

Bericht der Herbsttagung der Schweizerischen Physikalischen Gesellschaft vom 8./9. Oktober 1976 in Genf

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Bericht der Herbsttagung der Schweizerischen Physikalischen Gesellschaft vom 8./9. Oktober 1976 in Genf

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1. Mitglieder

Die folgenden Neumitglieder sind in unsere Gesellschaft aufgenommen worden: B. ADAM (Zürich), M. BERSSET (Lausanne), D. BESSET (Genf), W. BOLLI (Eglisau), P. BONNARD (Lausanne), M. BUETTIKER (Windisch), M. COHN (Zürich), M. DACAROGNA (Genf), R. DELESCLEFS (Genf), J. N. IMFELD (Lausanne), J. JOLLER (Bochum), H. KAUSCH (Lausanne), CH. PFERSICH (La Tour de Peilz), F. ROESEL (Basel), CH. ROSSEL (Genf), R. SCHALLER (Lausanne), E. STUERMER (Uetendorf), A. TRABER (Arlesheim), P. ZWEIACKER (Lausanne).

Die Gesellschaft besteht zur Zeit aus 1160 Einzelmitgliedern, 7 Kollektivmitgliedern und 13 Bibliotheken.

2. Generalversammlung

Prof. P. Erdős, Universität Lausanne, wird als Sekretär der SPG für Frühling 1977–1979 gewählt.

Die Fachgruppe Kern- und Teilchenphysik der SPG will ausnahmsweise statt mit der SPG gemeinsam mit den deutschen und holländischen Fachgruppen vom 21. bis 25. März 1977 in Konstanz tagen.

Die Deutsche Physikalische Gesellschaft wird 1977 nicht gemeinsam mit der Französischen Physikalischen Gesellschaft tagen. Damit wird eine Beteiligung der SPG hinfällig.

Die Oesterreichische Physikalische Gesellschaft (OePG) hat die SPG freundlich eingeladen, gemeinsam im Herbst 1978 in Innsbruck zu tagen. Die Generalversammlung nimmt diese Einladung an.

3. Nächste Tagung

Die Frühjahrstagung 1977 findet am 14./15. April in Lausanne statt.

Wissenschaftlicher Teil

Beiträge aus dem Gebiet der angewandten Physik erscheinen in der Zeitschrift für angewandte Mathematik und Physik (ZAMP) 28/2, 1977.

PHYSIK DER KONDENSIERTEN MATERIE

Metalle

Etude du pic de Bordoni de L'Ag et de L'Al de haute pureté, après déformation plastique et après recristallisation¹⁾

P. STADELMANN et C. BONJOUR

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Le modèle de la formation thermiquement activée des doubles décrochements donne une explication cohérente des propriétés primaires et secondaires du pic de Bordoni des métaux cfc [1, 2]. Cependant il a été observé dans des échantillons d'Ag et d'Al de pureté 99,999% que le pic de Bordoni ne disparaît pas pendant la recristallisation, mais qu'au contraire l'intensité de la relaxation croît au cours de la recristallisation. Après la recristallisation cette intensité est 6 fois plus grande qu'après déformation plastique et est très sensible à l'amplitude de la déformation de mesure [3]. Ces observations sont interprétées à l'aide de la condition de Paré [4] qui lie les contraintes internes (ou appliquées) à l'énergie de formation d'un double décrochement. Dans ces conditions la déformation plastique n'est plus nécessaire à l'observation de la relaxation de Bordoni, mais doit être remplacée par une contrainte agissant sur la dislocation. Cette contrainte, nécessaire à l'apparition de la relaxation de Bordoni, peut ainsi être d'origine interne (donc due à la déformation plastique) ou d'origine externe (donc due à la contrainte de mesure). C'est une nécessité expérimentale.

Références

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- [4] V. K. PARÉ, *J. Appl. Phys.* 3, 332 (1961).

¹⁾ Recherche partiellement subventionnée par le Fonds National de la Recherche Scientifique, subside no. 2.486-o. 75

Mesures de trainage anélastique dans un alliage Al-Ag 30%¹⁾

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La présence des précipités γ' (Ag₂Al-hc) dans un alliage Al-Ag 30% poids, se traduit dans les mesures de frottement intérieur par un pic de relaxation anélastique à une température de 160°C environ, pour une fréquence de vibration voisine de

¹⁾ Recherche partiellement subventionnée par le Fonds National de la Recherche Scientifique, subside no. 2.486-0.75.

1 Hz. Ce pic n'apparaît pas dans les échantillons trempés, soit qu'il est effectivement absent, soit qu'il est masqué par le fond important qui se développe dans ce domaine de température, où les transformations de l'alliage perturbent fortement la mesure.

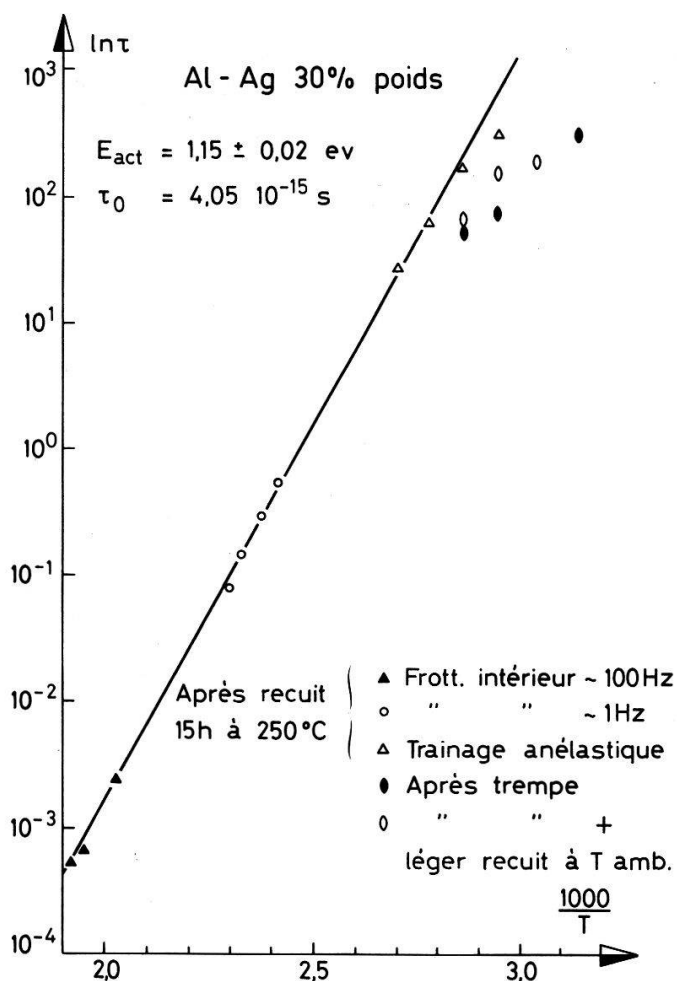


Figure 1

Diagramme d'Arrhénius du pic de relaxation observé en frottement intérieur et en trainage anélastique dans un alliage Al-Ag 30%.

Aussi, en vue d'étudier le pic précité à des températures inférieures où les évolutions internes de l'alliage sont minimisées, et d'en vérifier les conditions d'apparition, des mesures de trainage anélastique ont été entreprises.

Les résultats obtenus sont reportés sur la figure ci-jointe. Pour les échantillons recuits (15 h à 250°C), les temps de relaxation τ mesurés en frottement intérieur à basses et hautes fréquences, et en trainage, obéissent à la relation d'Arrhénius, $\tau = \tau_0 \exp(-E/kT)$, avec $E = 1,15 \pm 0,02$ eV et $\tau_0 \cong 4,05 \cdot 10^{-15}$ s, valeurs en bon accord avec les mesures d'autres auteurs [1] et [2].

Les échantillons trempés présentent aussi une relaxation, mais dont les caractéristiques sont différentes. En effet, les temps de relaxation τ sont plus petits, et même s'ils augmentent au cours de recuits à des températures voisines de l'ambiante, ce qui peut s'expliquer par un effet des lacunes en excès provenant de la trempe, ils n'atteignent pas les valeurs des échantillons recuits.

En conclusion, un phénomène de relaxation est observé dans les échantillons trempés, et qui semble être différent de celui qui se manifeste dans le pic de frottement intérieur associé à la présence des précipités γ' .

Références

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Microdéformation associée à la création de doubles décrochements dans l'or¹⁾

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Le pic de Bordoni observé à basses températures dans les métaux cfc est généralement attribué à la formation activée thermiquement des doubles décrochements sur les dislocations [1]. On en déduit alors une valeur de la contrainte de Peierls qui est de l'ordre de $5 \cdot 10^{-4} \mu$, où μ est le module de glissement. Ces résultats semblent être en contradiction avec les mesures usuelles de la limite élastique (définie à 0,2% de déformation permanente) qui, même à basse température, est bien inférieure à cette valeur.

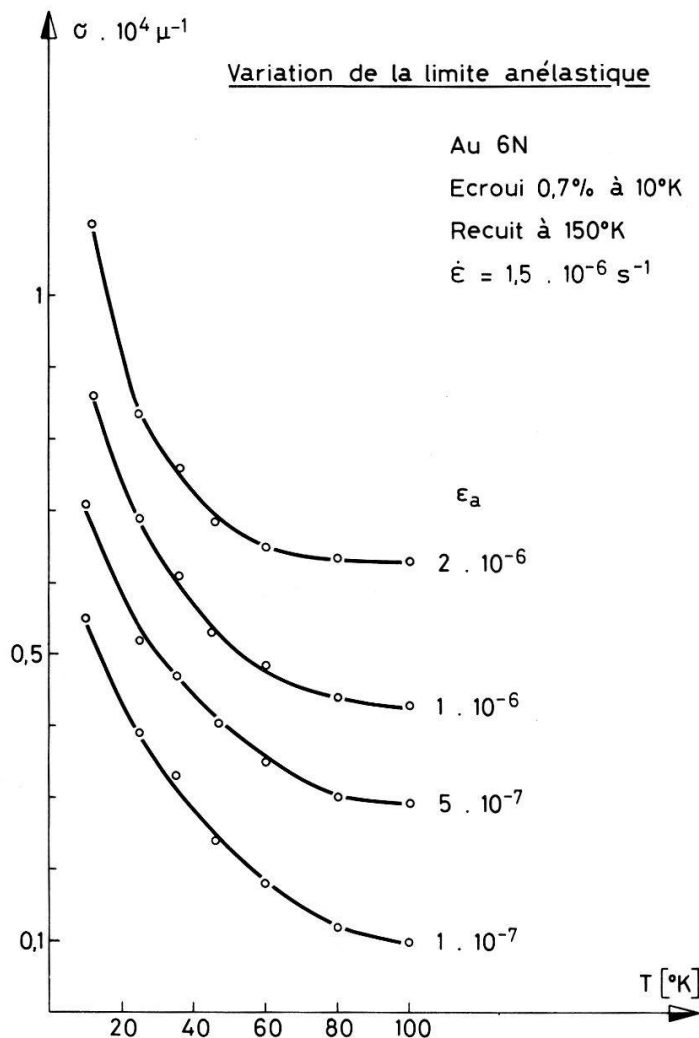


Figure 1

Variation de la limite anélastique σ en fonction de la température, pour différentes valeurs de la déformation anélastique ϵ_a .

Dans ce travail on étudie la microdéformation associée au pic de Bordoni, c'est-à-dire dans le même domaine de température (10 à 100 K). Les échantillons, sous forme de plaquettes $40 \times 5 \times 1 \text{ mm}^3$ sont soumis à une contrainte de torsion σ linéairement croissante dans le temps: $\dot{\sigma} = 1,5 \cdot 10^{-5} \mu \text{ (s}^{-1}\text{)}$. La contrainte σ , appelée limite anélastique, est reportée en fonction de la température, pour diverses déformation ε_a (Fig. 1).

