a string; (ii) the existence of an x-type neutral point behind Io (Figure 3). The x-type configuration is asserted to occur when the field lines leading Io are stretched too far; reconnection then takes place behind the moon, leading to the propagation of Alfvén waves up the IFT. In effect, Io "plucks" the magnetic field lines, like plucking a violin string. The angular limit of distortion before reconnection occurs is about 15°, accounting for the Io-phase asymmetry. Goertz and Deift estimate the maximum power radiated in Alfvén waves to be $\geq 2.7 \times 10^{15}$ W.

In Goertz and Deift's model, the wavelength of the Alfvén-waves generated by reconnection is taken to be 6 R_{Io} (i.e., $k \simeq R_{Io}^{-1}$). The resulting frequency of the waves is between 1-3 Hz, depending on the plasma density at Io. This density is determined from a magnetospheric model developed by Goertz [1973a]. Deift and Goertz [1973] consider the propagation of these waves along the IFT. They find a "transmission coefficient" for the waves into the Jovian ionosphere: the depth to which the Alfvén waves penetrate—and, therefore, the gyrofrequency at this depth—is a function of the wave frequency.

Goertz [1973b] equated the gyrofrequency at the local altitude of penetration, as a function of longitude, to the maximum DAM frequency at that longitude, showing good agreement with observation [Goertz and Haschick, 1972].

Goertz and Deift attempted to consider the three-dimensional nature of the MHD flow around Io. It is not clear, however, that they did so consistently. They assumed the conductivity of the ambient medium to be a scalar, whereas in a low-β plasma ($\beta \equiv 8\pi N K_B T_e / B^2$) the tensor properties of the conductivity are significant. Furthermore, even within the context of their model, the field lines which obey the geometry of Figure 3 are limited to those lines near the stagnation point of the flow; in principle, these lines degenerate to a set of measure zero, and it is then unclear to what extent their analyzed flow is dominated by the actual three-dimensional flow. In particular, the efficiency of flux reconnection at x-type neutral points may be drastically lower in three-dimensional flow than in two-dimensional flow. Furthermore, the back-flow towards Io associated with the reconnection is not taken into account.*

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*In defense of these authors, it must be noted that the three-dimensional MHD flow past a spherical body is an intractable problem, particularly in a low-β plasma. Nevertheless, the spirit of our comments is that a two-dimensional solution may be greatly misleading. In the common example, a tight-rope walker confined to the vertical plane could never fall off the tight-rope.

In addition, Deift and Goertz consider incompressible plane waves propagating along the IFT. It is not clear whether, in the situation of wave generation by flux reconnection, compressible modes or finite wave packets [Goldstein, Klimas, and Barish, 1974] might be more appropriate.

Finally, we remark that the entire discussion by Goertz and Deift [1973], Deift and Goertz [1973], and Goertz [1973b] is heavily dependent for its quantitative significance upon the magnetospheric model of Goertz [1973a].

None of the wave-coupling models discussed include a quantitative discussion of a consequent instability mechanism. In respect to the instability mechanism, however, the power available in the coupling mechanism is of vital importance because the efficiency of conversion to electromagnetic waves is necessarily low. We shall take 10^8 W to be a nomimal power requirement for electromagnetic waves in DAM bursts [Warwick, 1970]. Therefore, the power in the coupling mechanism must be $\gg 10^8$ W.

McCulloch [1967] estimates that 2×10^9 W may be transferred from Io via whistlers. Taken by itself this figure would indicate a conversion efficiency of 0.05, which is incredibly high. The presumed interaction mechanism, however, is the phase-organization of trapped particles (c.f. section III), and therefore the particles may be more important in the energy balance than the whistlers.

Because no instability mechanism has been demonstrated for coupling by low-frequency MHD-waves, we take the viewpoint that the conversion efficiency must be compatible with the power available in the MHD waves. (This viewpoint is consistent with the suggestion by Goertz [1973b] that the DAM radiation may originate from particles accelerated in the MHD-wavefronts.) In this light, Schmahl's estimate of $\simeq 10^{10}$ W in shear-Alfvén waves requires a very high conversion efficiency of 10^{-2}. On the other hand, the maximum-power estimates of 2.7×10^{15} W by Goertz and Deift and 4×10^{16} W by Warwick [1970] are much too high to agree with the recent determination by Goldstein [1975] that the secular acceleration of Io is limited to a value

$$\frac{\dot{p}}{p} < 3 \times 10^{-18}\ \mathrm{sec}^{-1} \quad ,$$

where p is the orbital period of Io. Because angular momentum is conserved in the system of Jupiter and its satellites, this finding limits the power dissipated by Io to about 10^{14} W. Although such a power is consistent with conversion efficiencies for DAM bursts of $\geqslant 10^{-6}$, it also indicates that the coupling mechanisms considered by Warwick [1970] and by Goertz and Deift cannot work at the maximum powers estimated by their authors.

It is certain, however, that Io must create some MHD disturbance in the magnetospheric plasma. Whether or not these waves are responsible for the coupling to DAM, they may have other significant effects. One such effect could be to provide local ionospheric heating near the base of the IFT. A crude calculation (Appendix A) yields an estimate that the waves might heat the ionosphere to temperatures of the order of 10^4 to 10^5 K. Because the plasma temperature may, under some conditions, be an important parameter in the contexts of instability mechanisms and wave propagation, such heating may have significant consequences for these aspects of the theory.

B. Acceleration Models

The starting point for acceleration models is that, owing to the motion of Io with velocity $|\vec{V}_I| = 56$ km/sec relative to the corotating plasma, an induced electric field $\vec{E}_{ind} = \vec{V}_I \times \vec{B}_I/c$ appears in Io's frame, where $\vec{B}_I$ is the ambient magnetic field at Io. If Io is not a good conductor, then this field would be attenuated only by volume polarization throughout its interior. Piddington and Drake [1968], however, suggested that Io might be highly conducting and so its interior would be screened from the induced field by surface polarization charges. As a result, they asserted that the induced emf of 400 kV could be mapped onto the field lines of the IFT, which are assumed to be equipotentials. The existence of the potential is communicated along the field line at the Alfvén speed $V_A = B_o/(4\pi\rho_m)^{\frac{1}{2}}$, where B_o is the magnetic field strength and

ρ_m the mass density of the plasma.* Provided that the conductivity σ is large enough so that the field does not diffuse through Io in a time shorter than the propagation time τ_A of Alfvén waves to the Jovian ionosphere, a DC current circuit which closes in the Jovian ionosphere will be established. The resistive diffusion time τ_m, over a scale length R_{Io}, is given by

$$\tau_m = 4\pi\,\sigma\,\frac{R_{Io}^2}{c^2} \quad .$$

The precise value of τ_A is not crucial in the present context. For purposes of estimating a constraint on σ, we may set $\tau_A \approx 60$ sec; this is the time required for Io to move a distance equal to its own diameter. Then for $\tau_m > 60$ sec, we require $\sigma > 1.5 \times 10^5\ \mathrm{sec}^{-1} = 1.7 \times 10^{-5}$ mho/m.

Goldreich and Lynden-Bell assumed sufficient conductivity for Io that the DC circuit model would be valid. They asserted that this requires $\tau_m > 2\tau_A$ as we have defined τ_A. In actuality, the DC circuit will be established for $\tau_m \sim \tau_A$, because only a one-way trip for an Alfvén wave is necessary to establish the potential in the ionosphere.

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*Strictly speaking, it is communicated at the phase velocity ω/k of an Alfvén wave, where $\omega/k = V_{ph} = c/(1 + c^2/V_A^2)^{\frac{1}{2}}$. For $V_A \ll c$, $V_{ph} \to V_A$; in the limit $V_A \to \infty$, $V_{ph} \to c$. The opposite limit, $V_A \to 0$, cannot be recovered from the Alfvén-wave dispersion relation.

(One may turn on a light switch and get light very quickly; no electron in the wires need traverse the entire circuit.) Goldreich and Lynden-Bell, however, actually required more than the existence of a circuit which closes in the Jovian ionosphere. They argued that if $\tau_m \gg 2\tau_A$, the entire IFT will be frozen to Io, only slipping in the ionosphere in order to follow rigidly Io's orbital motion. The importance of this assumption to Goldreich and Lynden-Bell was that it enabled them to argue that the field vanishes within the flux tube and that the current is carried in thin sheets on the boundary of the IFT. From this argument, and from some assumptions regarding the thickness of these sheets and the plasma density, they inferred a requisite energy for the current-carrying particles of a few keV.

Goldreich and Lynden-Bell assumed that there would be some potential drop along each field line. They then solved the current circuit equations to find the shape of the potential distribution across the IFT. Their solution has the correct form if one assumes that the current flow may be proportional to the voltage [Smith and Goertz, 1975]. For the potential to drop along the field line, however, requires some parallel resistivity. Goldreich and Lynden-Bell did not discuss the origin of this resistivity, and so their solution for the potential was left in terms of an unspecified constant V_o.

The models by Piddington and Drake and by Goldreich and Lynden-Bell are oversimplified in their assumptions regarding the MHD interaction of Io with the ambient plasma. Gurnett [1972] considered the effect of the formation of potential sheaths around Io. An object placed in a plasma attains a surface charge because the flux of thermal electrons hitting its surface is greater than the flux of ions. As a result, there develops a negative surface potential which is shielded from the ambient plasma by a Debye sheath. The equilibrium state is one in which the current to the surface is zero. If the surface emits a photoelectron current larger than the current it collects from the thermal flux, the potential at the surface becomes positive and a photoelectron sheath of the opposite polarity develops. Gurnett showed that in the case of a sunlit Io, part of which emits photoelectrons and part of which is in darkness, both types of sheaths occur. Their respective current densities J_D (for the Debye sheath) and J_p (for the photoelectron sheath) and areas A_D and A_p are related by the current balance condition:

$$J_p A_p = J_D A_D \quad .$$

Gurnett's model was considered in more detail by Hubbard, Shawhan, and Joyce [1974], and was extended to account for the existence of an ionosphere at Io by Shawhan <u>et al.</u> [1974]. The principal features of the latter modification are: (i) the conductivity required to close the current system is provided by the ionosphere of Io, rather than by the solid moon; (ii) the sheaths are assumed to form at some ill-defined height at the "top" of the ionosphere. The electron density profile of Io's ionosphere deduced by Kliore et al. [1974] shows a precipitous drop at an altitude of about 750 km, where the scale height changes abruptly from approximately 220 km to approximately 36 km. Shawhan et al. interpreted this as the sheath region. Adding this altitude to the radius of Io to give an "effective" radius for the conducting system, the total induced emf becomes 570 kV. The sheath model also assumes a current system that closes in the Jovian ionosphere; the model is schematically depicted in Figure 4.

Gurnett [1972] considered an idealized steady-state model of the current system (Figure 5) and concluded that if the height integrated Pedersen conductivity Σ_ρ of the Jovian ionosphere is lower than $\simeq$ 5 mho, the IFT would move with Io and the bulk of the potential drop in the external circuit would occur in the Jovian ionosphere, rather than in the sheaths. Conversely, Gurnett asserted that if $\Sigma_\rho \geqslant 5$ mho, the IFT slips by Io and substantial voltage drops occur across the sheaths. After noting these conclusions, Hubbard et al. [1974] assumed that the major potential drops occurred across the sheaths [c.f. their Eq. (6)]. This assumption was also made by Shawhan et al. [1974] and by Shawhan [1975]. All of these authors assumed that particles traversing the sheaths would gain potential energy equal to the sheath potential.

Gurnett estimated J_D and J_p independently, as did Hubbard et al. Similarly, Hubbard et al. estimated the sheath thicknesses as independent parameters, and similar estimates were made by Shawhan *et al.* These latter authors, however, recognized the sheaths to be potential double layers, which have been investigated by Block [1972], Carlqvist [1972], Knorr and Goertz [1974], and Goertz and Joyce [1975]. Such layers form when electrons and ions in a plasma with $T_e \leqslant T_i$ drift relative to one another with a velocity $U \geqslant V_e$, where V_e is the electron thermal velocity. A potential Φ_o then develops over a scale length L. Smith and Goertz [1975] noted that the development of the double layer was subject to the Buneman instability [Buneman, 1958], which produces acoustic turbulence in the frequency range $\omega_i \ll \omega \ll \omega_e$. The convection of the turbulence at the group velocity, which is given by $V_{g,B} \simeq 0.4\ U$ for the most unstable waves, leads to a natural scale length L for the sheaths, given by

$$L \simeq \frac{V_{g,B}}{\gamma_B} \quad ,$$

where $\gamma_B \simeq \omega_e/18$ is the maximum growth rate of the Buneman instability. Smith and Goertz also showed that the limiting drift velocity U was related to the sheath potential Φ_s by

$$\frac{1}{2}\, m\, U^2 \leqslant \frac{1}{2}\, e\, \Phi_s \quad .$$

By considering a steady-state current system similar to that analyzed by Gurnett, Smith and Goertz showed that the maximum potential drop across the sheaths is limited to about 140 kV.

We stress that all of these analyses regarding the sheath model assume the current circuit to be closed in the Jovian ionosphere. The self-consistent system of drift currents in the ionosphere of Io, the sheath regions, and the ambient plasma into which Io is penetrating (and which is therefore magnetodydrodynamically perturbed) is not analyzed, and it has not been demonstrated conclusively that the current system does not close locally near Io. We note, however, that the current system cannot close in a region which moves with Io, because in such a region $\nabla \times \vec{j} = 0$.

Related to this point is the question of whether the conductivity of Io's ionosphere is itself sufficient to carry the current without causing a large drop of the emf. Shawhan et al. [1974] estimated the conductance of Io's ionosphere to be 260 mho, a value sufficient by more than an order of magnitude. The atmospheric model of McElroy and Yung [1975] is also compatible with a conductance of this order. This model assumes N_2 and NH_3 to be the primary atmospheric constituents. Whitten, Reynolds, and Michelson [1975] developed a model in which the primary constituent is neon. In this model, the ionospheric structure inferred by Kliore et al. [1974] is obtained with a much smaller neutral density than in McElroy and Yung's model, and the ionospheric conductivity is correspondingly lower, perhaps becoming marginal with regard to the requirements of the sheath model [Shawhan, 1975].