On constate que pour des températures situées au dessous de 40 K, une forte dépendance en fonction de la température est observée. L'énergie et le volume d'activation de cette limite anélastique sont successivement mesurés en fonction de la température. La limite vers les basses températures de ces deux paramètres est respectivement de 0,1 eV et 100 b^3 ($\text{b} = \text{intensité du vecteur de Burgers}$), valeurs en parfait accord avec le modèle de la formation activée thermiquement des doubles décrochements [2]. D'autre part, la limite anélastique à basse température, de l'ordre de $10^{-4} \mu$ (Fig. 1), est en accord avec la valeur de la contrainte de Peierls.

Références

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¹⁾ Recherche partiellement subventionnée par le Fonds National Suisse de la Recherche Scientifique, subsidence no. 2.486-0.75.

Effets des impuretés sur le spectre de frottement intérieur d'alliages CFC déformés à basses températures¹⁾

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Le pic de Bordoni observé à basses températures dans les métaux cfc est généralement attribué à la formation activée thermiquement des doubles décrochements [1]. Ce mécanisme ne prévoit pas l'effet que peuvent avoir les impuretés et l'amplitude de mesure.

Le spectre de frottement intérieur que nous avons mesuré dans des alliages dilués d'or dopés au cuivre ou au platine se caractérise par :

- une diminution de la hauteur du pic de Bordoni qui peut être reliée au durcissement des solutions solides
- un maximum du fond de frottement intérieur mesuré à 150K qui peut être relié à une mobilité des dislocations inversement proportionnel à la concentration des substitutionnels

Nous décrivons ailleurs [2] un modèle permettant d'interpréter ces résultats. En

¹⁾ Recherche partiellement subventionnée par le Fonds National Suisse de la Recherche Scientifique, subsidence no. 2.486-0.75.

linéarisant l'équation du mouvement d'une dislocation de longueur L , d'intensité de vecteur de Burgers b (densité $\Lambda \text{ cm}^{-2}$), on montre que l'on peut déterminer l'expression de la hauteur Q^{-1} ($= \Delta/2$) du pic de Bordoni en fonction de la concentration c de substitutionnels. Cette expression peut être développée en série, pour les faibles valeurs de c . On obtient:

$$2Q^{-1} = \Delta = \Delta_0 \left(1 - \frac{L^2 K}{10\gamma} \right)$$

où: $K = \alpha c^{7/6}$; γ = la tension de ligne $= \frac{1}{2}\mu b^2$

$$\Delta_0 = \frac{\Lambda b^2 L^2}{12\gamma} \mu; \mu = \text{module de glissement}$$

A partir des mesures expérimentales du frottement intérieur à faible amplitude ($\sigma = 2 \cdot 10^{-6} \mu$), en fonction de la température autour du pic de Bordoni, on en déduit la hauteur du pic $Q^{-1} = \Delta/2$. Ces mesures sont faites pour diverses concentrations c des impuretés Cu, Pt dans Au, $c = 10^{-5}$ à 10^{-3} . Des mesures semblables faites dans des échantillons très purs ($c = 10^{-6}$) fournissent la valeur de référence $Q_0^{-1} = \Delta_0/2$. De cette détermination expérimentale de Δ et Δ_0 , on reporte $\log(1 - \Delta/\Delta_0)$ en fonction de c (Fig. 1). On obtient une droite, dont la pente est très voisine de $7/6$. Ce qui confirme le modèle utilisé.

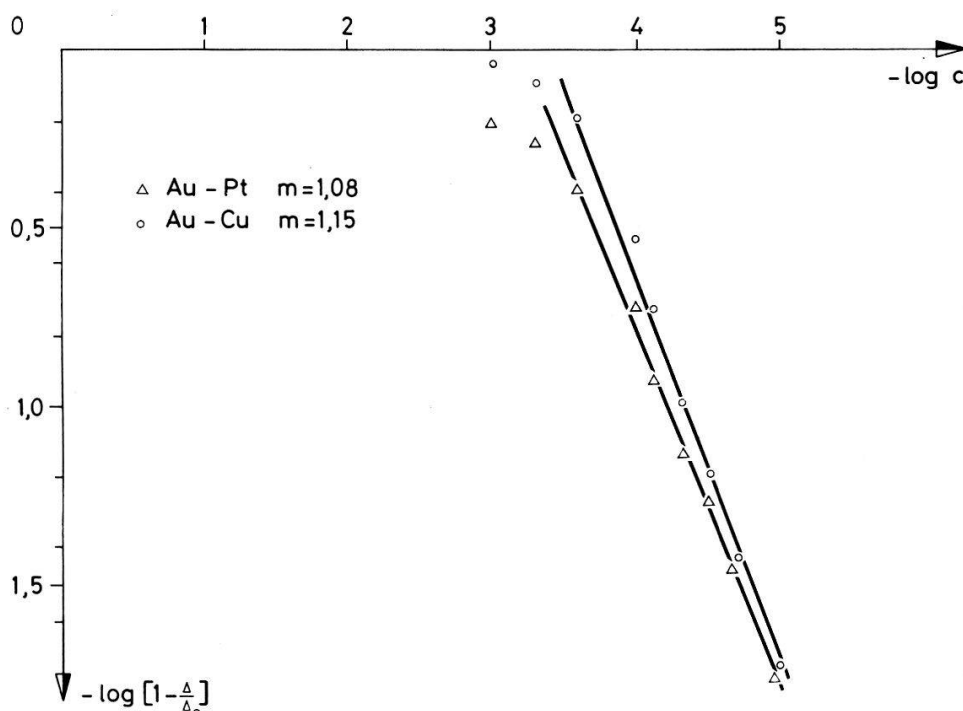


Figure 1

Détermination expérimentale de l'exposant m de la concentration c dans le terme $K = \alpha c^m$. Les mesures de la hauteur Q^{-1} du pic de Bordoni en fonction de la concentration c donnent $\Delta = 2Q^{-1}$ et $\Delta_0 = 2Q_0^{-1}$ pour $c = 10^{-6}$.

Mais les calculs obtenus en linéarisant l'équation d'un mouvement d'une dislocation ne peuvent pas faire apparaître une dépendance amplitudinale. Pour ce faire, on intègre, sans la linéariser, l'équation du mouvement par ordinateur (méthode d'Euler-Romberg [3]). Les résultats sont confrontés avec l'expérience (Fig. 2)

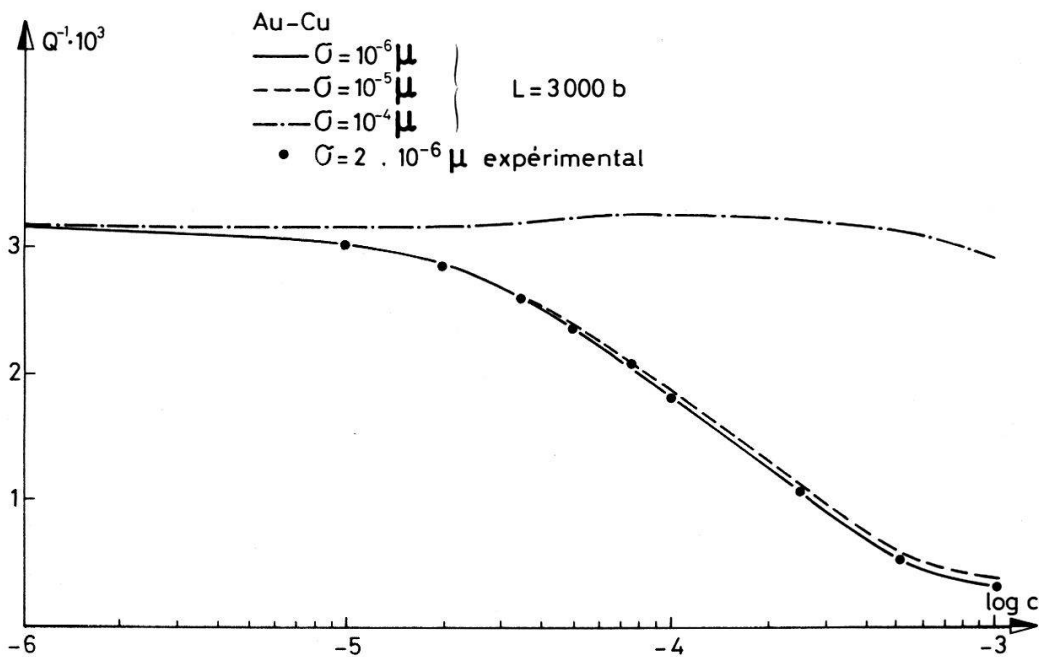


Figure 2

Valeurs théoriques et expérimentales de la hauteur Q^{-1} du pic de Bordoni, en fonction de la concentration c et pour différentes amplitudes de la contrainte appliquée σ (μ = module de glissement).

Ces résultats nous amènent à conclure :

- l'hypothèse de linéarisation de l'équation du mouvement d'une dislocation, pour de faibles contraintes, est justifiée quelle que soit la concentration des impuretés, au moins jusqu'à $c = 10^{-3}$
- l'effet de l'amplitude ne serait sensible qu'à fortes concentrations; ce qui est en accord avec des résultats expérimentaux antérieurs, mais où la concentration des impuretés ne variait que dans un faible domaine. Un programme expérimentale plus complet sera mis en œuvre.

Références

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Elasticité électronique de cisaillement dans les métaux de transition

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Les contributions électroniques aux modules de cisaillement sont calculées en fonction de la température pour quelques métaux de transition cubiques. On peut signaler en passant que dans notre approche du problème on obtient en plus des renseignements sur les potentiels de déformation pouvant être confrontés à l'expérience. Le modèle NFE-TB est utilisé pour la structure de bande. Les résultats sont

comparés avec des calculs récents de Peter *et al.* [1] obtenus par une méthode de perturbation.

Référence

- [1] M. PETER, W. KLOSE, G. ADAM, P. ENTEL et EWA KUDLA, *Helv. Phys. Acta* 47, 807 (1974).

High frequency cyclotron resonance in cadmium

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and

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The frequency ($80 \leq \Omega \leq 460$ GHz) and temperature ($1.3 \leq T \leq 7.5$ K) dependencies of the relaxation times and effective masses of the 3rd band electrons on circular ($\mathbf{H} \parallel \langle 0001 \rangle$ in $(11\bar{2}0)$) and lenticular ($\mathbf{H} \parallel \langle 1010 \rangle$ in $(11\bar{2}0)$) extremal orbits will be presented. The effective masses are independent of Ω to within 0.2% and vary as T^2 [1]. The electron-phonon relaxation time is found to vary as T^3 only at low T for $\Omega = 80$ GHz. A more rapid variation ($\propto T^{5-6}$) is observed at higher Ω and T when inter-sheet scattering becomes important [2, 3].

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Gantmakher-Kaner oscillations in cadmium

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Circularly polarized radio-frequency excitation has been used in measurements of the surface impedance of thin, monocrystalline plates of Cd as a function of the magnetic field \mathbf{H} (0–100 kG) perpendicular to the sample surface. Thus, we can distinguish impedance variations produced by collective, propagating modes of the electromagnetic field-dopplersons when $\mathbf{H} \parallel \langle 0001 \rangle$ [1] – from those produced by a localized group of electrons called Gantmakher–Kaner oscillations (GKO). When $\mathbf{H} \parallel \langle 11\bar{2}0 \rangle, \langle 10\bar{1}0 \rangle$ GKO are observed ($1 \lesssim H \lesssim 50$ kG, 10 kG), as the open orbits in the 2nd band strongly damp all collective modes. The field independent period of the GKO identifies the electrons responsible as those on the rim of the 3rd band lenticular Fermi surface sheet. The temperature dependence of the GKO signal amplitude is used to measure the scattering frequency of these electrons. The results of these measurements will be presented in a subsequent publication.

Reference

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Coincidence-site lattices and Burgers vectors of grain boundary dislocations

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A grain boundary often shows an energy minimum if the translation lattices of the two neighbouring grains have a common lattice with a small primitive cell [1]. This lattice is called the coincidence site lattice (CSL) of the two lattices, which we call 1 and 2. Let Σ_1 and Σ_2 be the numbers of primitive cells of lattices 1 and 2 respectively that are contained in a cell of the CSL. We may assume $\Sigma_1 \geq \Sigma_2$. Let \mathbf{T} be the matrix that expresses a basis for lattice 2 in terms of a basis for lattice 1,

$$\begin{aligned} \mathbf{b}_1^2 &= T_{11}\mathbf{b}_1^1 + T_{21}\mathbf{b}_2^1 + T_{31}\mathbf{b}_3^1 \\ \mathbf{b}_2^2 &= T_{12}\mathbf{b}_1^1 + T_{22}\mathbf{b}_2^1 + T_{32}\mathbf{b}_3^1 \\ \mathbf{b}_3^2 &= T_{13}\mathbf{b}_1^1 + T_{23}\mathbf{b}_2^1 + T_{33}\mathbf{b}_3^1 \end{aligned}$$

and let $\|\mathbf{T}\|$ be the absolute value of its determinant. Then it has been shown in [2] that a (3-dimensional) CSL exists if and only if \mathbf{T} is a rational matrix (i.e. a matrix with rational elements) and that Σ_2 is equal to the least positive integer n such that $n\mathbf{T}$ and $n\|\mathbf{T}\|\mathbf{T}^{-1}$ are integral matrices. Obviously, $\Sigma_1 = \|\mathbf{T}\| \cdot \Sigma_2$. The geometrically possible Burgers vectors of dislocations in the grain boundary are the vectors of the form $\mathbf{v}_1 + \mathbf{v}_2$, where \mathbf{v}_1 is a vector of lattice 1, \mathbf{v}_2 a vector of lattice 2. The lattice formed by these sums is called the DSC lattice [3]. Bonnet [4] and Iwasaki [5] have shown that Σ_2 is equal also to the number of primitive cells of the DSC lattice in a cell of lattice 1.

Bonnet and Durand describe in [6] a method to determine the DSC lattice, which does not always work as has been pointed out by Bonnet in [7], where he gives generally valid methods to determine DSC lattice and CSL. Knowing the value of Σ_2 , his methods, which are easily transformed into computer programs, can be simplified even a bit further.

- 1) Determine N_1 , the least positive integer such that $N_1 T_{11}$, $N_1 T_{21}$, and $N_1 T_{31}$ are integers.
- 2) Determine n_{12} and N_2 , where N_2 is the least positive integer such that $n_{12} T_{11} + N_2 T_{12}$, $n_{12} T_{21} + N_2 T_{22}$, and $n_{12} T_{31} + N_2 T_{32}$ are integers for a suitable choice of the integer n_{12} in the range $0 \leq n_{12} < N_1$.
- 3) Compute $N_3 = \Sigma_2 / N_1 N_2$.

Determination of the CSL. Determine the integers n_{13} and n_{23} satisfying $0 \leq n_{13} < N_1$ and $0 \leq n_{23} < N_2$, for which $n_{13} T_{11} + n_{23} T_{12} + N_3 T_{13}$, $n_{13} T_{21} + n_{23} T_{22} + N_3 T_{23}$, and $n_{13} T_{31} + n_{23} T_{32} + N_3 T_{33}$ are integers. A basis for the CSL is then given by $\mathbf{b}_1^c = N_1 \mathbf{b}_1^2$, $\mathbf{b}_2^c = n_{12} \mathbf{b}_1^2 + N_2 \mathbf{b}_2^2$, $\mathbf{b}_3^c = n_{13} \mathbf{b}_1^2 + n_{23} \mathbf{b}_2^2 + N_3 \mathbf{b}_3^2$.

Determination of the DSCL. (By the internal coordinates of a vector \mathbf{v}_2 of lattice 2 we shall understand the components of a vector $\mathbf{v}_1 + \mathbf{v}_2$ expressed in our basis of

lattice 1, where \mathbf{v}_1 is that vector of lattice 1 for which all the components of $\mathbf{v}_1 + \mathbf{v}_2$ become non-negative numbers less than 1.) Compute the internal coordinates of the vectors $n_1 \mathbf{b}_1^2 + n_2 \mathbf{b}_2^2 + n_3 \mathbf{b}_3^2$, where the n_i are integers satisfying $0 \leq n_i < N_i$, for $i = 1, 2, 3$. Determine among these Σ_2 triples of internal coordinates

- a) the one of form $(k_{11}, 0, 0)$ with least positive k_{11} .
- b) the one of form $(k_{12}, k_{22}, 0)$ with least positive k_{22} and $0 \leq k_{12} < k_{11}$.
- c) the one of form (k_{13}, k_{23}, k_{33}) with least positive k_{33} and $0 \leq k_{13} < k_{11}$, $0 \leq k_{23} < k_{22}$.

A basis for the DSC lattice is then given by $\mathbf{b}_1^D = k_{11} \mathbf{b}_1^1$, $\mathbf{b}_2^D = k_{12} \mathbf{b}_1^1 + k_{22} \mathbf{b}_2^1$, $\mathbf{b}_3^D = k_{13} \mathbf{b}_1^1 + k_{23} \mathbf{b}_2^1 + k_{33} \mathbf{b}_3^1$.

A paper where these results are applied to the study of DSC lattices and CSLs for hexagonal lattices 1 and 2 with various ratios of c/a is in preparation.

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Anomale optische Eigenschaften von stark überhitzten Metalloberflächen

B. ELIASSON, S. STRÄSSLER, J. WULLSCHLEGER und H. R. ZELLER

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Beim Laserbohren von Metallen zeigt sich oberhalb einer kritischen Leistung ein drastischer Abfall der Reflexion auf $\lesssim 0.1$.

In der Literatur wurde postuliert, dass der sich bildende überhitzte Flüssigkeitsfilm bis in die Nähe der kritischen Temperatur getrieben wird und dabei einen Metall-Isolator Uebergang macht. Wir zeigen, dass der Film bereits vorher instabil wird und in kleine Tröpfchen zerfällt. Unser Modell beschreibt die optischen Eigenschaften der Oberfläche *ohne* dass ein Phasenübergang im Metall postuliert werden muss.

Photoelektronen, Halbeiter

Photoélectrons induits par rayons X dans l'oxyde de nickel; influence de la température

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The energy and angular dependence of X-ray photoelectrons from NiO has been measured as a function of temperature. The energy distribution in the Ni_{2p} region is deeply modified following heating. The main oxide lines and their satellites on the high binding energy side are gone and a sharp peak attributed to Ni metal appears. The relevant surface layer can be removed by Argon ion bombardment. This behaviour has been observed in single crystals and oxidized metal samples as well.

1. Introduction

Les spectres XPS des oxydes de certains métaux de transition présentent, dans la région des niveaux $2p$ et $3d$, une structure assez inhabituelle. Les lignes de photoélectrons, respectivement la bande de valence, sont accompagnées, dans un domaine s'étendant à environ 10 eV du côté des énergies de liaison croissantes, de satellites très intenses. Ces derniers ne sont pas explicables par le modèle à une particule et plusieurs hypothèses ont été avancées pour expliquer leur origine. Dans le cas particulier de l'oxyde de nickel NiO (Fig. 1), qui est un isolateur antiferromagnétique, Novakov [1] a fait appel à une excitation multiple. Une transition entre la bande de valence et la bande de conduction, vide dans l'état initial, accompagne l'émission du photoélectron dont l'énergie cinétique est diminuée de la quantité correspondante, ce qui donne naissance au satellite. Cette interprétation à elle seule n'est pas à même

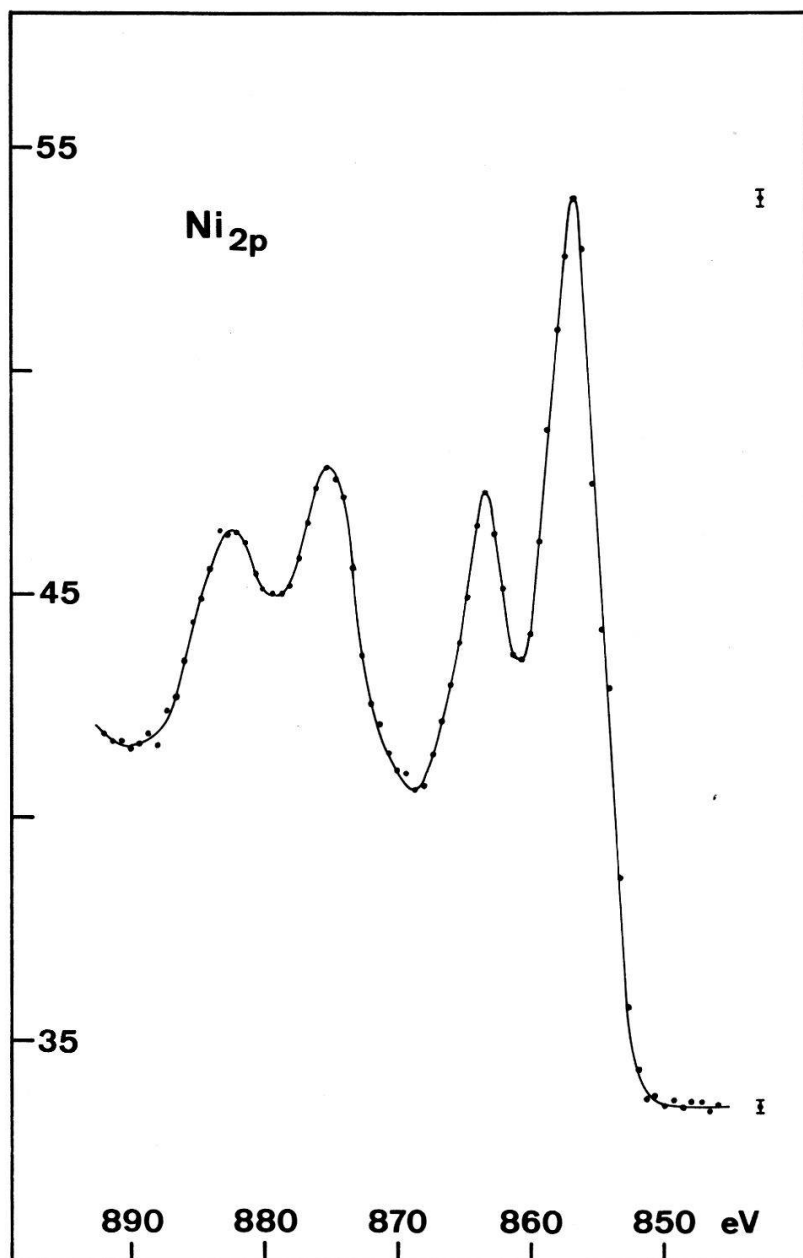


Figure 1
Spectre XPS du niveau Ni_{2p} avec ses satellites dans NiO.

d'expliquer les détails des spectres mesurés. Un point de vue différent est d'admettre que les spectres XPS sont le reflet de l'état initial [2], soit plutôt de l'état final, comme l'ont proposé Wertheim *et al.* [3]. Deux phénomènes sont alors à considérer. Le couplage du trou laissé derrière lui par le photoélectron avec la couche $3d$ incomplète conduit à une structure de multiplets. Il est cependant difficile d'expliquer de cette manière la forte intensité de certains satellites. A cette propriété de l'atome isolé se superpose l'effet du champ cristallin qui résulte de la présence des voisins immédiats de l'ion. Par suite de la rupture de la symétrie, la dégénérescence en énergie est partiellement levée et l'ion peut occuper dans l'état final des états excités discrets auxquels correspondent alors les satellites observés. Aucun modèle ne peut prévoir actuellement les énergies liées à ces deux modes d'excitation. Il est admis que, dans le cas des ions des métaux $3d$, l'effet du champ cristallin l'emporte sur le couplage qui