C. The Sweeping Model

The MHD-wave and acceleration models discussed above all entail the assumption that Io is a good conductor. In contrast, Wu [1973] assumed that the conductivity of Io was sufficiently low that the Jovian magnetic field would diffuse readily through it. In that case, particles trapped on the field lines threading Io would impinge on the moon and be "swept" from the magnetosphere. Sweeping had been previously discussed by Mead and Hess [1973] in the context of its effects on the structure of the Jovian radiation belts.

Energetic particles trapped in a dipole magnetic field execute a complete bounce, from a given mirror point to the conjugate mirror point and back again, in the time

$$\tau_b \simeq (1.3 - 0.56 \sin \alpha_o) \, 4L \frac{R_J}{v} \quad ,$$

where α_o is the equatorial pitch angle, L is the conventional magnetic shell parameter, and v is the speed of the particle. For particles at $L \simeq 6$, mirroring at gyrofrequencies corresponding to the DAM frequency range, $\alpha_o \leqslant 0.1$ rad, and so $\tau_b \simeq 30\ R_J/v$. In the frame of reference corotating with Jupiter, Io moves a distance equal to its own diameter in a time $\tau_s \simeq 60$ sec. Therefore, particles that execute a half-bounce in a time shorter than τ_s will impinge upon Io and be removed from the field lines described by their guiding center motions. The threshold velocity v_c above which particles will be swept in this way is determined by equating $\tau_b(v_c)/2 = \tau_s$. For L = 6, $v_c \simeq 1.8 \times 10^4$ km/sec, corresponding to threshold energies of approximately 0.8 keV for electrons and 1.6 MeV for protons. Particles above the appropriate threshold are effectively swept from the Io flux tube.

In contrast to the instability mechanisms associated with other coupling mechanisms, that due to the sweeping effect depends on the energetic protons, rather than electrons. In section III we shall consider the form of the proton distribution function; here it suffices to note that it includes as a parameter a characteristic energy E_p.

The sweeping of energetic particles by Io results in a cavity in the density distribution of such particles. In the orbital plane of Io, the cavity has a cross-section similar to a wake and remains for some time after the passage of Io; diffusion across the cavity boundary tends to replace the energetic particles on some time scale τ_D, long compared to the growth time of the electromagnetic instability. The cavity cross-section in Io's orbital plane is illustrated in Figure 6. The density of particles with energies above the sweeping threshold falls to zero across the distance δ. The bow portion of the cavity where the sweeping is incomplete has a characteristic dimension ℓ given by

$$\ell = \frac{\tau_b(E_p)}{2} V_{Io} \quad ;$$

thus the bow of the completely swept portion has a cross-section which is a semi-circle displaced from the bow of Io by the distance ℓ. The dimension δ of the lateral boundary layer is given by the intersection of this profile with the trailing limb of Io:

$$\delta = R_{Io} \left[1 - \cos \sin^{-1} \left(\frac{\ell}{2 R_{Io}} \right) \right] \quad . \tag{II-1}$$

Equation (II-1) is valid as long as it gives a value for δ greater than the characteristic proton Larmor radius $R_p = (E_p/M\Omega_p^2)^{\frac{1}{2}}$, where $\Omega_p = 2\pi f_{cp}$ is the angular proton gyrofrequency; R_p is the physical lower bound on δ. The equatorial boundary thickness δ maps to values $\delta'(\lambda) < \delta$ at magnetic latitudes $|\lambda| > 0$ (Figure 7):

$$\delta'(\lambda) = \delta - \int_0^{\lambda} \frac{d\delta}{d\lambda'} d\lambda' \quad .$$

Assuming $R_p(0)/\delta \ll 1$, we have $R_p(\lambda)/\delta' < 1$ also, and the boundary thickness at higher latitudes is given approximately by the mapping for a centered collinear dipole:

$$\frac{\delta'}{\delta} = \frac{\cos^3 \lambda}{(4 - 3 \cos^2 \lambda)^{\frac{1}{2}}} \quad .$$

Owing to the density gradient of energetic protons across the walls of the IFT, drift currents appear. These currents drive instabilities at the upper and lower hybrid resonances; these instabilities are discussed in later sections.

D. Evidence From Pioneers 10 and 11

The trajectories of Pioneers 10 and 11 did not pass through the instantaneous IFT, and so the evidence from these probes regarding the coupling mechanisms discussed above is indirect and inferential. Inbound, Pioneer 11 passed about 6000 km from the instantaneous IFT as defined by the D_2 field model of Smith _et al._ [1975] at about 0300^h on 3 December 1974 [Fillius, McIlwain, and Mogro-Campero 1975]. The spacecraft was then at approximately -40° magnetic latitude, and so its distance from the IFT was about twenty times the radius of the flux tube itself at that latitude. No distortions of the magnetic field near this region have been reported in the literature, indicating that the MHD perturbations arising from Io's penetration of the magnetospheric plasma are either unresolvable or are localized at this latitude to distances much less than 6000 km from the IFT. Thus, it is unlikely that Io warps its flux tube by 15°, as suggested by Goertz and Deift. No direct evidence about the actual MHD nature of the Io-magnetosphere interaction can be obtained from trajectories such as those of the Pioneers, however.

The particle data from both Pioneers 10 and 11 show that sweeping occurs and is effective. In this regard, it must be remembered that sweeping of the IFT is followed by diffusion of the energetic particles back into the swept region. The Pioneer 10 data were often interpreted to indicate that sweeping is much less effective than predicted by Mead and Hess; typical reductions in the particle fluxes at the IFT observed on Pioneer 10 were by factors of about two to five, and the sweeping effect seemed generally to be confined to lower-energy particles. Mead and Hess, however, considered only the slow processes of electric and magnetic diffusion, as they are understood to operate in the terrestrial magnetosphere, to balance the sweeping effects of the various Jovian satellites. They were led to predict large decreases, by several orders of magnitude, of the trapped particle fluxes at the orbits of each of the moons. Smith [1973], however, showed that the sweeping effect leads not only to the high-frequency electromagnetic instability described by Wu, but also to a low-frequency instability which produces strong diffusion ocross the boundaries of the IFT. (To the extent that sweeping is effective at other satellites, the same mechanisms operate there also.) Moreover, the diffusion coefficient is energy-dependent, increasing with the particle energy. Thus, the relative mildness of the sweeping effect as observed in the Pioneer 10 data may be attributed to the fact that the satellite crossed the orbit of Io several hours behind the satellite, and the density gradient had been eliminated in the case of higher-energy particles and considerably reduced at lower energies. Further evidence of this interpretation may be found in the Pioneer 11 data. The trajectory of Pioneer 11 was such that the spacecraft passed the Io orbit much closer to Io, both inbound and outbound, than did Pioneer 10; sweeping was found by all the particle experiments to be quite severe.

There is clear evidence for acceleration of electrons at Io, but its interpretation in terms of existing models is ambiguous. On both the inbound and outbound crossings of Pioneer 10, Fillius and McIlwain [1974] observed enhancements, by a factor of about two, in the electron flux at energies $E > 0.16$ Mev. These enhancements were observed just inside the orbit of Io, in accord with the sheath-model predictions of Shawhan et al [1973]. On the inbound crossing of Pioneer 11, Fillius et al. [1975] observed a dramatic increase -- by a factor of thirty -- in the electron flux at $E > 0.46$ Mev, an energy compatible with the interpretation of the sheath model by Shawhan et al. [1973,1974]. This "spike", however, occurred outside the position of Io's orbit as inferred from the location of the sweeping minimum. This is in contrast to the prediction of the sheath model, in which electrons are accelerated away from Io by the Debye sheath on the inward face of Io and accelerated toward Io, striking its surface, in the outward face. Thus, it is not clear that the evidence for particle acceleration at Io may be regarded as favorable evidence for the sheath model, although there exists the possibility that the combined effects of drift motions and diffusion may result in a displacement of the spike across the sweeping minimum. At this writing, however, the origin of these flux enhancements at Io seems ambiguous.

III. THE INSTABILITY MECHANISM

There is some variability in the literature in the nomenclature used for the various normal modes of plane waves in plasma. Therefore, we shall first define the nomenclature to be used here. We shall not be concerned with low-frequency modes for which the ion motion is important. In the high-frequency regime there are two normal modes in a uniform magnetized plasma. These are designated as the ordinary (o) and extraordinary (x) modes. The refractive indices $n_{o,x}$ ($\equiv ck/\omega$) of these modes are given by the well-known Appleton-Hartree formula:

$$n^2_{o,x} = 1 - \frac{2\omega_e^2(1 - \omega_e^2/\omega^2)}{2(\omega^2 - \omega_e^2) - \Omega_e^2 \sin^2\theta \pm \Omega_e \Delta} , \qquad \text{(III-1)}$$

where $\omega_e = 2\pi f_{p,e}$, $\Omega_e = 2\pi f_{c,e}$, θ is the angle between the wave vector $\vec{k}$ and the ambient magnetic field $\vec{B}_o$, and $\Delta \equiv [\Omega_e^2 \sin^4\theta + 4(\omega^2 - \omega_e^{\,2}) \cos^2\theta]^{\frac{1}{2}}$. The upper sign refers to the o-mode, and the lower sign to the x-mode. If variations in the magnetic field or plasma density occur on a length scale much greater than k^{-1}, Eq. (III-1) may be taken to describe the local propagation of the waves.

The ratio $\epsilon \equiv \omega_e^2/\Omega_e^2$ is an important parameter in determining the nature of the normal-mode propagation. Throughout the region of the Jovian plasmasphere of interest for DAM, $\epsilon \ll 1$. This is an unusual parameter regime in both laboratory and astrophysical plasmas; the only other known region in the solar system in which $\epsilon \ll 1$ is the high-latitude region above the terrestrial plasmasphere. Interestingly enough, this is the source region of the terrestrial kilometric radiation [Gurnett, 1974; Kaiser and Stone, 1975].

In Figure 8 we show the topology of the dispersion curves ω versus k and the refractive indices n^2 versus ω of the o- and x-modes for propagation at large angles to the magnetic field. The x-mode is seen to consist of two branches, the slow branch (over most of which $n_x^2 > 1$) and the fast branch ($n_x^2 < 1$). Both modes exhibit cutoffs, occurring at ω_e for the o-mode and $\omega_x = \Omega_e \{[1 + 2\epsilon \pm (1 + 4\epsilon)^{\frac{1}{2}}]/2\}^{\frac{1}{2}}$ for the x-mode, where the upper (lower) sign refers to the fast (slow) branch. For the fast x-branch, the cutoff frequency is approximately given by $\omega_x = \Omega_e(1 + \epsilon)$. In addition, the slow x-branch exhibits a resonance at $\omega = \Omega_e(1 + \epsilon/2 \sin^2 \theta)$.

At $\theta = 90°$ the o-mode is linearly polarized and the x-mode is right elliptically polarized, with $\vec{E} \perp \vec{B}_0$. As $\theta \to 0$, the x- and o-modes go over continuously to right and left circular polarizations, respectively. The portion a-b of the slow branch is the whistler regime.

The x-mode cannot propagate in the "stop band" between the slow-branch resonance and the fast-branch cutoff. The axes of Figure

8 are normalized to Ω_e; for a fixed wave frequency, increasing ω/Ω_e generally corresponds to moving towards lower magnetic fields. Therefore, we see that waves on the slow x-branch may not propagate from the Jovian plasmasphere into interplanetary space. Strictly speaking, the stop band exists only in the cold-plasma limit; thermal effects remove the resonance and lead to a non-zero minimum in the group velocity of the slow branch [Aubrey, Bitoun, and Graff, 1970]* The finite temperature of the plasma also implies a collision frequency ν_c, however, and waves with extremely slow group velocity V_g are attenuated by collisions over a scale length V_g/ν_c. If the group velocity is sufficiently small, or the spatial region corresponding to the stop band (the "stop zone") is sufficiently extensive, the attenuation in traversing the stop zone may still be large. Thus, as a general rule the existence of the stop band must be taken into account in assessing the propagation of waves. In particular, instability mechanisms in which the waves are amplified on the slow branch also imply the necessity of some mode-coupling process leading to the escape of waves in either the o-mode or the fast x-mode.

The cold-plasma approximation may be deduced from the more rigorous kinetic theory in the asymptotic limit of low temperatures, under the conditions

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*Other thermal effects on the dispersion relations are negligible for our purposes, provided $V_e^2/c^2 \ll \epsilon$. Although $\epsilon \ll 1$, the low temperature of the Jovian plasma ensures the validity of this inequality.

$$\frac{k_\perp V_e}{\Omega_e} \to 0 \quad , \tag{III-2a}$$

$$\left| \frac{\omega - \ell\,\Omega_e}{k_\parallel V_e} \right| \gg 1 \qquad (\ell = 0, 1, 2, \ldots) \quad , \tag{III-2b}$$

where $k_\perp = k \sin\theta$, $k_\parallel = k \cos\theta$. Equation (III-2a) is the condition that finite-Larmor-radius effects may be neglected. Equation (III-2b) is the condition that the thermal electrons not exhibit Landau and cyclotron resonance with the wave as they spiral along field lines. Such resonances may quickly damp the wave.

Inequalities (III-2) arise from the inclusion of thermal effects in the plasma dielectric tensor. In the cold-plasma limit the dielectric is Hermitian, giving refractive indices which are either purely real (oscillatory solutions) or purely imaginary (exponentially damped solutions). Thermal effects modify somewhat the Hermitian part of the dielectric, which determines the normal-mode frequencies. As noted above, however, these modifications are of order V_e^2/c^2 for the electromagnetic modes, and are generally negligible. The thermal effects also introduce a non-Hermitian part to the dielectric. It is this part which implies damping of the normal modes; the damping is weak under the conditions (III-2).

In the same manner, the inclusion of a population of energetic particles also modifies the plasma dielectric, in general introducing a frequency shift into the real part of the refractive index and, in addition, an imaginary part to the frequency, which may then be written

$$\omega(\vec{k}) = \omega_r(\vec{k}) + i\gamma(\vec{k}, \omega_r) \quad . \tag{III-3}$$

The usual approach adopted in linear instability analysis is to consider an energetic particle population of low density relative to the thermal background plasma. Then the real frequency ω_r is governed by the cold-plasma dispersion relation, while the growth or damping of the waves is determined by the energetic particle distribution. The dispersion relation is then of the form

$$D(\vec{k}, \omega) = C(\vec{k}, \omega) + \sum_s g_s(\vec{k}, \omega) \quad , \tag{III-4}$$

where $C(\vec{k}, \omega)$ is the cold-plasma dispersion relation and $g_s(\vec{k}, \omega)$ represents the contribution of energetic particles of species s. In principle, the thermal damping effects may be included in $\Sigma_s\, g_s$, but this is not usually done; in particular, it is not done in the theories we shall discuss here. Therefore, the tendency for instability must be compared a posteriori with conditions (III-2), to determine that the instability is stronger than the competing damping by thermal particles.

All of the instability mechanisms we shall consider depend on an anisotropic distribution of a small energetic-particle population. Instabilities may exist in various parameter regimes in both ω_r and in the parameters that describe the anisotropy, e.g., $T_\perp/T_\parallel$. The nature of the instability may be closely related to the nature of the coupling, and some coupling mechanisms are compatible with only certain types of instability mechanisms. To some extent, individual coupling mechanisms have generally been proposed with some particular type of consequent instability mechanism in mind, even when the coupling and instability have not both been investigated by the same author. Conversely, in some cases a proposed instability mechanism includes the tacit assumption that there exists an unspecified coupling mechanism which produces either the required anisotropy in the energetic-particle distribution or the conditions necessary for an existing marginally stable anisotropy to be driven unstable to electromagnetic wave amplification.

It is sometimes asserted in the literature that the DAM emission mechanism must be cyclotron emission. There is no compelling evidence, either observational or theoretical, that this is the case. Goldreich and Lynden-Bell argue for coherent cyclotron emission primarily on the basis of the sharp beaming pattern of the emission, and also on the high intensity. Such beaming, however, is an inherent feature of all of the mechanisms discussed below. In addition, there is the complicating factor that, if amplification occurs on the slow x-branch, the beaming pattern may be closely related to or determined by the mode-coupling process by which the DAM escapes from Jupiter. Thus, cyclotron emission is only one of the possibilities for the instability mechanism.

Before discussing individual instability mechanisms, it should be noted that a complete understanding of DAM may involve more than one mechanism. There are at least two distinct types of Io-modulated DAM, the decasecond (L) and millisecond (S) bursts. They have quite distinct time scales—from about 0.5-3 sec for L-bursts and 1-10 msec for S-bursts—which do not merge continuously into one another. Moreover, the S-bursts exhibit true frequency drifts, which are always negative. In contrast, L-bursts exhibit frequency drift only in their overall envelope; this is probably a geometric effect of the beaming pattern. Although both L- and S-bursts share the same general morphology of Io-modulation, these facts imply that they may be caused by separate instability mechanisms.

It is also unclear whether the Io-modulated and Io-independent L-bursts imply different emission mechanisms. With the recent recognition of an Io-independent component of the early source region [Carr and Desch, 1975; Bozyan and Douglas, 1975], one might be tempted to hypothesize that the Io-modulated DAM is simply an enhancement of a weak continuous emission. It would then seem necessary, however, to shift from the common interpretation that the morphology of the Io-modulation implies a beaming pattern, and instead to suppose that the Io-modulation is indeed only possible for certain geometrical configurations relative to the observer. This latter hypothesis is physically untenable. Moreover, the declination effect in CML is related to the Io-independent DAM [Goertz, 1971], while the modulated component exhibits a declination effect in Io phase [Conseil, 1972; Lecacheux, 1974; Thieman, Smith, and May, 1975; Bozyan and Douglas, 1975]. Thus, it seems probable that the Io-independent and -modulated components are due to distinct mechanisms, although these arguments do not conclusively rule out the possibility that they may be closely related.

Thus, in assessing the viability of a given instability mechanism, one should attempt to do so in the several contexts of Io-independent L-bursts and Io-modulated L- and S-bursts. In the first of these contexts, the lack of a demonstrable coupling mechanism is irrelevant.

Finally, we note that the mechanisms we shall discuss have generally been analyzed only in the linear regime. Linear instability analysis may be quite misleading, however, because it provides no information on the saturation level of the instability. In general, a complete stability analysis must include consideration of the nonlinear regime.

A. The Sweeping Model

Of the instability mechanisms we discuss, the sweeping model of Wu [1973] is unique in that the instability is driven by anisotropy in configuration space rather than in velocity space. As was mentioned in section II, the density gradients $\nabla N_{e,i}$ of trapped energetic particles across the boundary of the IFT produce drift currents in the cavity "wall"; these currents are perpendicular to both $\vec{B}_o$ and N. Wu considered the motion of each energetic particle species transverse to $\vec{B}_o$ to be characterized locally by a mean energy $E_{\perp,s}$. The resulting drift velocity of the s-th species is then given by

$$\langle v_{x,s} \rangle = - \frac{q_s}{|q_s|} \frac{1}{\Omega_s} \left(\frac{d \ell n N_s}{dy} \right) \frac{E_{\perp,s}}{m_s} \equiv v_{D,s}, \qquad \text{(III-5)}$$

where q_s is the charge, the density gradient is taken to be positive in the y-direction, and the drift motion is in the +x-direction for electrons and the -x-direction for protons.

Wu considered instability at the upper hybrid frequency ω_{UH} and for very short wavelengths, $ck/\Omega_e \gg 1$. Under these conditions, the energetic electrons contribute only a negligible shift in the real frequency, and the instability is produced by the non-resonant contribution of the protons. Taking an isotropic gaussian for the local distribution of the energetic protons, Wu found the growth rate of waves with $\vec{k} \perp \vec{B}_o$ to be given by

$$\gamma = \frac{1}{(8\pi)^{\frac{1}{4}}} \frac{1}{(kR_p)^{\frac{3}{2}}} \left(\frac{N_p}{N_e} \frac{M}{m} \epsilon \right)^{\frac{1}{2}} \left(\frac{k_x v_D}{\Omega_e} - 1 \right)^{\frac{1}{2}} \omega_e \quad , \qquad \text{(III-6)}$$

where subscripts p,e denote energetic protons and thermal electrons, respectively, $R_p = (E_{\perp,p}/M)^{\frac{1}{2}}/\Omega_p$, and k_x is the x-component of the wave-vector. Maximum growth occurs at

$$\frac{k_x v_D}{\Omega_e} \simeq \frac{v_D}{c}\frac{ck}{\Omega_e} = \frac{3}{2} \quad .$$

Taking $v_D = E_\perp / M\Omega_p \delta'$, where δ is the thickness of the density-gradient layer (see section II-C), one finds v_D in the range 1-4 x 10^3 km/sec for 5 MeV $\leqslant E_\perp <$ 10 MeV. Thus, the wavenumber at maximum growth is of order $cK/\Omega_e \simeq 100$. Wu did not address the question of how the waves penetrated the stop band.

Wu's analysis was done for $\vec{k}$ strictly perpendicular to $\vec{B}_o$. Because the analysis was a local one, it also requires a finite component of k in the y-direction, so that

$$|k_y| \left|\frac{\partial \ln N_p}{\partial y}\right|^{-1} \gg 1 \quad . \tag{III-7}$$

Although it is clear from Eq. (III-6) that γ increases as k_x/k increases, the inequality (III-7) is not sufficient to allow a precise determination of the azimuthal width of the beaming pattern; at best, one can say only that it is a pencil beam. In the direction parallel to $\vec{B}_o$, the width of the pencil beam is sharply delimited by Eq. (III-2b), which requires

$$\sqrt{\left|\frac{k_\parallel}{k}\right|} \ll \frac{\epsilon}{2}\frac{1}{kR_e} = \frac{\epsilon}{2}\frac{c}{V_e}\left(\frac{\Omega_e}{ck}\right) \quad .$$

The current-driven instability also contains the feature of a low-frequency cutoff, because an essential requirement for the instability to occur is that, locally, $\epsilon > 2m/M \simeq 10^{-3}$. In general, ϵ decreases with increasing altitude, at least until the point of maximum zenopotential height [Melrose, 1967; Smith, 1973; Goertz, 1975].

The power available in this mechanism was estimated by Wu [1973] to be about 10^9 W, adequate for the DAM requirement. Smith [1973] estimated the power to be $10^9\eta$ W, where $\eta \ll 1$ is an efficiency factor. Smith also pointed out that the estimated requirement of $\simeq 10^8$ W depends on an assumed beaming into a solid angle of approximately 0.1 sr; from the above discussion this solid angle might be much too large in the context of the sweeping instability. Both of these estimates assumed growth rates of order $\gamma \sim 10^{-2}$ sec^{-1}. Because the growth rate depends on the energetic-proton density the precise value of the estimated power may have little significance. The ability of the mechanism to meet the power requirement of Io-modulated DAM, however, seems marginal.

Comparison of the predictions and requirements of the sweeping model with Pioneer observations tends to cast doubt on the viability of the mechanism for Io-modulated DAM. As was noted above, sweeping seems to occur effectively enough to establish the necessary density gradient. Because the mechanism relies on trapped protons, however, and the higher-order moments of the magnetic field make the surface field highest in the northern hemisphere, the instability is limited at any longitude to the region between the topside Jovian ionosphere in the southern hemisphere and the conjugate mirror point. Because the growth rate of the instability is proportional to N_e, the emission is essentially limited to the southern hemisphere. In order to span the range of DAM frequencies, therefore, the peak ionospheric electron density, N_{eo}, must be sufficient to raise the upper hybrid frequency in the ionosphere up to about 40 MHz. The maximum gyrofrequency at the foot of the IFT is about 27 MHz, giving the requirement that $N_{eo} \gtrsim 10^7$ cm^{-3}. The maximum density inferred by Fjeldbo et al. [1975], however, is of order 3 x 10^5 cm^{-3}. Thus, unless this determination is too low by a factor of 30, the sweeping mechanism is essentially limited to frequencies below 30 MHz.

Other aspects of the sweeping mechanism will be discussed below. In particular, it may be quite relevant to the question of the "Io-independent" DAM.

B. "Maser-Like" Mechanisms

A number of authors have considered wave amplification by electron distributions which are anisotropic in velocity space. Such mechansim are often called "maser-like", because the anisotropic distribution is analogous to an inverted level population. The most commonly analyzed situation is that in which the mean perpendicular energy $T_\perp$ is greater than the mean parallel energy, $T_{\parallel}$.

Chang [1963] discussed the amplification of x-mode waves by an anisotropic distribution of relativisitic electrons. Although he referred to these waves as whistlers, his analysis was in fact valid for the entire x-mode. He only considered the cases of propagation parallel and anti-parallel to $\vec{B}_o$, however. Chang found two parameter regimes in which amplification could occur. Of these, the only case relevant to the Jovian exosphere, in which $\epsilon \ll 1$, is the frequency regime

$$\omega_e^2 \gg \Omega_e \omega - \omega^2 > 0 \quad , \qquad \Omega_e/\omega \simeq 1 \quad . \tag{III-8}$$

Because $\epsilon \ll 1$, however, inequality (III-8) is equivalent to the condition

$$\frac{\omega}{\Omega_e} = 1 - \xi \quad , \qquad \xi \ll \epsilon \quad .$$

Thus, as in the case of the sweeping model, Chang's mechanism requires $ck/\Omega_e \gg 1$ (see Figure 8). For parallel propagation, however, these waves are strongly damped by the thermal electrons because, for $\xi \ll \epsilon$ and $\epsilon \ll 1$, we have

$$\left| \frac{\omega - \Omega_e}{k_{\parallel} V_e} \right| \ll \epsilon \frac{c}{V_e} \frac{\Omega_e}{ck} \ll 1 \quad ,$$

in violation of inequality (III-2b).

In an analysis complementary to that of Chang, Goldstein, and Eviatar [1972] discussed amplification at $\theta = 90°$ on the slow x-branch by a loss-cone distribution of trapped electrons. Goldstein and Eviatar, however, explicitly recognized the necessity for the waves to penetrate the stop zone. They argued that the amplification would be sufficiently strong so that the radiation would simply tunnel through the stop band, suffering exponential attenuation but emerging into the outer region of the exosphere with observable intensity. They asserted that this mechanism could be responsible for the Io-independent DAM. As we shall discuss below, the stop zone is probably sufficiently thin that exponential attenuation of waves incident upon it is negligible. Assuming this to be so, the instability analysis of Goldstein and Eviatar might account for Io-independent DAM by emission from a loss-cone distribution near marginal stability.

Hirshfield and Bekefi [1963] considered the amplification of x-mode waves at Ω_e, with $\theta = 90^o$. In their model the distribution function was assumed to be isotropic; the "maser-like" effects were due to the dependence of the gyrofrequency on the relativisitic electron mass. This dependence leads to coherent interaction between the particles and the wave field through phase synchronization, such as occurs in synchrotron accelerators [Kuckes and Sudan, 1971].

Although it was apparently not recognized by Hirshfield and Bekefi, the unstable waves at the gyrofrequency are below the stop band. They also predicted amplification at higher harmonics of the gyrofrequency, but at unobservable levels. Moreover, their work was done before the discovery of the Io-modulation, and so the distribution function used in their analysis was based on the assumption of direct access to the Jovian magnetosphere of streams of low-energy particles of solar origin.

Goldreich and Lynden-Bell also postulated coherent cyclotron amplification for the emission mechanism, but recognized the problem of traversing the stop zone. They discussed the instability as seen in the frame of the streaming particles; in this frame their results are essentially similar to those of Hirshfield and Bekefi. Goldreich and Lynden-Bell asserted that in Jupiter's frame the waves would be Doppler-shifted above the stop band and so would escape the magnetosphere. This conclusion was erroneous, however, because a wave with phase velocity smaller than c in one frame will not have phase velocity greater than c in any other frame. Thus, the Doppler-shifted waves postulated by Goldreich and Lynden-Bell would not be normal modes in the magnetospheric plasma, and would be heavily damped away from the source region.