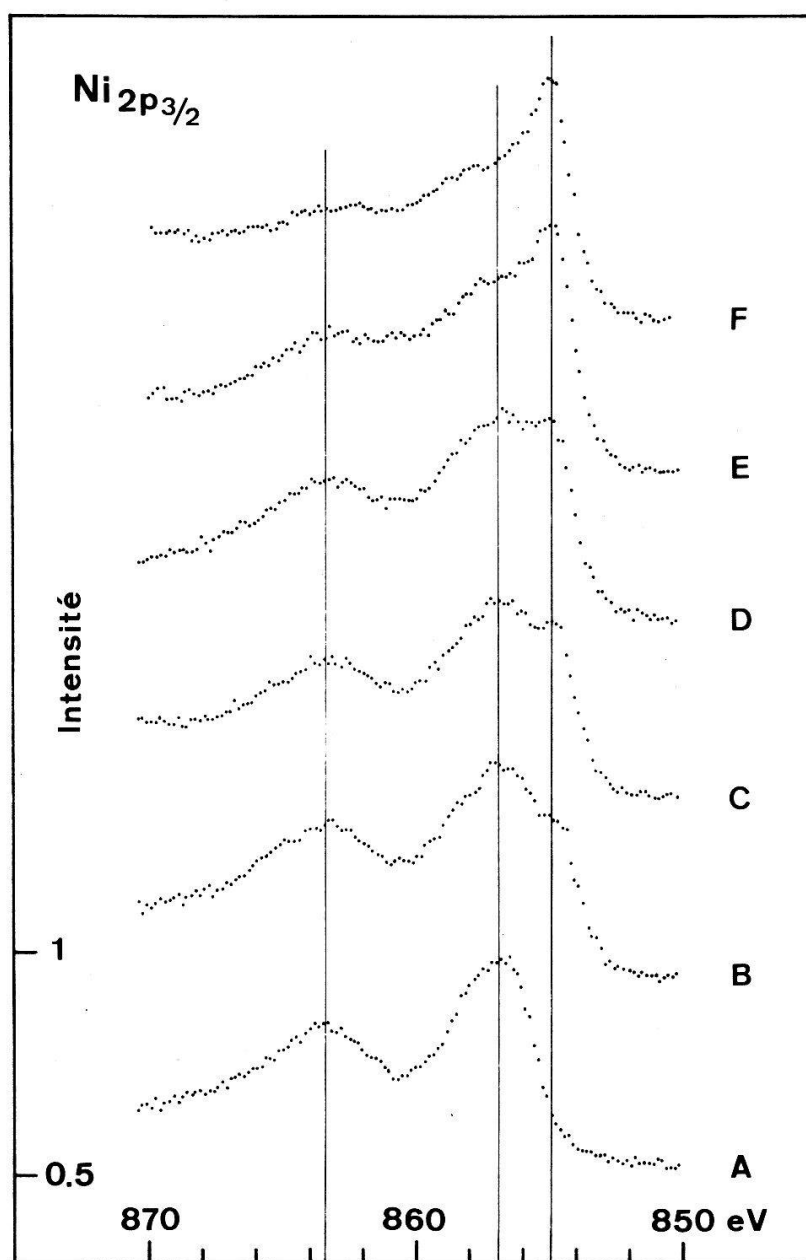


Figure 2

Effet de la température dans la région $Ni_{2p_{3/2}}$. Le spectre A a été relevé à température ambiante, les suivants entre 250 et 300°C à des intervalles d'une demi'heure.

conduit à la structure de multiplets. Sugano *et al.* [4] ont déterminé les énergies d'excitation à partir de mesures optiques. Il faut cependant remarquer qu'une comparaison directe avec les spectres XPS n'est pas possible puisque les états qui interviennent ne sont pas les mêmes, l'état final étant ionisé en photoémission. Il semble donc probable que les deux modes d'excitation coexistent et une interprétation correcte doit reposer sur un traitement global de ces phénomènes.

Le calcul de structure de bande [5] est en désaccord avec les résultats expérimentaux, particulièrement du fait qu'il conduit à attribuer un caractère métallique à l'isolateur NiO. Une autre approche théorique [6], qui applique la méthode du champ self-consistant SCF- $X\alpha$ à un cluster NiO_6^{10-} , prévoit des états très localisés qui décrivent les propriétés optiques de manière satisfaisante et d'autre part des transferts de charge métal-ligand qui pourraient être à l'origine des satellites les plus fortement liés. Une analyse détaillée basée sur ce modèle a été faite par Kim [7]. Il faut remarquer que les transitions vers des états localisés obéissent aux règles de sélection dipolaires, alors que les transferts de charge sont du type monopolaire et suivent les règles de sélection correspondantes. Il est ainsi possible d'expliquer quantitativement les intensités observées.

2. Mesures XPS

Une application immédiate de ces prévisions à l'interprétation des spectres XPS est cependant sujette à caution. En effet, les résultats expérimentaux ne sont pas tous en accord. L'intensité et même l'existence des satellites ont été reliées étroitement à un excès d'oxygène, adsorbé ou dû à la présence de vapeur d'eau à la surface de l'échantillon. En conséquence, Novakov *et al.* [2] ont attribué l'existence des satellites dans NiO à un phénomène de surface mais, contrairement au cas de CuO, le chauffage de l'échantillon à 200°C n'a pas été accompagné d'une réduction et le spectre dans la région 2p est demeuré inchangé. Rosencwaig *et al.* [8] ont chauffé leur échantillon pendant plusieurs heures à 200°C. Le changement observé dans la ligne 1s de l'oxygène est attribué à la désorption de vapeur d'eau. Le fait que le spectre n'est pas modifié dans la région 2p conduit les auteurs à admettre que les satellites sont dus à un effet de volume et non de surface. Kim [9] attribue certains des effets observés à des défauts de structure, en particulier à la présence en surface de Ni_2O_3 .

Les investigations de Hirokawa *et al.* [10] sont en contradiction avec ces travaux en ce qui concerne la stabilité thermique des échantillons. Le spectre d'un échantillon chauffé sous vide pendant 10 heures à 350°C est profondément modifié. Les lignes de l'oxyde et leurs satellites disparaissent au détriment de lignes caractéristiques du nickel métallique. Les différences de comportement ne s'expliquent pas par le mode de préparation puisqu'elles ont été observées sur des échantillons obtenus par oxydation d'une feuille ou d'une poudre métallique ou par réduction de l'hydroxyde.

3. Résultats et discussion

Le présent travail a pour but d'étudier le comportement de différents échantillons de NiO en fonction de la température. Pour mettre en évidence les phénomènes qui se jouent au voisinage immédiat de la surface, nous avons utilisé la dépendance angulaire de la photoémission [11]. Les spectres que nous présentons ont été obtenus sur des échantillons monocristallins de NiO vert et sur des plaques métalliques

traitées plusieurs heures à 800°C sous atmosphère d'oxygène, à l'air libre ou oxydées *in situ* à 300°C sous pression réduite. Immédiatement après l'introduction dans le spectromètre, tous les échantillons montrent un excédent d'oxygène. La ligne 1s correspondante présente deux composantes, à environ 2 eV l'une de l'autre. L'intensité de la composante la plus liée augmente pour les photoélectrons émis à angle rasant. Cette composante est donc bien due à des atomes de surface. Le fait qu'elle soit présente sur les échantillons préparés *in situ* avec de l'oxygène sec montre qu'il ne s'agit pas de vapeur d'eau. Son intensité relative diminue après chauffage modéré. Après deux heures à 300°C, l'intensité de l'autre composante, attribuée à l'oxyde, est réduite environ de moitié. Après prolongation du chauffage pendant plusieurs heures, cette ligne n'est pratiquement plus discernable. Ce comportement est accompagné d'une modification sensible du spectre dans la région 2p, voir Fig. 2. L'éta-lonnage est basé sur une énergie de liaison de 854,7 eV de la ligne 2 p_{3/2} du métal. Le spectre A est celui d'un échantillon n'ayant pas subi de traitement, mesuré à la température ambiante. Les spectres consécutifs ont été relevés entre 250 et 300°C à des intervalles d'environ une demi-heure. On voit qu'il apparaît, du côté des faibles énergies de liaison, une ligne qui croît au dépend des lignes de l'oxyde. Son énergie correspond à celle de la ligne 2p_{3/2} du nickel, comme l'a montré une mesure de contrôle sur un échantillon métallique. La séparation en énergie, de l'ordre de 2 eV, est attribuée au déplacement chimique entre le métal et l'oxyde, augmentée éventuellement d'un faible déplacement dû à la charge de l'échantillon. Le spectre F correspond à celui d'un échantillon métallique faiblement oxydé. La présence d'une couche métallique en surface est confirmée par l'augmentation de l'intensité quand on s'approche de l'émission sous un angle rasant. La croissance de cette couche n'est pas réversible, le spectre ne change pas quand la température est abaissée. Elle peut par contre être éliminée par un bombardement d'ions d'argon de faible énergie. Un traitement pendant 5 min. avec un courant de 10 µA suffit à l'éliminer complètement. Le spectre mesuré immédiatement après a de nouveau l'allure A et le cycle peut être recommencé. Répétons que ces observations ont été faites sur tous les échantillons examinés. Des résultats analogues ont été obtenus dans le domaine des électrons 3d.

Les mesures de dépendance angulaire montrent que l'intensité relative du satellite situé à 6.5 eV au-dessous de la ligne 2p ne dépend pas de l'angle d'émission, ce qui ne serait pas le cas pour un effet de surface. On peut donc en conclure que ce satellite est bien dû à un effet de volume.

Le mécanisme responsable de la réduction sous vide de l'oxyde, par ailleurs stable à l'atmosphère, n'est pas clair. L'hydrogène présent dans le gaz résiduel de notre système, évacué par une pompe ionique, pourrait en être la cause. Cependant, un essai de réduction par l'hydrogène *in situ* n'a pas eu d'influence sur le spectre mesuré. D'autre part, nous n'avons jamais observé de réduction des oxydes d'autres métaux. Rappelons encore que Hirokawa *et al.* [10] ont observé la réduction de NiO dans un système évacué au-dessous de 10⁻⁹ Torr par des pompes à diffusion munies de trappes à azote liquide.

En conclusion, les mesures que nous présentons montrent que la stabilité de certains composés n'est pas forcément garantie dans les mesures XPS. Les réactions chimiques dans les couches superficielles, seules accessibles par cette méthode, ne sont pas claires. La dépendance angulaire de la photoémission apporte une aide précieuse dans l'étude de tels problèmes. Dans le cas de NiO, l'effet d'une mince couche métallique à la surface de l'échantillon pourrait conduire à une interprétation erronée du spectre de l'oxyde.

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Temperature dependence of photoelectron spin polarization on Fe₃O₄

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The first temperature measurements of the photoelectron spin polarization (ESP) on a ferrimagnet are reported. The (ESP) decreases faster with temperature than the bulk magnetization. We interpret the data with the characteristic magnetic properties of the surface and photoelectron escape depth.

Photoemission aus UC, UN, US mit Synchrotronstrahlung

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Die ersten Energieverteilungskurven der Photoelektronen aus UC, UN, US mit Synchrotronstrahlung der Energie bis 90 eV werden vorgestellt. Die beobachteten Strukturen werden den elektronischen Zuständen zugeordnet und im Falle des Ferromagneten US mit bereits vorhandenen Ultraviolett-Photoemissionsspektren verlichen.

Photoemission studies on lead chalcogenides

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Photoemission spectra for lead chalcogenides PbS, PbSe, and PbTe which are polar semiconductors, obtained in the photon energy range of $5 < h\nu < 90$ eV using synchrotron radiation from a storage ring are interpreted with emphasis on band widths and critical-point positions with respect to valence-band edge. The results are compared with former ones obtained with X-rays and resonance lamps, as well as with theoretical predictions based on OPW, EPM, and APW calculations. Besides their technological importance as detectors of infrared radiation, lead chalcogenides

are recognized as potential candidates for efficient spin-polarized electron sources when irradiated with circularly polarized light.

Photoleitung von SmSe, SmTe, YbSe, YbTe und FeS₂

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Die Photoleitung von YbSe- und YbTe-Einkristallen ergibt gleiche Energielücken für die Beweglichkeit wie für die Zustandsdichte (1.5 bzw. 1.7 eV) und ist oberhalb der Kante maximal bei minimalem Absorptionskoeffizienten, verläuft also qualitativ gleich wie bei den Eu-Monochalkogeniden. Dieses Verhalten kann erklärt werden, wenn man für die Ladungsträger an der Oberfläche eine kürzere Lebensdauer als im Inneren annimmt. Auch SmTe kann grösstenteils in diesem Sinne interpretiert werden. Dagegen ändert sich das Spektrum von SmSe-Kristallen oberhalb von 40 K drastisch und wird, wie bei dünnen Filmen, proportional zur Absorption, was auf eine ortsunabhängige Lebensdauer hinweist. Die Photoleitung aller dieser Substanzen ist unterhalb *ca.* 150 K gering temperaturaktiviert (typisch: $E_a = 20$ meV) und erreicht bei höheren Temperaturen ein relatives Maximum.