All of the mechanisms so far discussed provide for amplification on the slow x-branch. With the exception of the sweeping-model instability, which is non-resonant and is carried by protons, these mechanisms all invoke resonant-electron effects. These are possible only on the slow x-branch, where the phase velocities of the waves may be exceeded by particle velocities. For the o-mode and the fast x-branch, however, the phase velocities are greater than c, and resonant interaction is impossible. There are, nonetheless, maser-like mechanisms to amplify these waves; they are sometimes called subluminous instabilities. Such mechanisms were hypothesized by Ellis [1962,1963,1965] and by Ellis and McCulloch [1963], who assumed the existence of electron "bunches" precipitating into the Jovian ionosphere. The mechanism was investigated quantitatively in a series of papers by Fung [1966a,b,c]. The first of these papers [Fung, 1966a] was concerned with terrestrial VLF emission and so neglected terms of order Ω_e^2/ω_e^2, which are large in the Jovian exosphere. Fung [1966b] extended the analysis to include these terms. Both of these analyses assumed precipitating electron distributions of the form

$$F_e(p) \sim \delta(p_{\parallel} - p_{\parallel o})\ \delta(p_{\perp} - p_{\perp o}), \qquad \text{(III-9)}$$

and allowed for arbitrary angles of wave propagation. In his third paper, Fung [1966c] extended the formalism to include finite temperatures $T_{\perp}$ and $T_{\parallel}$ in the precipitating stream.

In this mechanism the emission is Doppler-shifted to a normal mode above the stop-band by the streaming velocity of the particles. Radiation at the fundamental is confined to a finite cone, in the forward direction, of wave-vector directions about the magnetic field, while emission at higher harmonics may occur in all directions but is much weaker.

Melrose [1973, 1975] also considered subluminous instability from a streaming anisotropic distribution, which he assumed to be a shifted bi-Maxwellian:

$$F_e(\beta, \alpha) = \frac{1}{(2\pi)^{3/2}\beta_{\perp o}^2 \beta_{\parallel o}} \exp\left[\frac{-\beta^2 \sin^2 \alpha}{\beta_{\perp o}^2} - \frac{(\beta \cos \alpha - \beta_s)^2}{\beta_{\parallel o}^2}\right] . \tag{III-10}$$

Here $\beta = v/c$, α is the particle pitch angle, and $\beta_s = \langle\beta \cos \alpha\rangle$ is the normalized streaming velocity. The Doppler-shifted frequency of the s-th harmonic, radiated in the direction θ by a particle of pitch angle α and Lorentz factor γ, is

$$\omega = \frac{s\Omega_e}{\gamma(1 - n_x \beta \cos \alpha \cos \theta)} .$$

At the fundamental, the condition that the radiated wave be Doppler-shifted above the stop band is then

$$n_x \beta_s \cos \theta > \epsilon ,$$

which is not excessively stringent in view of the smallness of ϵ.

Melrose finds a maximum growth rate of order

$$\gamma_{max} \simeq \frac{\eta\epsilon}{n_x \, \beta_{\parallel o}} \, \Omega_e \quad , \tag{III-11}$$

where $\eta \equiv N_s/N_e$ is the ratio of the stream density to that of the background plasma. The bandwidth of the growing waves is of order

$$\frac{\Delta\omega}{\Omega_e} \simeq n_x \, \beta_{\parallel o} \ll 1 \quad ,$$

which is in qualitative accord with observations.

Melrose finds a threshold condition on the degree of anisotropy necessary in order for amplification to occur; the condition is

$$\beta_{\perp o}^2 > \beta_{\parallel o} \quad . \tag{III-12}$$

Although this inequality may seem stringent, it is actually fairly easy to meet under certain models for the acceleration of the particles. For instance, the accelerated spectrum in the sheath model is essentially a delta-function in energy at any given point at Io; the particular injection energy depends on the spatial structure of the sheath potential [Smith and Goertz, 1975]. Because the energy is in general much larger than the temperature of the particles before acceleration, the accelerated distribution will map at high latitudes essentially to a distribution of the form (III-9) assumed by Fung.

The parallel motion may be equated locally to Melrose's β_s, while his $\beta_{\perp o}$ may be equated to the characteristic transverse velocity. In the limit of this model, inequality (III-12) is satisfied virtually trivially. This condition appears to be an artifact of the assumed form (III-10) for the distribution function; Fung [1966b] found instability for distributions of the form (III-9) in which $\beta_\perp^2 < \beta_o$.

Melrose [1975] used quasilinear arguments to estimate the saturation energy density W_w of the amplified waves as

$$\frac{W_w}{W_s} = \frac{2}{\gamma_{max}} \frac{\beta_s \; 2 \; \beta_{\perp o}^2}{L_B \left(\beta_s^2 + \beta_{\perp o}^2 \right)} ,$$

where L_B is the gradient scale length of the magnetic field and $W_s = N_s \, mc^2(\beta_{\perp o}^2 + \beta_s^2)/2$ is the energy density in the stream. He evaluated this for two different sets of beam parameters, obtaining results which bracket an observational estimate of 4×10^{-12} erg cm^{-3} required. This latter estimate is made using a very narrow beaming cone of solid angle 0.01 sr. In general, however, the energetic requirement of DAM appears to be well met within the context of Melrose's calculations.

The maximum growth rates predicted by Melrose's theory are of the order of $10^2 \leqslant \gamma \leqslant 10^3$ sec^{-1}. Thus the mechanism may be most relevant to L-bursts, rather than to S-bursts, which are already fully developed on this time scale.

Finally, we note that Melrose did not quantitatively address the question of the beaming pattern of his mechanism. As we shall discuss in section IV, this is an important topic for comparison of the predictions of the mechanism with the observational phenomenology.

We have discussed only mechanisms that have been treated quantitatively in the literature. Other suggestions have been advanced by Warwick [1961, 1963], Zheleznyahov [1966], and Marshall and Libby. For a discussion of these works, the reader is referred to the review by Warwick [1967].

IV. WAVE PROPAGATION IN THE JOVIAN PLASMASPHERE

Compared with the aspects of coupling and instability mechanisms, propagation of the DAM waves in the Jovian plasmasphere has received relatively little attention. It is an important topic, however, which must be considered as an integral part of a complete theory of DAM. The beaming pattern implied by the DAM morphology is undoubtedly due to a combination of the inherent beaming pattern of the instability mechanism and the subsequent propagation of the waves. Conversely, the phenomenology may provide the basis for inferences regarding the plasma medium and the source regions.

As an example of the latter approach, some authors have considered the polarization. Warwick and Dulk [1964] argued that the highly elliptical polarization of decametric waves upon emergence from the Jovian magnetosphere indicates that the base mode of the emission is the extraordinary mode, an interpretation which has generally been accepted in the literature. The sense of polarization of this mode is right (left)-handed when the wave vector has a component parallel (antiparallel) to $\vec{B}_o$. Although the limiting polarization depends upon the total ray path, the former case would generally be expected to occur for emission from the northern hemisphere, and the latter for emission from the southern hemisphere. At high frequencies (> 18 MHz), DAM is predominantly right-polarized, while at lower frequencies the degree of left-polarization increases as the frequency decreases. The axial vector exhibits a quasi-sinusoidal rotation profile, which is conventionally interpreted as indicating that at a given longitude one or the other hemisphere dominates the beaming into a given direction. This reasoning should be considered with caution, however, for within storms that are rich in time structures there may exist considerable polarization diversity in the frequency-time plane.

Notwithstanding the argument of Warwick and Dulk about the base mode of the emission, Warwick [1970] pointed out that the radiation incident upon the terrestrial ionosphere is not in a base mode. Nearly all of the Faraday rotation observed, however, is attributable to propagation through the terrestrial ionosphere; Warwick estimated that less than 360° of Faraday rotation occurs in the Jovian magnetosphere. Warwick argued that this fact placed a small upper limit on the column density $\int N_e\, ds$ through which the radiation propagates at Jupiter, and that this limit on the column density was incompatible with density models such as those of Gledhill [1967] and Goertz [1973, 1975], in which there is a disk-shaped concentration of plasma in the equatorial plane.* Goertz [1974] challenged this argument;

*Although the density models of Goertz and Gledhill are topologically similar, they are derived from different considerations and so differ by orders of magnitude in the actual densities predicted at a given radial distance in the plasma disk.

he showed that, if the radiation is generated entirely in a base mode, then during propagation through a spatially varying medium it will couple to the other mode just sufficiently to preserve its polarization, and so Faraday rotation could not occur.

As we have discussed in section III, several proposed instability mechanisms actually amplify waves on the slow x-branch; such waves cannot escape directly to free space. The general topology of the stop zones has been discussed by Liemohn [1973] and by Gulkis and Carr [1966]; these latter authors suggested that the dimension of the stop zone might be longitudinally asymmetric, so that the DAM beaming pattern--as manifested in the source morphology--might be governed by the location of the thinnest stop zones. They noted, however, that the mechanism by which waves could penetrate the stop zones was unclear.

Oya [1974] attempted to account for the escape of waves through the stop band by considering mode conversion of extraordinary waves propagating inward towards the planet. The conversion occurs by the reflection of the o-mode at the plasma level $\omega_e = \omega$. There are a number of difficulties with the geometry inherent in Oya's model. It suffices to note, however, that if the peak ionospheric electron density is 3×10^5 as reported by Fjeldbo <u>et al.</u>, then the corresponding plasma frequency is about 5 MHz, and so the critical condition is not meet at DAM frequencies.

The importance of the stop zones, however, is greatly dependent on the temperature and peak density of the Jovian ionosphere.

Smith [1975] showed that for $N_o = 3 \times 10^5$ [Fjeldbo et al., 1975] and $T_o = 750°$ [Kliore et al., 1975], the attenuation of waves during traversal of the evansecent region is negligible for frequencies below about 20 MHz. This fact has profound consequences for the explanation of Io-independent DAM, which we discuss in section V.

As we have seen in section III, theories in which instability occurs on the slow x-branch, such as those of Goldstein and Eviatar [1972] and of Wu [1973], amplify waves of very low group velocity $V_g \ll c$. Smith et al. [1972] and Wu et al. [1973] attempted to account for the Io-phase asymmetry (referred to in section II) in terms of the low group velocity of the waves before their escape from the Jovian magnetosphere. In addition, these authors suggested that the sense of frequency drift of the modulation lanes--negative in the main and third sources and positive in the early source--was caused by the vectorial combination of a low group velocity in the medium relative to the observer. This latter suggestion stemmed from an interpretation of the modulation lanes as drifts as a modulated structure in the emission itself. Smith [1973] showed that this interpretation could not be correct if the modulation structure in the emission were impressed upon the IFT; instead, a longitudinal modulation is required within the framework of the group-delay hypothesis.

Smith [1973] also found that refractive effects upon upper-hybrid waves were strong, tending to make the waves propagate into the stop zone even when the wave vector is initially perpendicular to the normal to the stop zone surface. In view of Smith's later results concerning the thickness of the stop zones, it is unlikely that group-delay effects play any significant role in either the phenomenology or morphology of DAM.

One caveat to be made about this last statement is that it implicitly assumes geometrical optics to be valid everywhere. All of the instability mechanisms discussed in section III imply either subsequent propagation through a thin evanescent region or, as in the case of Melrose's mechanism, that the radiation is intially beamed toward the planet and subsequently reflected toward the observer. All studies of wave propagation to date have been in the context of geometrical optics, in which the medium is slowly varying and the propagation at any point is governed by the Appleton-Hartree equation with the local plasma parameters. The condition for the validity of geometrical optics, which is essentially a WKB approach, is

$$\frac{c}{\omega n^2} |\nabla n| \ll 1 \quad . \tag{IV-1}$$

Near resonance and reflection, inequality (IV-1) breaks down. Nevertheless, application of geometrical optics in the case of reflection, for example, usually yields good results for the ray trajectory even when the solution has been carried through a thin region where WKB analysis is invalid [Ginzburg, 1970]. Strictly speaking, however, the breakdown of inequality (IV-1) in some region implies that a full-wave solution must be carried out in that region and matched to the WKB solution in regions where geometrical optics is valid. It is possible that such an analysis would explain some of the fine structure observe in decametric spectra.

V. SUMMARY AND DISCUSSION

For the most part the discussion above has treated coupling, instability, and propagation separately, although occasional references have been made to the relations between them. In this section we attempt to draw the different aspects together in evaluating the present state of the theory. We shall also indicate some of the major unsolved problems, and suggest some possible approaches for future work.

A. Io-Independent DAM

As we have already remarked, there are not one but three distinct decametric emissions: Io-independent L-bursts and Io-modulated L- and S-bursts. Since the discovery of the Io-modulation only one paper, that of Goldstein and Eviatar, has been addressed specifically to the question of an emission mechanism that cannot be related to a coupling mechanism. This mechanism appears to be viable for the Io-independent DAM, especially in light of the indications that tunneling of the radiation through the evanescent regions does not lead to severe attenuation. The radiation may emanate from particles trapped on L-shells close to but inside the orbit of Io, and so may be unaffected by any magnetospheric perturbations local to Io. The waves are inherently beamed sharply perpendicular to the magnetic field, in accord with the declination effect. The chief shortcoming of the work of Goldstein and

Eviatar is that they did not estimate the level at which the instability would saturate due to pitch-angle scattering by the radiation.

Another theory which may account for the Io-independent DAM is the sweeping model of Wu. Although this model was originally proposed to explain Io-modulated emission, there are several reasons why it may be more relevant to the unmodulated radiation. Among these are the sharp beaming pattern perpendicular to the magnetic field, and the fact that, as discussed in section III, the instability is limited to frequencies below 30 MHz within the context of the density inferred by Fjeldbo et al. In addition, the growth rates evaluated by Wu and by Smith [1973] are rather low, and cannot account for the time scales of Io-modulated L-bursts.

Although sweeping is certainly a coupling mechanism, the real question is whether its effects are localized close to Io. Smith [1973] showed that in addition to the upper-hybrid instability, the drift currents induced by sweeping also drive a long-wavelength instability at the lower hybrid frequency $\omega_{UH} = \Omega_i[1 + \epsilon(M/m)/(1+\epsilon)]^{\frac{1}{2}}$. . Smith evaluated the spatial diffusion coefficient [Dupree, 1967]

$$D_{\perp} = \max\left[\frac{\gamma(k)}{k^2}\right] \quad ,$$

and showed that it depended on both the thermal plasma density and the local gyrofrequency, in addition to its dependence upon the energetic-proton distribution. In general, the diffusion is slower at high altitudes, so that the density gradient is maintained longer at low frequencies. Smith suggested that the Io-independent DAM is in fact due to sweeping, but is poorly correlated to the position of Io because of the slow diffusion at low frequencies, and thus appears statistically independent of Io. Another diffusion mechanism has recently been investigated by Huba and Wu [1975].

Wu [1973] and Smith [1973] neglected the longitudinal drift motion of the protons in the dynamics of the sweeping process. This drift is oppositely directed to Io's orbital motion relative to the corotating plasma. As shown by Thomsen and Goertz [1975], the swept cavity drifts in the same way. Thus, the effective sweeping time is reduced from 13 hours to about 4 hours, and it is likely that the cavity persists along the entire Io orbit. Although the diffusion tends to reduce the magnitude of the density gradient during this period, the gradient is never completely eliminated, and so the drift-current-driven instability is always occurring.

A possible impediment to the sweeping model as the cause of unmodulated DAM might exist if in fact Io causes local heating of the ionosphere by MHD waves, as suggested in section II. Such heating would drive thermal plasma up into the IFT, increasing the scale height H of the thermal-plasma density gradient. If $H \gg \epsilon R_J$, then the thickness of the stop zone becomes equal to H, so that x-mode waves could not penetrate the stop zone. We might expect this enhanced density "filament" to cool quickly by thermal conduction to the ionosphere, however, and so be localized close to Io.

B. Io-Modulated L-Bursts

A complete theory for Io-modulated DAM would synthesize all three aspects of the problem. Starting from a particular demonstrated coupling mechanism, the theory would predict the form of the resulting free-energy source, for example the velocity-space distribution of accelerated particles. It would then incorporate this distribution into an instability analysis, demonstrating that the instability could account for the various aspects of DAM phenomenology such as time scales, bandwidth, polarization, and spectral power. The inherent beaming pattern of the radiation in the source region would be evaluated, and its subsequent propagation in and emergence from the Jovian plasmasphere would be compared with DAM morphology. In addition, the theory should contain some natural explanation for the probabilistic nature of the emission.

Needless to say, such a comprehensive theory has not yet been attained. Nevertheless, some promising work has been done on different aspects of the problem. Smith and Goertz have given the sheath model of acceleration a firm theoretical basis. The emission mechanism of Melrose is compatible with the sheath coupling model, and seems able to account for most of the necessary features of the phenomenology. It yet remains, however, to synthesize these two models, and to investigate the propagation and reflection of the resulting radiation.

Continued investigations of other mechanisms for both coupling and instability might also yield viable new models. For example, Palmadesso _et_ _al_. [1975] have recently proposed a parametric instability mechanism for the terrestrial kilometric radiation. This theory seems attractive, and might also have application in the parameter regime of DAM.

There can be no question that propagation effects contribute significantly to the DAM morphology, and probably also to the phenomenology. Any instability mechanism has an inherent beaming pattern in the source region. The structures of the magnetic field and the exospheric density distribution combine to modify this pattern into the directivity pattern that is conventionally inferred from the morphology. A complete theory of DAM must include a study of the "antenna pattern" of the Jovian plasmasphere to show that it is compatible with this morphology. Such a study would serve two purposes. First, to the extent that it is dependent upon a model of the exospheric structure, it would help to elucidate that structure. Second, the emerging radiation pattern depends upon the beaming pattern of the instability in the source region. Thus, study of the propagation may place boundary conditions upon the instability mechanism which allow one to discriminate between likely mechanisms. Because the instability mechanism is, in general, also dependent upon the ambient plasma regime, we might expect that the synthetic study of coupling, instability, and wave propagation will lead to a convergent picture.

C. Three Major Problems

None of the models we have discussed appear able to account for the Io-modulated S-bursts. Ellis [1965] interpreted them as being due to streams of 30-60 keV electrons moving outward. It is unclear why such streams should always be accelerated outward, or why they should occur primarily in the early-source region. It is also unclear why they should give rise to a different instability than that responsible for the L-bursts. As noted in section III, the well-separated time scales of the L- and S-bursts, and the time frequency drifts of the S-bursts, indicate that different mechanisms produce the decasecond and millisecond emissions. For the moment the S-bursts remain enigmatic.

A second major question is the explanation of the asymmetry in Io phase for emission viewed from east and west of the Earth-Jupiter line. This may be a local-time propagation effect; if so, its explanation would have important implications about the structure of the Jovian plasmasphere.

Finally, one of the most intriguing problems regarding DAM is the explanation of the modulation lanes observed in both L- and S-bursts. Riihimaa [1971, 1974] has documented an extensive morphology for these features. The modulation lanes probably contain a great deal of information about the mechansim and propagation of DAM.

Epilogue

We began this review with the observation that a comprehensive understanding of DAM has yet to be attained. Nevertheless, many authors have contributed a great deal of insight to our perspective of the problem. The conceptual framework for the theory is largely in place, and we may reasonably expect much progress in the future.

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APPENDIX A
HEATING OF THE JOVIÂN IONOSPHERE BY MHD WAVES

In section II-A we mentioned the possibility that MHD waves generated by Io might provide local heating of the Jovian ionosphere. This comment was based on the following rough but indicative calculation, due to Wu and Smith [unpublished].

Assume the waves steepen into shocks during their propagation along the IFT. The typical dimension of a collisionless shock is c/ω_i, where $\omega_i = 2\pi f_{pi}$ is the angular ion plasma frequency. Then for a wave amplitude of δB, the current density j is given by

$$j = \frac{c}{4\pi} |\nabla \times \vec{B}| \simeq \frac{\omega_i \delta B}{4\pi} \quad .$$

The Joule heating rate of the plasma is then

$$NK_B \frac{\partial T}{\partial t} = \frac{j^2}{\sigma} = \frac{\omega_i^2 \delta B^2}{16\pi^2 \sigma} \quad . \tag{A-1}$$

The conductivity σ can be either anomalous or the usual Coulomb conductivity

$$\sigma_c \simeq \frac{K_B T^{\frac{3}{2}}}{m_e^{\frac{1}{2}} e^2 \ln \Lambda} \quad , \tag{A-2}$$

where $\Lambda \equiv (K_B T)^{\frac{3}{2}} / (4\pi N)^{\frac{1}{2}} e^3$. Use of σ_c will lead to an upper bound for the heating. Substituting (A-2) with (A-1), we obtain

$$\frac{\partial}{\partial t} T^{5/2} = \frac{5}{32} \frac{m_e^{\frac{1}{2}} e^2 \ln \Lambda \, \omega_i^2 \, \delta B^2}{\pi^2 N K_B^{5/2}} \quad . \tag{A-3}$$

Because the temperature dependence on the RHS (A-3) is logarithmic, we may obtain an approximate solution by taking $\ln\Lambda$ to be constant; we assume $\ln\Lambda \simeq 20$. Assuming that all the heating occurs in the ionosphere, in a time interval τ, we find

$$T(\tau) \simeq T_o \left[1 + \frac{1.7 \times 10^{13}}{T_o^{5/2}} \delta B^2 \tau \right]^{2/5} \quad . \tag{A-4}$$

Taking $T_o \simeq 750$ K [Kliore et al., 1975] and $\tau \simeq 60$ s (the time for which the IFT threads Io), (A-4) becomes

$$T(\tau = 60 \text{ sec}) \simeq 10^6 \, \delta B^{4/5} \quad . \tag{A-5}$$

We may estimate δB in terms of the amplitude δB_o of the perturbation by Io. For a rough estimate, we consider a generalized Poynting theorem. Let U be the total energy density (electric, magnetic, and kinetic) in the wave, and let $\vec{V}_g$ denote the group velocity. The generalized Poynting theorem is:

$$\frac{\partial U}{\partial t} + \vec{V}_g \cdot \nabla U + U \nabla \cdot \vec{V}_g = 0 \quad . \tag{A-6}$$

We assume that a WKB solution is possible. (Strictly speaking, this is not consistent with our estimate of $|\nabla \times \vec{B}|$ from the shock condition; we discuss this below.) The Alfvén wave dispersion relation is

$$\omega = \frac{\vec{k} \cdot \vec{V}_A}{\left(1 + V_A^2/c^2\right)^{\frac{1}{2}}} \quad ;$$

therefore,

$$\vec{V}_g = \alpha \vec{B} \quad ,$$

where $\alpha \equiv [4\pi N m_p (1 + V_A^2/c^2)]^{-\frac{1}{2}}$, and so

$$\nabla \cdot \vec{V}_g = B(s) \frac{\partial}{\partial s} \alpha(s) \equiv f(s) \quad ,$$

where $\partial\alpha/\partial s$ is the directional derivative along $\vec{B}$. Assuming a solution of the form

$$U = W \exp - \int_0^s dx \frac{f(x)}{V_g(x)} \quad , \tag{A-7}$$

Eq. (A-6) transforms to

$$\frac{\partial W}{\partial t} + V_g \frac{\partial W}{\partial s} = 0 \quad , \tag{A-8}$$

the solution of which is simply

$$W(s, t) = W\left[s', t - \int_{s'}^{s} \frac{dx}{V_g(x)}\right] \quad .$$

The boundary condition on the solution is

$$W(0, t) = U_{Io}(t) \quad .$$

Furthermore, the integrand in (A-7) is simply

$$\frac{f(s)}{V_g(s)} = \frac{\partial}{\partial s} \ln \alpha \quad ,$$

so that

$$U(s) = U_{Io} \frac{\alpha_{Io}}{\alpha(s)} = U_{Io} \left[\frac{N(s)}{N_{Io}} \frac{1 + V_{AIo}^2/c^2}{1 + V_A^2(s)/c^2}\right]^{\frac{1}{2}} \quad . \tag{A-9}$$

In an Alfvén wave, energy is partitioned in the ratios

$$\frac{E^2}{8\pi} : \frac{B^2}{8\pi} : \frac{1}{2}\,\rho V^2 = \frac{V_A^2}{c^2} : 1 : 1 \quad .$$

Thus, neglecting V_A^2/c^2 throughout, we may combine (A-5) and (A-9) to find

$$T(\tau) \simeq 10^6 \left(\frac{N_o}{N_{Io}}\right)^{2/5} (\delta B_{Io})^{4/5} \quad ,$$

where N_o is the ionospheric electron density. Assuming $N_{Io} \simeq 30\ \mathrm{cm}^{-3}$ and $N_o \simeq 3 \times 10^5\ \mathrm{cm}^{-3}$, we have $T(\tau) \simeq 4 \times 10^7\ (\delta B_{Io})^{4/5}$. The power radiated by Io in Alfvén waves is approximately

$$P \simeq V_{A,Io}\ \pi R_{Io}^2 \left(\frac{\delta B_{Io}^2}{8\pi}\right) \quad .$$

Taking $V_{A,Io} \simeq 7 \times 10^3$ km/sec, we require $\delta B_{Io} \simeq 2 \times 10^{-5}$ g if $P = 10^{10}$ W, and 2×10^{-3} g if $P = 10^{14}$ W. (The former figure is Schmahl's estimate of the power; the latter is the upper limit for which the effect on Io's orbital period would not be observable.) The corresponding temperature, according to (A-10) are 7×10^3 K and 3×10^5 K.

The estimate (A-10) is quite crude, however, because the arguments leading to it were not entirely consistent. On the one hand, we assumed that the waves steepen into shocks in order to estimate the Joule heating current. On the other hand, we assumed that the wave propagation was compatible with a WKB solution in order to find the steepened amplitude. The effect of the first assumption is to overestimate the numerical coefficient in Eq. (A-10), the effect of the second is to underestimate δB. The value of the estimate should not be taken too seriously; rather, the calculation should be viewed in the spirit that the large discrepancy between the temperature as given by (A-10) and the nominal ionospheric temperature of 750 K indicates that transient local heating by Io may be significant.

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FIGURE CAPTIONS

Figure 1. The density of protons in the energy range 108 eV—4.80 keV near Io, as measured by Frank _et al._ [1975]. These authors fit their proton density measurements throughout the inner plasmasphere to Maxwellian distributions of characteristic temperature $\simeq$ 100 Ev and densities between 50 and 100 cm^{-3}.

Figure 2. The geometrical configuration of the Earth-Jupiter-Io system for peak occurrence probability in the various sources. Note the asymmetry in Io phase with respect to the Earth-Jupiter line. The CML ranges of the sources are in System III (1957.0) coordinates, at the Pioneer 11 epoch (1974.9).

Figure 3. (After Goertz and Deift [1973].)

(a) The magnetic field topology near Io.

(b) Detail of the x-type neutral point reconnection region behind Io.

Figure 4. Schematic of the Io sheath model [after Shawhan _et al._ [1974]).

Figure 5. The idealized circuit analyzed by Gurnett [1972]. The inset shows the plasma flow velocity in the context of Gurnett's analysis. Σ_ρ is the height-integrated Pedersen conductivity of the Jovian ionosphere. For Σ_ρ less than a

critical value $\Sigma_c \simeq 5$ mho, the electric field in the ionosphere is large and the flux tube tends to corotate with Io. For $\Sigma_\rho > \Sigma_c$, the voltage drop across the ionosphere is small and the plasma slips by Io.

Figure 6. Cross-section of the energetic-particle cavity in Io's orbital plane.

Figure 7. Mapping of the density-gradient layer along the Io flux tube. Also shown are the directions of the density gradients and drift currents at the inner and outer walls of the cavity.

Figure 8. Dispersion curves and indices of refraction of the ordinary and extraordinary modes for $\epsilon \ll 1$. The approximate regions of instability found by various authors are indicated.

Figure 9. The gyrofrequency at the feet of the Io flux tube, in the magnetic field model of Smith et al. [1975]. The approximate sub-Io longitude ranges of the various sources are indicated (epoch 1974.9).