Die spektrale Abhängigkeit der Photoleitung von FeS₂ bestätigt den Wert für die Energielücke, wie er aus der optischen Absorption bekannt ist (0.95 eV).

Resonante Ramanstreuung an Eu- und Yb-Chalkogeniden

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Die resonante Ramanstreuung an Eu- und Yb-Chalkogeniden verletzt die üblichen Auswahlregeln. Ähnlich wie bei der Lumineszenz wird das Streulicht als inkohärent aufgefasst. Damit wird die resonante Ramanstreuung zum Zweistufenprozess. Die Streuung ergibt sich aus der spontanen Emission eines durch die Absorption eines Photons präparierten Zustandes. Ein einfaches Modell für das System liefert insbesondere Temperaturabhängigkeiten der Streuintensität.

Site dependence of impurity ionization energies in semiconductors

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The binding energy of shallow impurities in semiconductors is theoretically shown to depend on the particular location of the impurity atom in the unit cell. This site-dependence results from corrections to the effective-mass-approximation which are related to the Bloch nature of the electronic wave functions. The effect is larger in more ionic host crystals and for heavier bound particles. Application to isocoric acceptors in tetrahedral semiconducting compounds gives a larger binding energy for substitutions on the anion site, in agreement with experimental data.

Beweglichkeit der Ladungsträger in Quecksilberjodid- und Bleijodid-Kristallen

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HgJ₂ und PbJ₂-Kristalle weisen eine ausgeprägte Schichtstruktur auf, was sich u.a. in einer Anisotropie der Gleichstromleitfähigkeit äussert. Im Vortrag soll über unsere Messungen der Temperaturabhängigkeit der Driftbeweglichkeit von Löcher und Elektronen in Richtung senkrecht zu den Schichten berichtet werden.

Varia

Hydride von La-Ni-Verbindungen

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Wasserstoff hat als möglicher Energieträger der Zukunft zunehmendes Interesse gefunden. Als vielversprechende Speichermöglichkeit für Wasserstoff bieten sich Metallhydride an, die z.T. grosse Mengen von Wasserstoff reversibel absorbieren und wieder desorbieren können. Besonders interessant ist hierbei die intermetallische Verbindung LaNi₅, die bei Zimmertemperatur und 2 bar Wasserstoffdruck etwa 6 Wasserstoffatome pro Formeleinheit zu speichern vermag [1]. Die daraus resultierende Wasserstoffdichte könnte durch Kompression von gasförmigem Wasserstoff erst bei einem Druck von 1600 bar erreicht werden.

Im Rahmen von Untersuchungen zum Verständnis der Wasserstoffabsorption in LaNi₅ und der Suche nach neuen Metallhydriden haben wir La₇Ni₃ und LaNi hergestellt und hydriert. Die nickelreichen La-Ni-Verbindungen wurden von Maeland *et al.* [2] untersucht. Die von uns hergestellten La₇Ni₃-Proben weisen eine hexagonale Struktur des Th₇Fe₃-Typs auf. Die Hydrierung lief bei 20°C unter 20 bar Wasserstoffdruck nach einer Aktivierungszeit von einigen Minuten in wenigen Sekunden unter deutlicher Wärmeentwicklung ab. Sie führte zu einem thermisch stabilen Hydrid, das mit der Formel La₇Ni₃H₁₉ beschrieben werden kann. Dies entspricht in Atomprozenten einem etwa doppelt so hohen Wasserstoffgehalt wie in LaNi₅H₆. Bis 700°C war keine Desorption gegen Atmosphärendruck feststellbar. Die Bildungswärme wurde kalorimetrisch zu 36,5 kcal/mol H₂ bestimmt. Die Debye-Scherrer Röntgenaufnahmen des Hydrids können mit einem kubischflächenzentrierten Gitter mit $a = 5.62 \pm 0.04$ Å indiziert werden. Die Hydrierung führt also zu einer wesentlichen Strukturänderung. Der Vergleich der gemessenen Gitterkonstante mit derjenigen des ebenfalls kubisch-flächenzentrierten Lanthanhydrids (für LaH₂ $a = 5.67$ Å, für LaH₃ $a = 5.60$ Å [3]) legt die Vermutung nahe, dass sich La₇Ni₃ beim Hydrieren in Lanthanhydrid und Nickel zersetzt. Die Röntgenfilme enthalten jedoch keine Nickellinien (vergl. auch [2]).

Zur Abklärung einer eventuellen Nickel-Ausscheidung wurden die Temperatur- und Feldabhängigkeit der magnetischen Suszeptibilität gemessen. Die Resultate schliessen jedoch eine Nickelphase aus. Sie lassen aber die Möglichkeit offen, dass

sich das Hydrid in Lanthanhydrid und LaNi_5 zersetzt. Die Abklärungen hierüber sind noch im Gang.

Die Bildungsenthalpie von $\text{La}_7\text{Ni}_3\text{H}_{19}$ und $\text{LaNiH}_{3.9}$ wurde mit derjenigen von LaNi_5H_6 [1] und LaNi_2H_2 [2] verglichen. Der Betrag der Bildungsenthalpie nimmt linear mit steigendem La-Gehalt zu. Die Anwendung einer phänomenologischen Theorie von Miedema [4] führt unter Berücksichtigung der Grössenunterschiede der La- und Ni-Ionen zu einer guten Beschreibung der gemessenen Werte.

Eine ausführlichere Beschreibung der Arbeit ist in Vorbereitung.

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Thermoanalytical study of phase transitions in $(\text{C}_n\text{H}_{2n+1}\text{NH}_3)_2\text{MCl}_4$ compounds

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A thermoanalytical study of the cadmium, manganese and copper series with alkyl chains of up to 18 C-atoms revealed phase transitions with enthalpies of up to 80 kJoule/mol. Possible transition mechanisms are discussed. The observed large effects are probably due to a hydrocarbon-chain 'melting' in a rigid halide octahedra matrix.

Infrared study of a phase transition in $(\text{C}_{10}\text{H}_{21}\text{NH}_3)_2\text{CuCl}_4$

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IR transmission measurements were carried out on single crystals of $(\text{C}_{10}\text{H}_{21}\text{NH}_3)_2\text{CuCl}_4$ between 4000 cm^{-1} and 300 cm^{-1} in the temperature range from 280 K to 315 K. On heating drastic changes in the spectrum take place at 307 ± 1 K. Most lines associated with the molecular vibrations of the alkyl chains have lost their intensity above the transition point. The thermal hysteresis indicates a first order phase transition. On heating the transition takes place as a one step process while the reordering of the alkyl chains requires two steps as the temperature is lowered through the transition point.

Ein universeller Zusammenhang zwischen den Parametern der dielektrischen Dispersion von Eis I_h

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Die Leitfähigkeit und die dielektrische Dispersion von Eis wird auf die Bewegung von Orientierungs- und Ionendefekten zurückgeführt. Unter wenig einschränkenden Bedingungen wird aus der Theorie zwischen der statischen Dielektrizitätskonstanten ϵ_s und dem Quotienten σ_0/σ_∞ aus Niederfrequenz- und Hochfrequenzleitfähigkeit die universelle Relation $(\epsilon_s - \epsilon_\infty)\epsilon_0 T = K\Lambda'(\sigma_0/\sigma_\infty)$ hergeleitet. Anstelle von $\epsilon_s - \epsilon_\infty$ kann auch das Produkt aus σ_∞ und der Relaxationszeit τ benützt werden, was zur analogen Beziehung $\sigma_\infty \tau \cdot T = K\Theta'(\sigma_0/\sigma_\infty)$ führt. K ist eine Konstante und Λ' resp Θ' sind Funktionen des Leitfähigkeitsverhältnisses σ_0/σ_∞ . Beide Relationen gelten im Existenzbereich der Phase I_h für alle Temperaturen, Drucke und weitgehend alle Dotierungskonzentrationen. Anhand experimenteller Daten von reinem sowie von HF- und NH_3 -dotiertem Eis wird dieser Zusammenhang verifiziert. Dies ergibt eine sehr gute quantitative Bestätigung der Theorie. Einige weitere Schlüsse, die gezogen werden können, werden diskutiert.

Kristallfeld-Aufspaltung in CeAs

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Unter dem Einfluss des oktaedrischen Kristallfeldes spaltet das ${}^2F_{5/2}$ Grundmultiplett von CeAs in ein Dublett Γ_7 und ein Quartett Γ_8 auf. Das temperaturabhängige Verhalten dieses Kristallfeld-Uebergangs wurde mittels inelastischer Neutronenstreuung untersucht. Im Gegensatz zu anderen Seltenen Erden Monopnictiden ist die beobachtete Kristallfeld-Aufspaltung stark temperaturabhängig und zudem wesentlich kleiner als man nach einem Punktladungs-Modell erwarten würde. Auch die gemessene Linienbreite ist stark temperaturabhängig. Sie wird anhand verschiedener Relaxations-Mechanismen diskutiert.

Nouveaux effets qui permettent d'augmenter la réflectivité de quartz dans des expériences de diffraction

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Des lames monocristallines de quartz sont fréquemment utilisées dans les spectromètres à cristal. Nous avons observé que l'application d'un champ électrique alternatif ou continu influence les propriétés de diffraction des plans (110). Nous avons constaté dans certains cas une augmentation importante de la réflectivité. L'interprétation des résultats obtenus sera discutée.

Résonances magnétiques à 1, 2 et 3 quanta en champ oscillant de direction quelconque

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Résumé. La position des transitions à 1, 2 et 3 quanta ainsi que la distance relative de ces transitions sont étudiées expérimentalement pour diverses valeurs de l'angle θ entre le champ statique et le champ oscillant par prépolarisation et précession libre – Les résultats obtenus sont comparés aux données théoriques antérieurement évaluées. On note l'apparition de nouvelles transitions pour des valeurs données de θ , différentes de 0 ou 90°.

1. Introduction

Cette communication complète des travaux antérieurs concernant l'analyse théorique et l'observation expérimentale de transitions dipolaires magnétiques, obtenues hors résonance pour des protons de substances liquides. Ces protons sont soumis à un champ constant H_0 (le champ magnétique terrestre) et un champ alternatif $H_1(\omega)$ dont la fréquence varie entre 0 et le quadruple de la fréquence de Larmor et l'amplitude entre 0 et $H_0/4$. Les orientations mutuelles de ces deux champs étaient soit perpendiculaires [1, 2] soit parallèles [3].

Les conditions expérimentales – précession libre dans le champ exciteur [$H_0 + H_1(\omega)$] après prépolarisation – sont décrites ailleurs [4].

2. Position des transitions

A) Lorsqu'on superpose au champ constant, un champ alternatif qui lui est perpendiculaire, on peut approximer ce dernier comme la somme de deux champs tournants.

Le système est alors caractérisé par 3 fréquences propres, dont l'une est simplement la fréquence ω du champ exciteur, et dont les deux autres s'écartent de ω_0 par les expressions

$$-2\omega_0 + [(\omega - \omega_0)^2 + \omega_1^2]^{1/2} + [(\omega + \omega_0)^2 + \omega_1^2]^{1/2} \quad (1)$$

$$\text{et } -2\omega_0 + [(\omega + \omega_0)^2 + \omega_1^2]^{1/2} - [(\omega - \omega_0)^2 + \omega_1^2]^{1/2} \quad (2)$$

Ces expressions, pour $\omega_1 \ll |\omega - \omega_0|$ se réduisent à :

$$\frac{1}{2} \left[\frac{\omega_1^2}{\omega_0 - \omega} + \frac{\omega_1^2}{\omega_0 + \omega} \right] \quad (1')$$

$$\text{et } \frac{1}{2} \left[\frac{\omega_1^2}{\omega_0 - \omega} - \frac{\omega_1^2}{\omega_0 + \omega} \right] \quad (2')$$

(dans ces formules $\omega_1 = \gamma H_1$, γ rapport gyromagnétique des protons [5]).