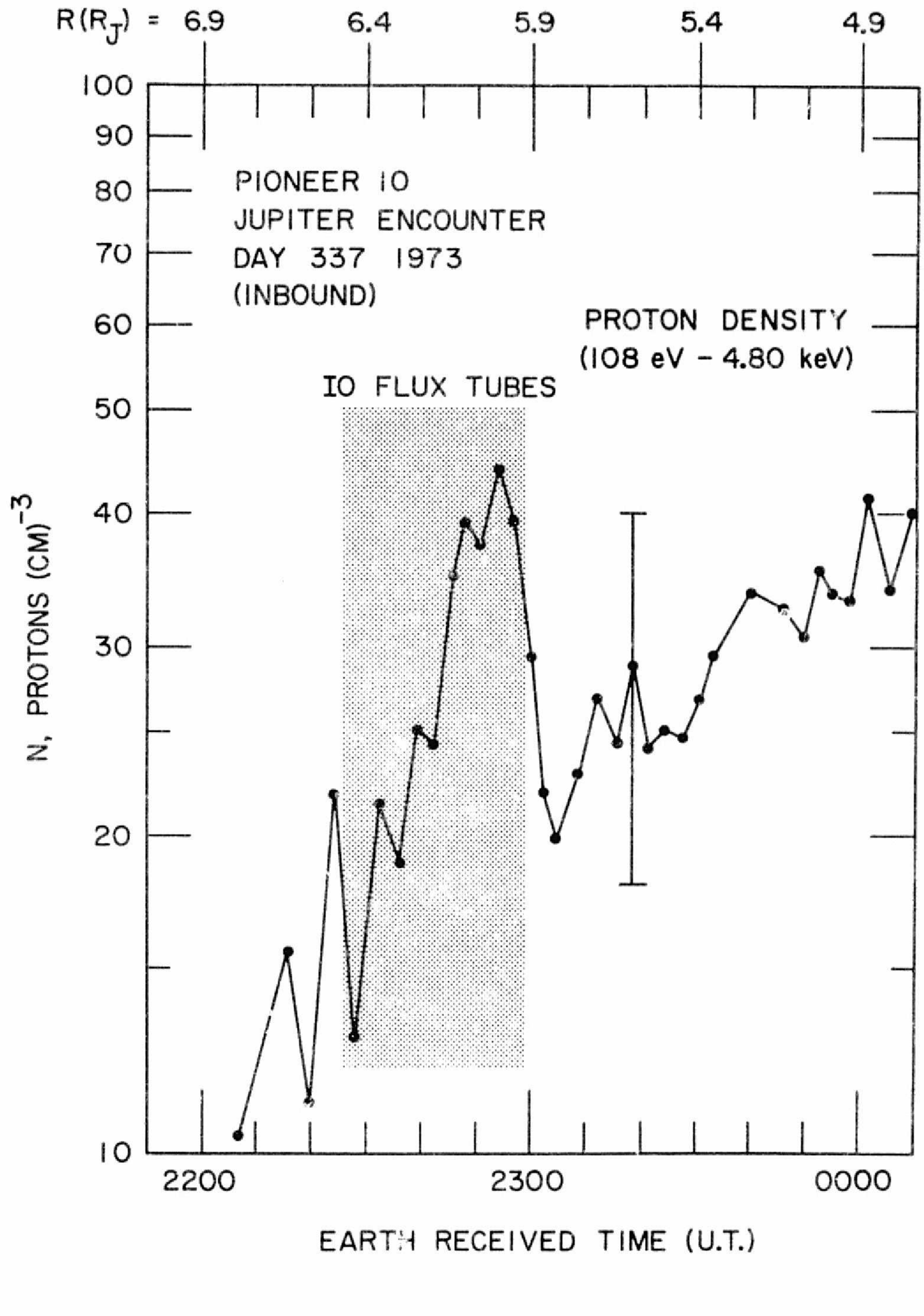

R(R_J) = 6.9
6.4
5.9
5.4
4.9
100
90
80
70
60
50
40
30
20
10
PIONEER 10
JUPITER ENCOUNTER
DAY 337 1973
(INBOUND)
PROTON DENSITY
(108 eV - 4.80 keV)
IO FLUX TUBES
N, PROTONS (CM)^-3
2200
2300
0000
EARTH RECEIVED TIME (U.T.)

Figure 1

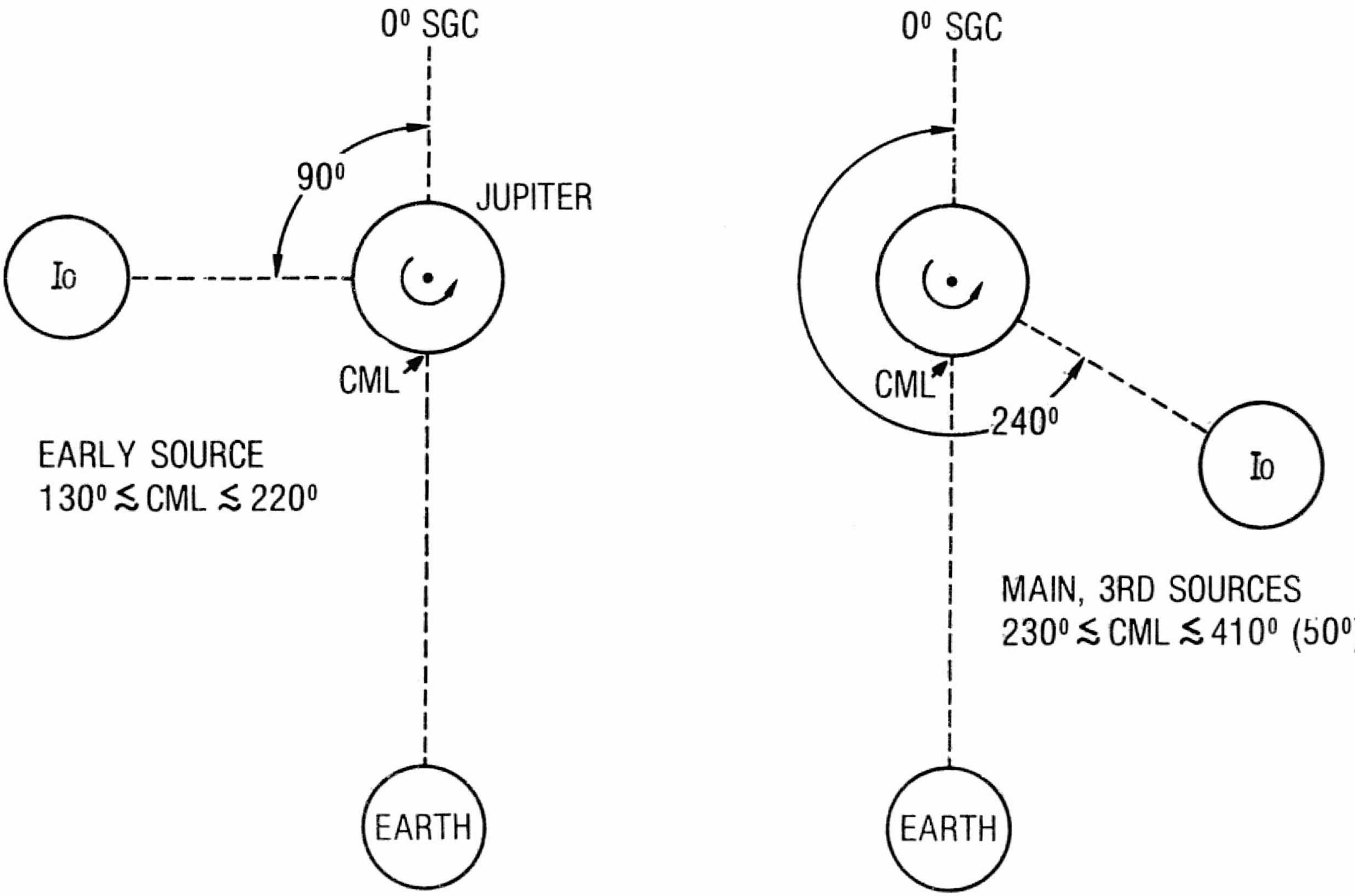

0⁰ SGC
90⁰
JUPITER
Io
CML
EARLY SOURCE
130⁰ ≲ CML ≲ 220⁰
EARTH
0⁰ SGC
CML
240⁰
Io
MAIN, 3RD SOURCES
230⁰ ≲ CML ≲ 410⁰ (50⁰)
EARTH

Figure 2

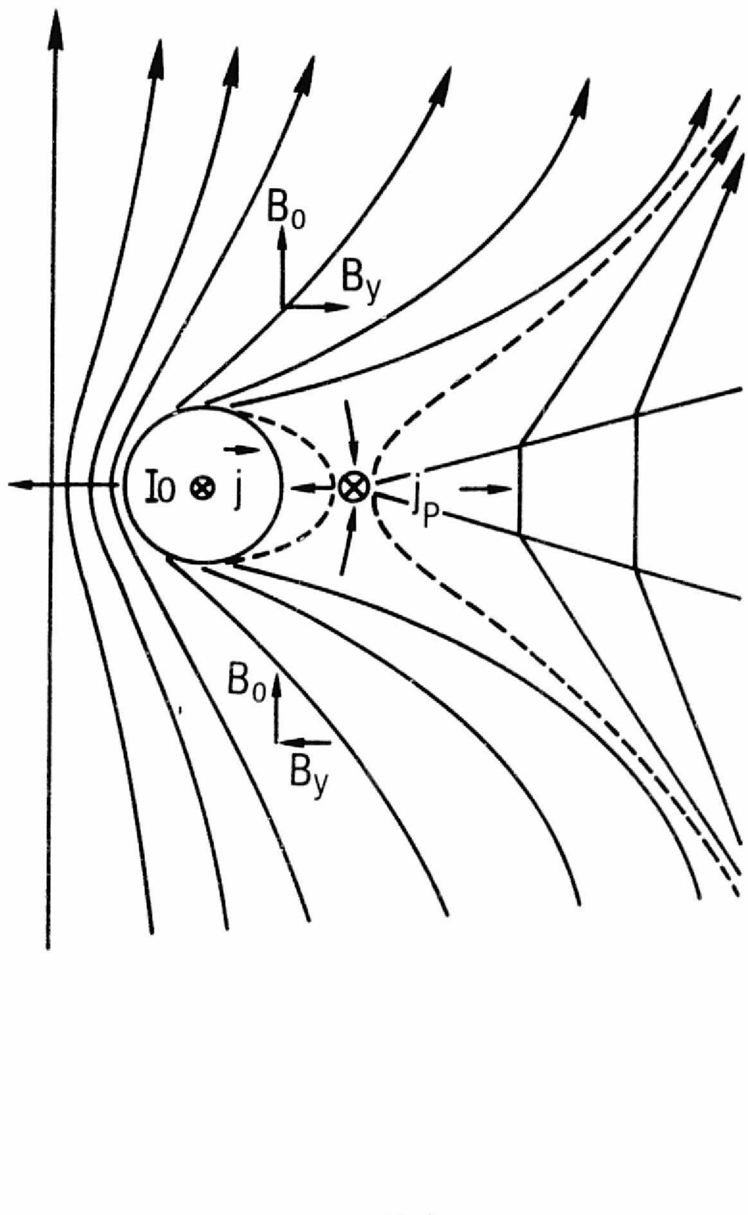


Figure 3(a)