La Figure 1 montre les déplacements de fréquence (effet Bloch–Siegert [6]) pour $0 < \omega < 3000$ Hz et $\omega_1 = 250$ Hz.

Cette figure montre approximativement la position des transitions à 2 et 3 quanta; Il suffit d'ajouter à $\omega_0/2$ et $\omega_0/3$ les déplacements respectifs $\Delta\omega_2$ et $\Delta\omega_3$ correspondant à ces deux transitions.

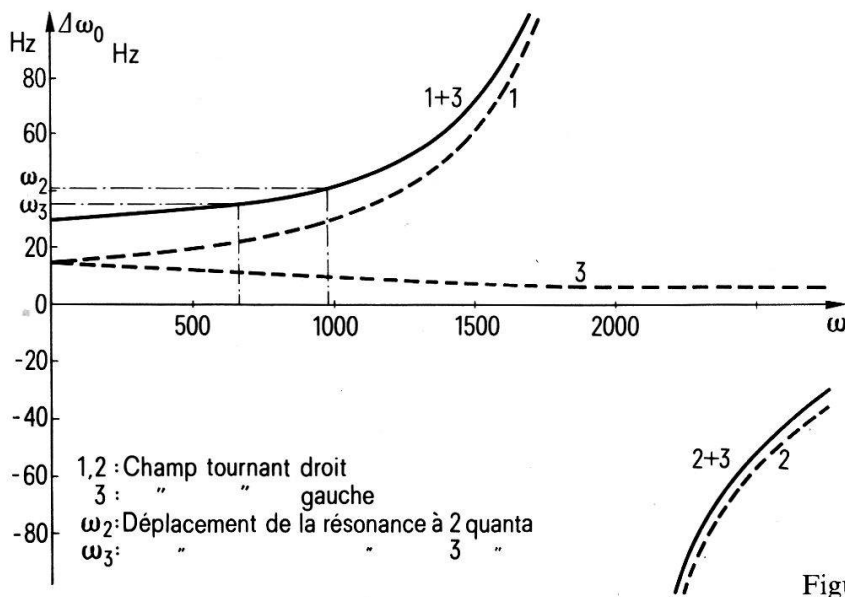


Figure 1

La transition à 1 quantum ($\omega = \omega_0$) n'est déplacée que d'une valeur de l'ordre de $\omega_1^2/4\omega_0$ qui, ici, vaut environ 8 Hz.

B) Si le champ alternatif est parallèle au champ constant, on observe des fréquences latérales $\omega_0 \pm p\omega$ (p entier) et les résonances de cohérence, $\omega_c = \omega_0/n$ (n entier positif). Pour les résonances de cohérence, il n'y a aucun effet Bloch-Siegert.

C) Si l'orientation relative des deux champs H_0 et $H_1(\omega)$ est quelconque, tous les déplacements observés peuvent s'interpréter à une bonne approximation: Il suffit d'admettre que pour une orientation donnée, seule la composante de $H_1(\omega)$ perpendiculaire à H_0 est la cause du déplacement.

Si θ est l'angle de $H_1(\omega)$ avec H_0 , il suffira donc de remplacer ω_1 par $\omega_1 \sin \theta$ (par exemple dans les formules 1, 2, 1' et 2') pour calculer les déplacements effectifs.

3. Fréquence des transitions à 2 et 3 quanta

Nous n'avons pas examiné l'évolution de l'effet Bloch-Siegert avec l'angle θ pour la transition à 1 quantum en raison de la très faible valeur maximum de cet écart.

Pour les transitions à 2 et 3 quanta, l'approximation des formules 1' et 2' donne une évolution de l'écart en $\sin^2 \theta$; comme pour les transitions choisies ω_1 (300 Hz pour la transition à 2 quanta et 400 Hz pour cette à 3 quanta) n'est pas très petit devant $|\omega_0 - \omega|$ cette approximation n'est pas tout à fait valable comme on peut le voir sur la Figure 2.

4. Largeurs de raie (ou distance entre les deux transitions extrêmes, distance qui se réduit, dans une expérience classique de résonance, à la distance des deux extrema de la courbe de dispersion)

A) Transition à 1 quantum

La Figure 3 donne la position des deux transitions lorsque θ varie entre 10° et 90° . On évaluera la distance des deux raies correspondant à un angle donné en faisant la différence des expressions (1) et (2) dans lesquelles on remplacera ω_1 par $\omega_1 \sin \theta$ et on fera

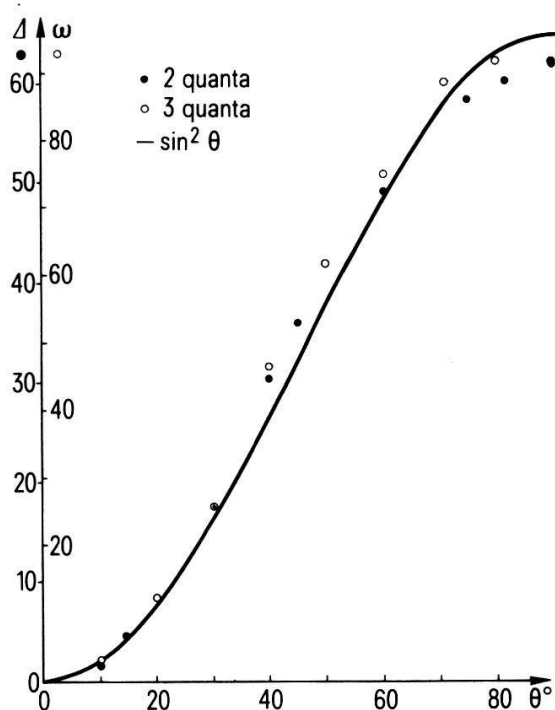


Figure 2

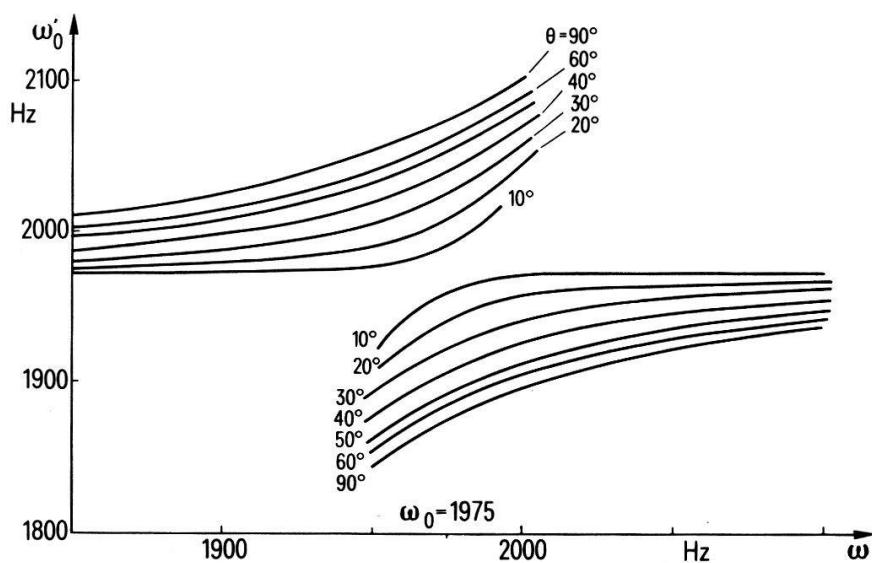


Figure 3

$$\omega = \omega_0$$

Ce calcul donne, pour l'écart, la valeur :

$$2\omega_1 \sin \theta$$

La Figure 4 montre un excellent accord avec l'expérience.

B) Transitions à 2 et 3 quanta

Nous avons comparé, sur cette même Figure 4, les écarts observés à des lois en $|\sin 2\theta|$ et $|\sin 3\theta|$.

L'accord avec l'expérience est moins bons que pour 1 quantum.

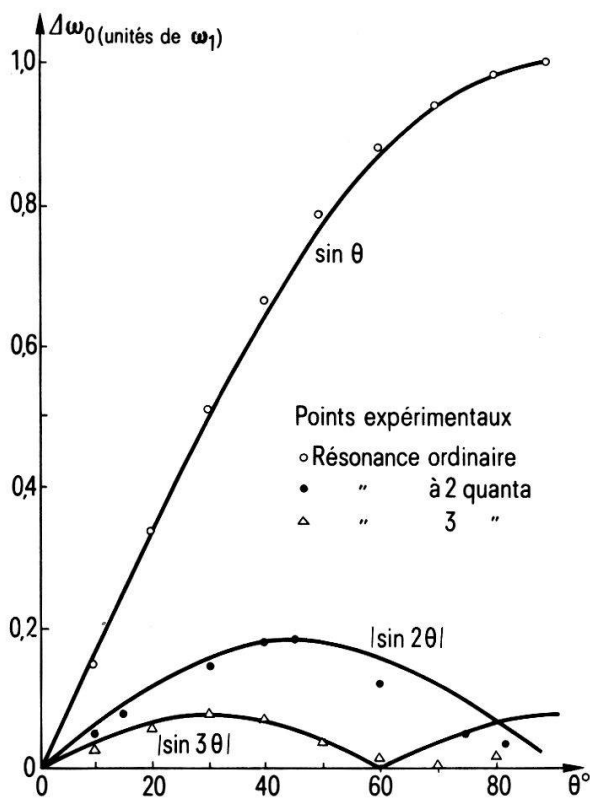


Figure 4

Remarquons que la transition à deux quanta est une résonance interdite (résonance de cohérence) pour $\theta = 0^\circ$ et 90° et qu'elle est permise entre ces valeurs (max. pour $\theta = 45^\circ$).

De même la transition à 3 quanta est une résonance pour $\theta = 0^\circ$ et interdite pour $\theta \approx 60^\circ$; elle est par contre permise pour $\theta = 30^\circ$ et 90° .

On voit ainsi apparaître un spectre nouveau en fonction de l'angle θ du champ constant et du champ alternatif.

L'analyse théorique de ces cas est en cours.

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Lineare Antwort für ein stochastisches Modell einer bistabilen Tunnel-Diode

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Das Verhalten der nichtlinearen Fluktuationen in einer bistabilen Tunnel-Diode kann mit Hilfe einer Mastergleichung von Landauer [1] beschrieben werden. Es wird das stationäre Verhalten als Funktion des Steuerstromes (Pumpparameter) exakt berechnet und mit Approximationen (Fokker-Planck) verglichen. An diesem Beispiel eines Nichtgleichgewichtsphasenübergangs wird die Suszeptibilität $\chi(\tau)$ für eine Störung des Steuerstromes im kritischen und ausserkritischen Gebiet als Funktion der Systemgrösse mittels einer numerischen Eigenwertanalyse untersucht. Im kritischen Gebiet ergibt sich eine Quasientartung zum Eigenwert 0, die einer extrem langsamen Relaxation eines metastabilen Zustandes entspricht.

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Elektrische Leitfähigkeit

Mesures de densités de courant critique et de champs de pénétration dans des échantillons de $\text{Nb}_{0.75}\text{Pt}_{0.75}$ et $\text{Nb}_{0.75}\text{Au}_{0.125}\text{Pt}_{0.125}$ par méthode d'écrantage

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et

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Nous avons mesuré des densités de courant critique et champs de pénétration par méthode d'écrantage sur des échantillons massifs de $\text{Nb}_{0.75}\text{Pt}_{0.75}$ et $\text{Nb}_{0.75}\text{Au}_{0.125}\text{Pt}_{0.125}$. Nous avons obtenu respectivement pour les densités de courant critique $2,5 \cdot 10^4$ A/cm² et $8,6 \cdot 10^4$ A/cm² à 4,2°K sans champs magnétique et pour les champs de pénétration 1,2 kOe et 2,8 kOe pour une épaisseur de 1,5 mm. Les valeurs ainsi trouvées, sont comparées avec celles d'autres supraconducteurs massifs de type A 15.