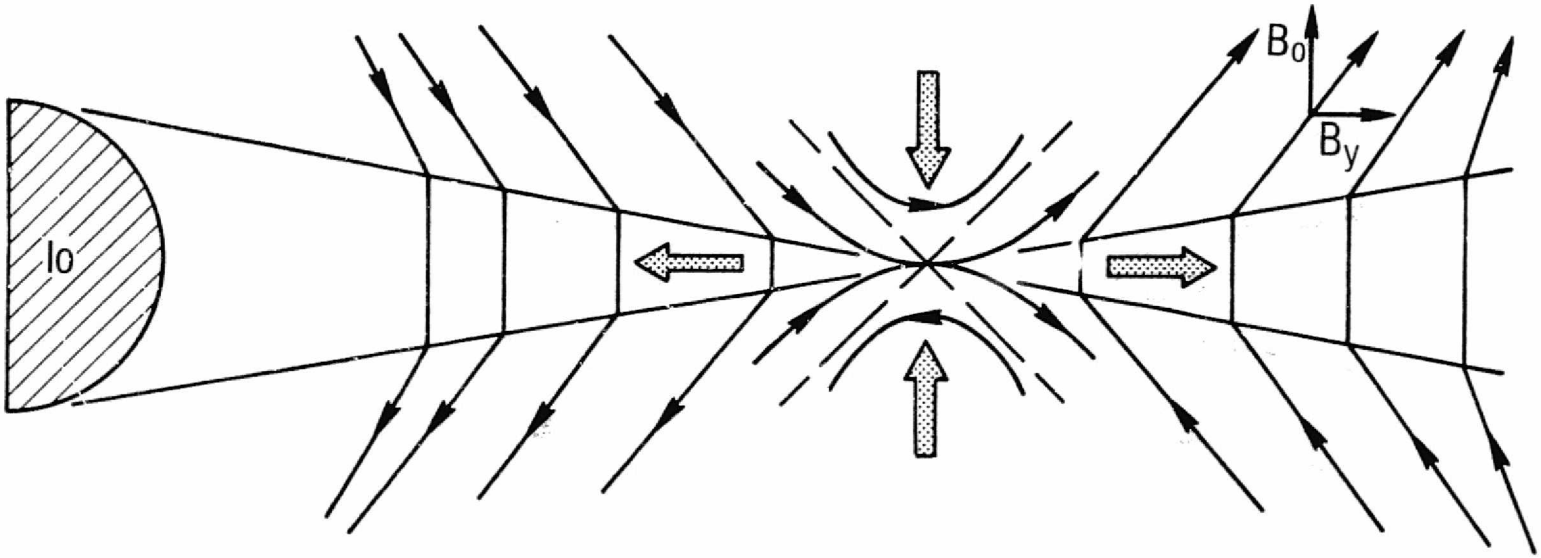
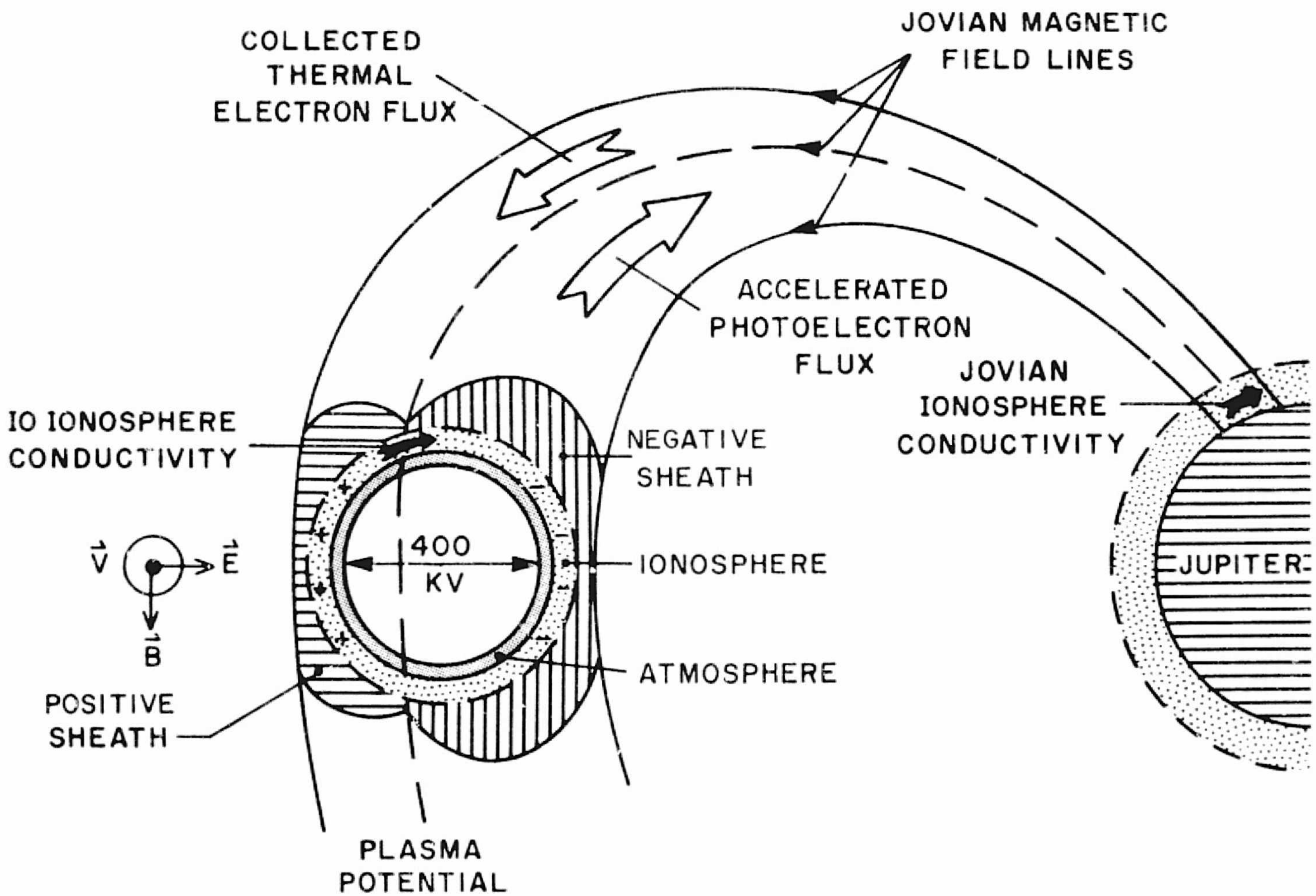
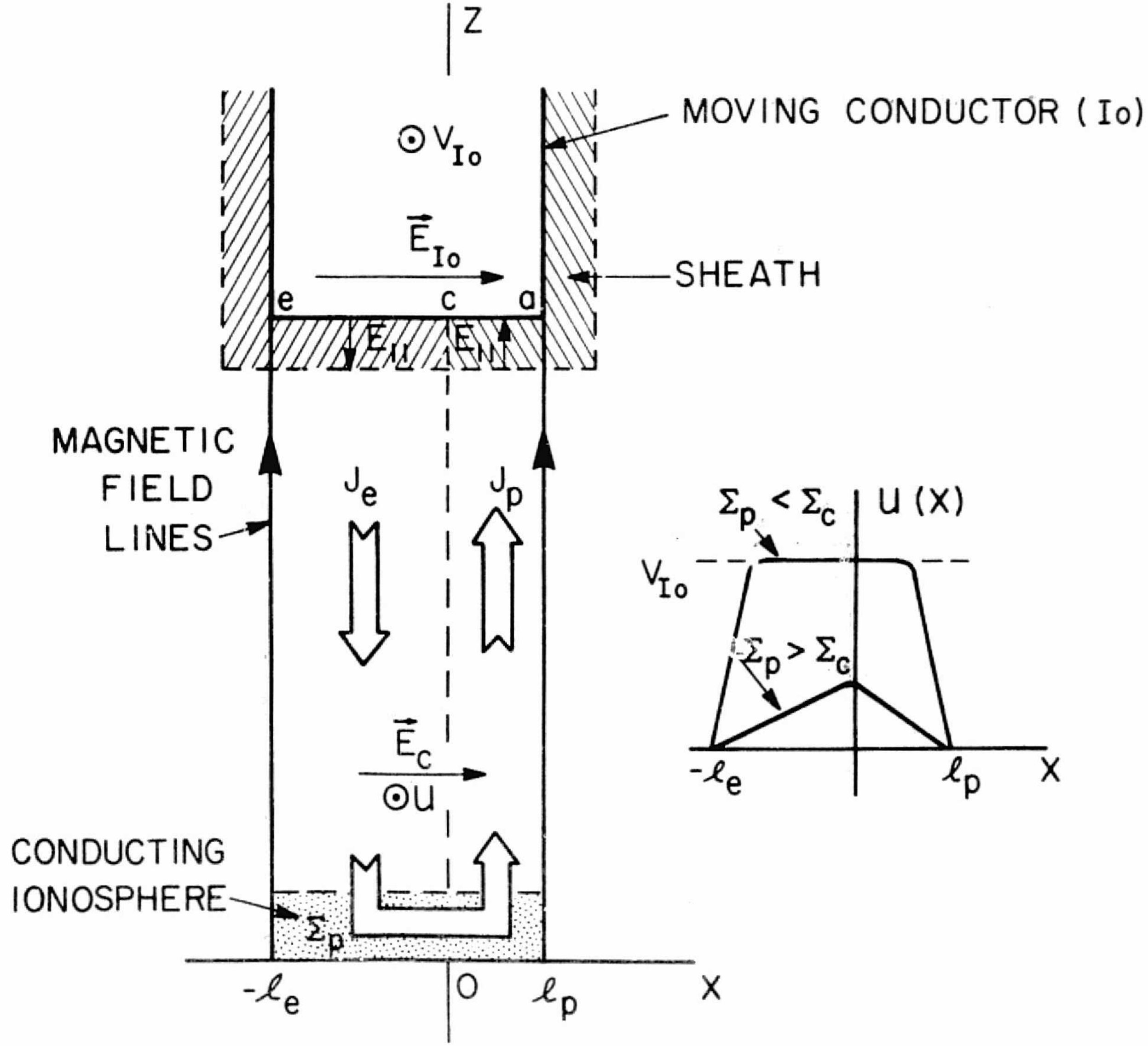


Figure 3(b)



COLLECTED THERMAL ELECTRON FLUX
JOVIAN MAGNETIC FIELD LINES
ACCELERATED PHOTOELECTRON FLUX
JOVIAN IONOSPHERE CONDUCTIVITY
IO IONOSPHERE CONDUCTIVITY
NEGATIVE SHEATH
400 KV
IONOSPHERE
JUPITER
ATMOSPHERE
POSITIVE SHEATH
PLASMA POTENTIAL

Figure 4



Z
MOVING CONDUCTOR (Io)
$\odot V_{Io}$
$\vec{E}_{Io}$
SHEATH
e
c
a
$E_{\parallel}$
$E_{\parallel}$
MAGNETIC
FIELD
LINES
J_e
J_p
$\vec{E}_c$
$\odot u$
CONDUCTING
IONOSPHERE
Σ_p
$-\ell_e$
O
ℓ_p
X
$\Sigma_p < \Sigma_c$
U (X)
V_{Io}
$\Sigma_p > \Sigma_c$
$-\ell_e$
ℓ_p
X