Instabilités et supraconductivité dans les phases rhomboédriques $\text{Cu}_x\text{Mo}_3\text{S}_4$ et $\text{Cu}_x\text{Mo}_3\text{Se}_4$.

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Les deux composés $\text{Cu}_x\text{Mo}_3\text{S}_4$ et $\text{Cu}_x\text{Mo}_3\text{Se}_4$, rhomboédriques à température ambiante, subissent des transformations de phases à des températures $T < 273\text{K}$.

Nous avons déterminé le domaine d'homogénéité de la phase rhomboédrique (phase de Chevrel) dans ces deux systèmes à l'aide d'une étude micrographique effectuée sur des échantillons fondus sous plusieurs atmosphères d'Argon. Les transformations de phase pour différentes compositions x de Cu ont été observées à l'aide de la résistivité électrique et par diffractométrie aux rayons X à basse température. Différents éléments ont été substitués au Cu: leur influence sur la température de transformation et sur la température supraconductrice est étudiée dans ce travail.

Relation entre le domaine d'homogénéité de la phase A15 et la température supraconductrice du système Nb–Ga

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Nous présentons le diagramme de phase du système Nb–Ga, étudié par micrographies, microsonde, microdureté, diffraction par $R \cdot X$ et ATD Haute Température (2100°C).

A l'aide d'échantillons trempés par jet d'argon nous prouvons que la déviation vers la stoechiométrie de la limite riche en Ga de la phase A15 est responsable de la forte variation de la température supraconductrice observée après différents traitements thermiques. ($9 < T_c < 18$ K).

La température supraconductrice maximum (20, 7 K) observée dans ce système ne peut être obtenue qu'après des recuits prolongés à basse température (650 K).

Temperature dependence of the electrical resistivity of Cu–Ni alloys

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The resistivity data [1, 2, 3, 4] of $\text{Cu}_x\text{Ni}_{1-x}$ alloys ($0.28 \leq x \leq 1.0$) is analysed in terms of a (decreasing) T dependent impurity resistivity and an (increasing) ideal resistivity. As indicated by Ahmad and Greig [2], this reproduces the resistance minimum observed at low and high T . However, at low T , in alloys near the critical concentration for ferromagnetism, there is an additional contribution which is reminiscent of the magnetic resistivity calculated by Levin and Mills [5]. This latter contribution is not present in the $\text{Pd}_{69}\text{Ag}_{40}$ alloy [2].

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Résistivité électrique des alliages de Ni–Cu à concentration critique

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Nous présentons des mesures de résistivité électrique en fonction de la température pour des alliages de Cu–Ni de concentration 52, 53, 55, 57% at. de Cu soumis à différents traitements thermiques.

Nous remarquons que le recuit à température standard de 900–1000°C a un faible effet sur la résistivité, par contre nous observons une grande influence de recuit à 350°C sur la résistivité des alliages faiblement ferromagnétiques.

Untersuchung des elektronischen Phasenüberganges in $\text{Sm}_{1-x}\text{La}_x\text{S}$

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Der druckinduzierte Halbleiter-Metall Übergang in $\text{Sm}_{1-x}\text{La}_x\text{S}$ wurde mittels Reflexions- und Kompressionsmessungen untersucht. Der Radius des 3-wertigen La-Ions liegt zwischen den Radien des 2- und des 3-wertigen Samariums. Der 'chemische Druck' des Lanthans kann den Übergang zur metallischen Phase deshalb nicht induzieren, d.h. den kritischen Druck nicht auf Null reduzieren. Stattdessen beobachtet man mit zunehmender Konzentration eine abnehmende Druckhysterese und, ab ca. 30%, einen kontinuierlichen Übergang. Aufgrund von Kompressions- und Reflexionsmessungen lässt sich auf einen kritischen Punkt bei $x \simeq 0.28$ schliessen.

Falicov-Kimbal Modell für $\text{Sm}_{1-x}\text{R}_x\text{S}$ ($R =$ Seltene Erde oder Übergangsmetall)

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Der 'Halbleiter-Metall Übergang' in Verbindungen $\text{Sm}_{1-x}\text{R}_x\text{S}$ wurden in letzter Zeit eingehend untersucht. Er kann in der Mean-Field Näherung des Falicov-Kimbal Modells halbquantitativ beschrieben werden [1, 2].

In dieses Model habe ich die Konzentration x als neue Variable eingeführt, um einen besseren Einblick in das Verhalten von SmS-Legierungen zu gewinnen. Die Resultate sind in befriedigender Übereinstimmung mit experimentellen Ergebnissen. Insbesondere finde ich eine kritische Konzentration $x_{cr} \simeq 0.28$, unabhängig von der Art des Ions R , oberhalb der der Übergang kontinuierlich wird. Wenn das Material jedoch bei $x < x_{cr}$ kollabiert, ist p_{cr} negativ und der kritische Punkt nicht beobachtbar.

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Une anomalie dans la magnétorésistance de l'étain ultra-pur

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Des mesures de magnétorésistance effectuées sur des fils polycristallins d'étain ultrapur (rapport de résistivité résiduelle 10^4 à 5×10^4) d'un diamètre de 2 mm ont montré que la règle de Kohler était en défaut entre 2 et 4.2 K, pour des valeurs du champ magnétique comprises entre $H_c(T)$, le champ critique supraconducteur, et 400 Oe [1]. Les écarts par rapport à la règle de Kohler étaient plus importants en champ transversal qu'en champ longitudinal et augmentaient avec la pureté de l'échantillon.

Poursuivant ces études, nous avons étendu le domaine de mesure jusqu'à 20 kOe et à d'autres valeurs du diamètre de l'échantillon. Les résultats antérieurs se sont trouvés confirmés. Mais nous avons remarqué en outre que le *signe* de l'écart (par rapport à la règle de Kohler) était anormal; dans un diagramme de Kohler, la courbe de magnétorésistance transversale isotherme mesurée à 4.2 K est située *au-dessous* de celle qu'on obtient à 3.8 K. D'ordinaire, c'est le comportement inverse que l'on obtient lorsque la température est modifiée [2]. Une anomalie analogue est observée dans l'indium polycristallin entre 7 et 8 K [3]. Nous avançons différentes hypothèses pour tenter d'expliquer cette anomalie de magnétorésistance.

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Etat intermédiaire: étude de la résistance thermique d'une frontière de phase

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Après avoir mesuré la conductivité thermique 'anormalement' faible [1] d'une structure lamellaire statique perpendiculaire au courant de chaleur, contrôlée par une méthode magnéto-optique [2, 3], nous avons déduit de ces mesures (effectuées entre 1.5 et 4.2 K sur un échantillon de plomb de haute pureté) la résistance thermique d'une frontière entre phase normale et phase supraconductrice. Il ressort de ces

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mesures préliminaires que les réflexions des quasiparticules par les frontières de phase [4] sont la partie principale de cette résistance pour $T \gtrsim \frac{1}{2}T_c$, mais qu'à plus basse température ($T \lesssim \frac{1}{3}T_c$), la limitation du libre parcours moyen des phonons dans l'état supraconducteur [5, 6] joue aussi un rôle important. Nous présentons un premier essai de séparation de ces deux contributions.

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Absorptivité de films de niobium dans l'infra-rouge lointain sous l'effet conjugué d'un champ magnétique extérieur perpendiculaire et de la proximité d'un film de cuivre

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L'influence d'un champ magnétique perpendiculaire sur le spectre d'absorption d'un film supraconducteur dans l'infra-rouge lointain est connue, de même que celle d'un film de métal normal en contact étroit avec le premier film (effet de proximité). On peut décrire ces deux effets au moyen de la même théorie d'Abrikosov et Gorkov ('pair-breaking effects'). Dans le but de vérifier leur additivité (au niveau du paramètre Γ de la théorie), nous avons mesuré les spectres d'absorption de films doubles Nb – Cu dans un champ magnétique perpendiculaire.

Direct observation of the perturbation of the flow of the intermediate state structure in pure indium by helicons

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The magnetic structure of type-I superconductors in the intermediate state moves in the presence of an electric current. In very pure uncompensated metals, this flow is along the direction of the current when the Hall angle φ approximates $\pi/2$. If the structure moves past a static magnetic field perturbation, helicon waves will be excited in the specimen. As the damping of these excitations is weak ($\propto 1/\tan \varphi$), they will propagate across the sample and influence the configuration of the moving structure. We report here direct observation of this effect. The intermediate state flow pattern observed is similar to that produced at the surface of a flowing liquid by an obstacle lying well below the surface of the liquid.

Optische Eigenschaften und diffuse Röntgenstreuung des Supraionenleiters (C₅H₅NH) Ag₅I₆

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Im Supraionenleiter Pyridinium-Silberiodid, (C₅H₅NH)Ag₅I₆, beobachtet man einen von der Temperatur kontinuierlich abhängigen Unordnungsgrad der Ag⁺-Ionen (praktisch vollständige Ordnung bei -40°C bis fast völlige Unordnung bei 125°C). Um Informationen über die Dynamik der Leitungslionen Ag⁺ zu gewinnen, haben wir die Reflexion im fernen Infrarot für parallel und senkrecht zur optischen Achse polarisiertes Licht bei verschiedenen Temperaturen gemessen. Von besonderem Interesse ist eine sehr tieffrequente Eigenschwingung mit einer Energie von ca. 2 meV. Die in der Röntgenstreuung beobachteten diffusen Ringe weisen ebenfalls auf sehr tieffrequente optische Modes hin, bei welchen die Ag⁺- gegen die I⁻-Ionen schwingen. Wir diskutieren die Bedeutung dieser Modes für den Mechanismus der Ionenleitfähigkeit.

Frequency dependent conductivity for *N* interacting Brownian particle memory function approach

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A method is presented to compute the current-current correlation function and therefore the conductivity for a system of *N* particles described by *N* coupled Langevin equations in the memory function approximation. In particular the attention is focused on the connection of the high frequency (oscillatory) regime to the low frequency (diffusive) one. The case of the lattice with two different atoms is presented as an example and its relation to the experimental results on simple superionic conductors is discussed.

Cluster expansion for many particle many site systems. Application to the solid electrolyte (C₅H₅NH) Ag₅I₆

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The statistical mechanics of a system with different types of interacting particles which can occupy lattice sites of various kinds is studied. The method is based on a generalized cluster expansion. In lowest order it is identical to the quasi-chemical method. The specific heat of (C₅H₅NH)Ag₅I₆ is computed and compared with experiments.

Korrelierte Teilchensprünge in Supraionenleitern

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Die Konzentration von beweglichen Ionen in Supraionenleitern ist genügend hoch, dass aufeinanderfolgende Teilchensprünge korreliert sein können. Wir geben einfache Beziehungen für den Einfluss der Korrelation auf experimentelle Grössen und zeigen, dass im allgemeinen die interessierenden Korrelationsfunktionen leicht aus dem tieffrequenten Teil von $\sigma(\omega)$ extrahiert werden können. Wir diskutieren exakt lösbare Modellbeispiele und experimentelle Resultate an verschiedenen Ionenleitern.

Strominstabilität und wandernde Dipoldomänen aufgrund von Bragg-Streuung

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Ein neuer Mechanismus, der in Eintal-Halbleitern zu einer Strominstabilität führt, wird vorgeschlagen. Bragg-Streuung heisser Elektronen am Brillouin-Zonenrand führt zu einem Impulsverlust der Ladungsträger und makroskopisch zu negativer differentieller Leitfähigkeit. Ein einfaches Modell wurde konstruiert, welches die Instabilität des homogenen stationären Stromzustandes bei einem kritischen Feld E_c zeigt, wo die dielektrische Relaxationsfrequenz weich wird und die Debyelänge divergiert. Die Hochfeldzustände sind wandernde Dipol-Domänen. Die Möglichkeit diesen Mechanismus zu beobachten wird diskutiert.