Figure 5

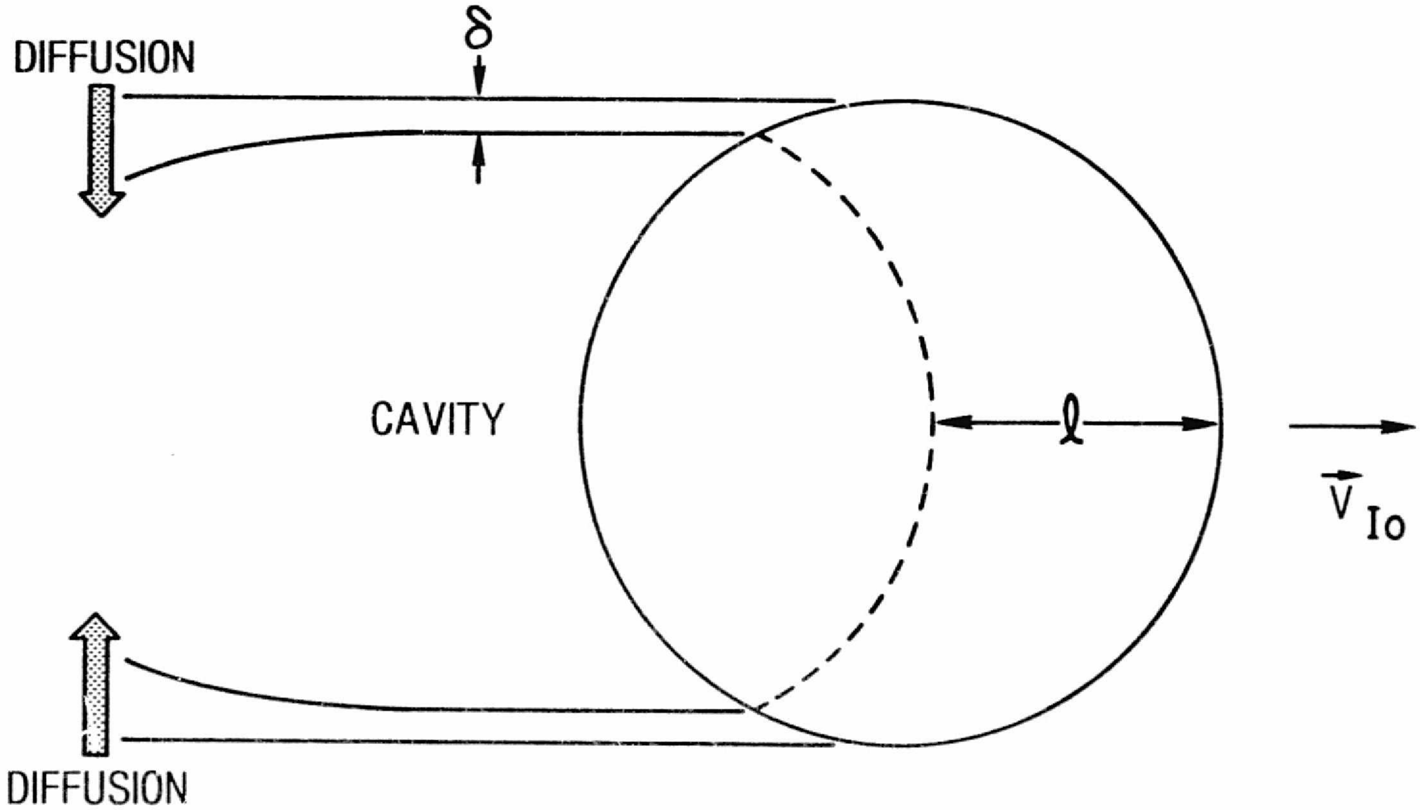


Figure 6

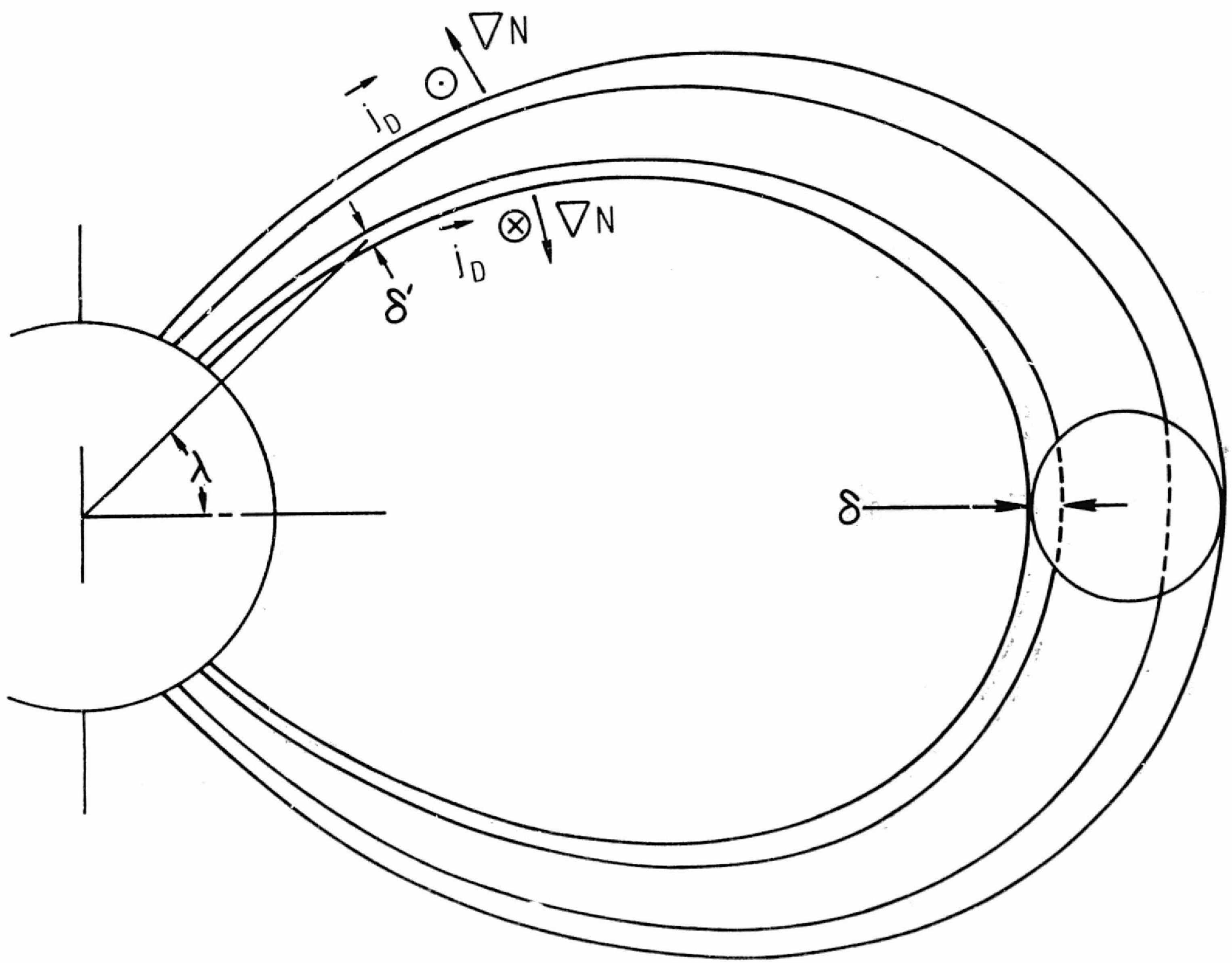
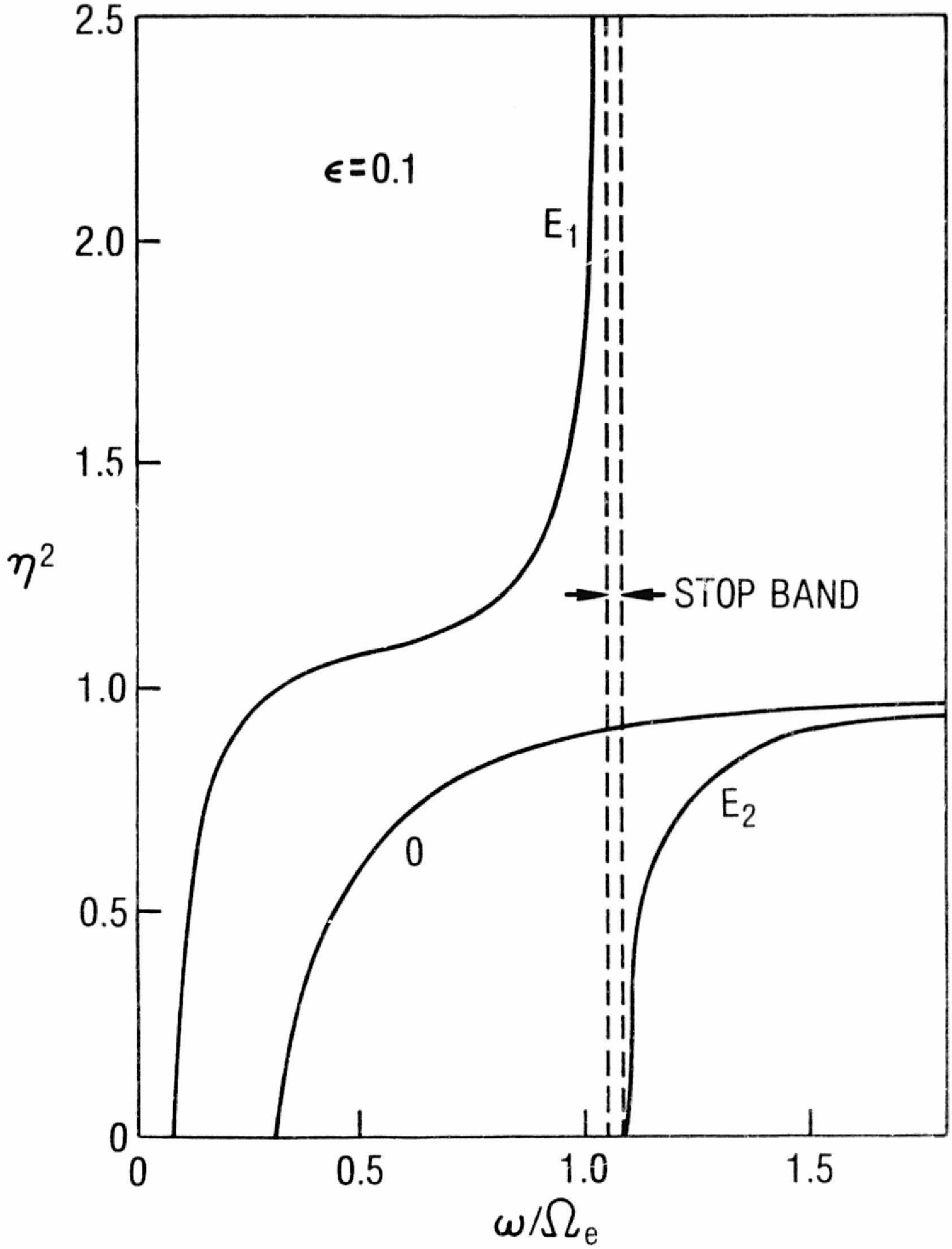
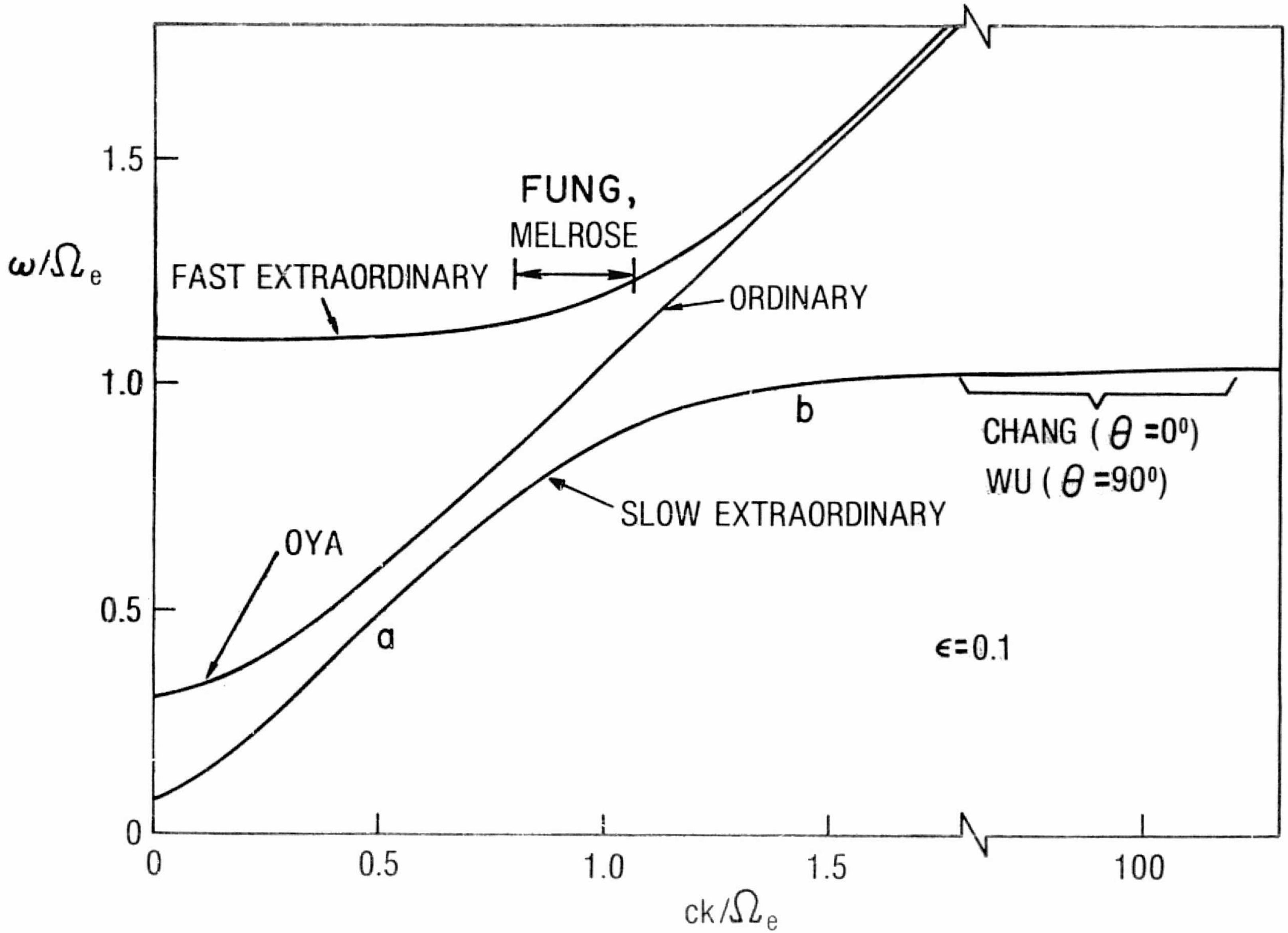


Figure 7



2.5
2.0
1.5
1.0
0.5
0
η^2
ϵ=0.1
E$_1$
STOP BAND
0
E$_2$
0
0.5
1.0
1.5
ω/Ω_e

Figure 8(a)



FUNG,
MELROSE
FAST EXTRAORDINARY
ORDINARY
ω/Ω_e
1.5
1.0
0.5
0
b
CHANG (θ =0°)
WU (θ =90°)
SLOW EXTRAORDINARY
OYA
a
ϵ=0.1
0
0.5
1.0
1.5
100
ck/Ω_e

Figure 8(b)

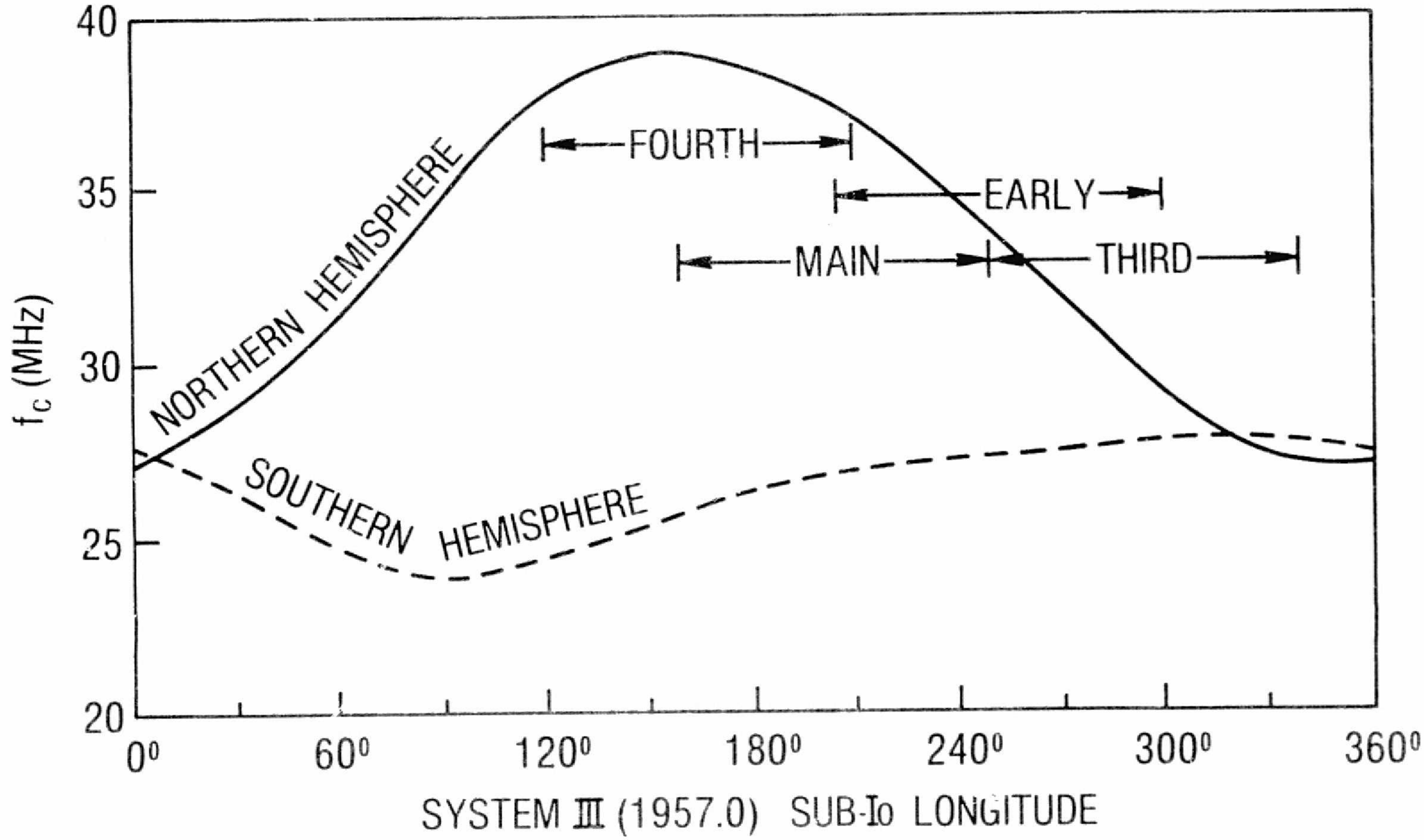

40
35
30
25
20
f_c (MHz)
NORTHERN HEMISPHERE
SOUTHERN HEMISPHERE
FOURTH
EARLY
MAIN
THIRD
0°
60°
120°
180°
240°
300°
360°
SYSTEM III (1957.0) SUB-Io LONGITUDE

Figure 9