Magnetismus

Magnetische Phasenübergänge von CeSb

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Der metallische Metamagnet CeSb besitzt wegen anisotropen Austauschwechselwirkungen und Kristallfeldeffekten interessante magnetische Eigenschaften. Es wird über neue Einkristall-Neutronenbeugungsuntersuchungen an CeSb in äusseren Magnetfeldern und als Funktion der Temperatur berichtet. Abgesehen von Domäneneffekten wurden in schwachen Magnetfeldern ferrimagnetische Strukturen analog zu antiferromagnetischen Phasen bei $H = 0$ gefunden. Eine Konfiguration $++-$ ist bei mittleren Feldstärken und 4.2 K stabil.

Magnetische Suszeptibilität und Valenzfluktuationen in Sm_3S_4

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Die gemischt valente Verbindung Sm_3S_4 zeigt eine temperaturaktivierte Valenzfluktuation zwischen Sm^{2+} und Sm^{3+} . Die magnetische Suszeptibilität von Einkristallen wurde zwischen 1.8 K und 300 K untersucht und ein qualitativ und quantitativ anderes Verhalten gefunden als von Escorne *et al.* [1] an polykristallinen Proben beschrieben.

Unterhalb etwa 70 K lässt sich die Suszeptibilität durch $\chi = \chi_{\text{Sm}^{2+}} + 2\chi_{\text{Sm}^{3+}}$ erklären, entsprechend einem Van Vleck Term von Sm^{2+} (${}^7\text{F}_0$) und einem Curie Weiss und Van Vleck Term des kristallfeldaufgespaltenen Sm^{3+} (${}^6\text{H}_{5/2}$). Oberhalb etwa 100 K wird die Suszeptibilität von Sm^{3+} gequencht und sie benimmt sich ebenfalls wie ein $J = 0$ System. Weil die Fluktuationsrate und damit die Lebensdauer einer Konfiguration temperaturabhängig ist, entspricht 100 K etwa 10^{-6} sek. Es folgt, dass für kürzere Lebensdauern einer Konfiguration, Sm^{3+} nicht mehr voll zur Suszeptibilität beitragen kann. Sm_3S_4 kann damit zu einer Modellsubstanz auch für die Suszeptibilität von metallischem SmS werden, ist dort doch die Lebensdauer immer etwa 10^{-12} sek und folglich $J_{\text{Sm}^{3+}}$ etwa gleich null.

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Study of valence fluctuations in Eu_3S_4

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By means of inelastic light scattering and $5d-4f$ photoluminescence we have investigated the thermally activated valence fluctuations in pressed powder samples of Eu_3S_4 . The light scattering measurements reveal a phonon mode, $\nu_0 = 425 \text{ cm}^{-1}$, which is not infrared active and which can possibly be described by a vibration alone in the anion sublattice with the cation positions being fixed. This mode can be related to the frequency factor of electron hopping, and ν_0 fits well with an extrapolated value determined from Mössbauer measurements [1]. At 300 K the photoluminescence spectrum extends to the absorption edge (1.7 eV) and shifts to lower energies by about 0.25 eV upon cooling to 4.2 K. This anomalous feature can be attributed to the valence fluctuations and can be explained by comparing the luminescence lifetime, 10^{-6} sec, with the temperature dependent fluctuation rate. Hence, the temperature shift in the luminescence results from driving the fluctuation rate through the reference time scale.

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Elektronenstruktur ferromagnetischer Halbleiter

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Durch approximative Bestimmung von relevanten Spektraldichten wird der Einfluss der fd – Elektron-Ion – Kopplung auf die elektronischen Eigenschaften eines ferromagnetischen Halbleiters untersucht. Die magnetischen Rotverschiebungen in der Reihe der Eu – Chalkogenide erklären sich mit einer einzigen fd – Kopplungskonstanten als Funktion der Bandbreiten.

Propriétés magnétiques d'un alliage eutectique orienté contenant des fibres de $\text{Co}_{17}\text{Sm}_2$

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Un alliage eutectique à 44% en volume de fibres de $\text{Co}_{17}\text{Sm}_2$ a été mis en évidence dans le système Co—Sm—Sn. Après solidification orientée à des vitesses comprises entre 1.7 et $550 \mu\text{m s}^{-1}$, des fibres ayant des diamètres entre 1.6 et $0.14 \mu\text{m}$ ont été élaborées. Un champ coercitif de 2.3 kOe a été mesuré pour les fibres les plus minces. On comparera ces résultats avec ceux obtenus pour l'eutectique Bi—MnBi et avec le comportement des poudres de ces composés. Ces résultats peuvent être interprétés à l'aide d'un modèle simple montrant que la limite pour l'obtention d'un champ élevé est donnée:

- dans le cas des poudres, par la densité des défauts de surface induits lors du broyage,
- pour les eutectiques par la régularité et l'orientation des structures à haute vitesse de solidification.

Zentren und Sonden

Radioluminescence de CsI irradié avec des protons de 1 MeV

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Le rendement des émissions intrinsèques (recombinaisons radiatives $V_K + e^-$) observées dans CsI pur ou dopé en Na^+ sous irradiation de protons de 1 MeV diminue en fonction de la dose d'irradiation. Ce phénomène est expliqué par la présence de centres interstitiels halogènes. Ces derniers seraient créés à la suite de collisions proton-halogène, le site halogène vacant étant occupé finalement par l'implantation de l'ion hydrogène H^- (centre U) dont la formation est observée simultanément par absorption optique. Dans CsI:Na on observe l'apparition d'une nouvelle composante luminescente à $2,7 \text{ eV}$. Son rendement varie quadratiquement avec la densité d'irradiation protonique. On peut penser que cette émission est due à

la recombinaison de biexcitons. Différentes mesures montrent cependant qu'elle résulte probablement de la recombinaison d'excitons perturbés par la présence de centres V_K ou Na^0 .

Etude de la polarisation des luminescences intrinsèques dans le iodure de césium

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Une irradiation X de CsI:Na à 4 K crée des centres V_K (trous self-trappés) et Na^0 . Il est possible d'orienter quasi totalement les centres V_K parallèlement à leur axe de symétrie [100] par excitation en lumière polarisée dans la transition UV $^2\Sigma_u^+ \rightarrow ^2\Sigma_g^+$. Les deux luminescences intrinsèques à 290 et 338 nm (recombinaison $V_K + e^-$) apparaissent cependant partiellement polarisées perpendiculairement à l'axe de symétrie [100] des V_K . La polarisation de l'émission 338 nm est indépendante de la température tandis que celle à 290 nm en dépend fortement. Un modèle théorique sera présenté expliquant ces propriétés ainsi que les dépendances en température des intensités et des temps de déclin. On considère en particulier que le 'hot exciton' créé pendant la relaxation ($V_K + e^-$) peut se réorienter partiellement ($\sim 22\%$) avant d'atteindre deux états excitoniques luminescents de symétries différentes.

Zeeman-effect of free and iodine bound excitons in AgBr

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We have observed the splitting of free excitons and of iodine bound excitons in AgBr in magnetic fields up to 80 kG. The fourfold degenerate free excitons show an exchange splitting into two twofold states without field, while the magnetic field splits the residual degeneracies. The iodine bound exciton splits into 8 components and has no exchange splitting. Both systems are described by one set of 3 g -values. The properties of these systems will be discussed.

Annihilation von Triplett-Exzitonen in Naphthalin bei verschiedenen Exzitonendichten

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Wir haben die Wechselwirkung insbesondere die gegenseitige Annihilation von Triplett-Exzitonen in Naphthalin untersucht. Die Annihilationsrate ändert sich im Magnetfeld. Wir haben die Anisotropie dieser Änderung und ihre Abhängigkeit

von der Triplett-Exzitonendichte gemessen. Die Messresultate zeigen deutliche Abweichungen von der allgemein anerkannten Suna-Theorie (Phys. Rev. *B1*, 1716, 1969).

Verzögerte Fluoreszenz in dotierten Anthrazen-Kristallen bei verschiedenen Temperaturen und Anregungsbedingungen

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Die verzögerte Fluoreszenz in dotierten Molekülkristallen gibt Aufschluss über den Einfluss von Haftstellen auf die Uebertragung von Anregungsenergie zwischen Wirt- und Gastmolekülen in diesen Systemen. Das Verhältnis der verzögerten Fluoreszenz von Gast und Wirt ist in mit Tetrazen dotierten Anthrazenkristallen in Funktion der Temperatur gemessen worden. Die Anregung des Wirtes erfolgte bei zwei verschiedenen Wellenlängen innerhalb des Absorptionsspektrums des Wirtes, wobei ein grosser Unterschied im Energie-Uebertragungsverhältnis beobachtet wird.

Magnetfeldabhängigkeit der prompten und der verzögerten Fluoreszenz in Anthrazen bei tiefen Temperaturen

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Die Anisotropie der Fusion von Triplett-Exzitonen und der Fission von Singlett-Exzitonen in einem äusseren Magnetfeld werden für tiefe Temperaturen (10 K) untersucht. Aus den Messresultaten werden Rückschlüsse auf die Exziton-Kinetik wie auch auf den Annihilationsprozess selbst gezogen.

Annihilation de positrons: une nouvelle approche expérimentale

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Nous avons étudié l'annihilation de positrons dans un monocristal de cuivre à l'aide d'un système expérimental basé sur deux chambres proportionnelles multifils munies de convertisseurs gamma-électrons. L'originalité de cette expérience réside dans la possibilité de résoudre simultanément deux composantes de l'angle séparant les deux photons émis lors de l'annihilation. Les statistiques bidimensionnelles ainsi obtenues reflètent clairement la topologie de la surface de Fermi (SF) du cuivre. L'accord des résultats avec un calcul de la SF déterminée par l'effet de Haas van Alphen prouve l'efficacité de la méthode.

Détection optique de la RMN par l'émission et l'absorption de paires de centres F dans KCl

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Alors que la détection optique de paires 'courtes' de centres F (d_{FF} : 2–3 nm) est révélée par une *augmentation* de la luminescence lorsque la condition de résonance est remplie, il existe un autre mécanisme qui fait apparaître la RMN comme une *diminution* de la luminescence et de l'absorption dans la bande F. Cela se produit pour des cristaux moyennement dotés ($5 \cdot 10^{17}$ F/cm³), trempés et examinés à basse température (12 K; $f = 2,3$ MHz). Bien que la luminescence et l'absorption manifestent la forte dépendance du champ magnétique caractéristique des paires 'longues' (d_{FF} : 7–9 nm), le temps de relaxation de 10–20 ms implique un type de paire intermédiaire, pour lequel un modèle est développé.

Electronic properties of pure and Cu-doped KTaO_3

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Incipient ferroelectrics are materials characterized by a dielectric constant which increases rapidly with decreasing temperature in a manner consistent with a transition to a ferroelectric state near 0 K. Such a transition, however, is not observed at ambient pressure [1]. The purpose of the present investigation is to gain insight into the electronic properties of the incipient ferroelectric potassium tantalate (KTaO_3) and into the role played by impurities in determining the behavior of this substance. This work has been carried out using the techniques of electron paramagnetic resonance to monitor the characteristics of deliberately added paramagnetic impurities, and of acoustic flexure resonance to determine the space-charge density in both pure and doped KTaO_3 .

In order to exercise the maximum possible control over the sample history and impurity content of the material, a simple technique for the growth of KTaO_3 single crystals was developed [2]. The crystal growth was carried out in a platinum crucible containing high-purity KTaO_3 to which 15 to 20% by weight K_2CO_3 had been added. The desired dopant was added in the oxide form, and the resulting mixture was heated to approximately 1400°C. A programmable temperature control system was then used to cool the material slowly to room temperature. Using this procedure, it was possible to produce single crystals as large as $1.0 \times 1.0 \times 0.2$ cm. The best crystal-growth results were obtained with a 48-hour cooling cycle, and although it was possible to grow single crystals with significantly shorter cooling times, the resulting material appeared to be highly strained and was characterized by different electronic properties. Single crystals which were intentionally doped with iron-group im-

purities were invariably transparent and were sometimes colored depending on the specific impurity which had been added. Attempts to grow transparent undoped single crystals resulted in the production of opaque blue-white material. The dielectric properties of this undoped material differed markedly from those of doped crystals and, additionally, no EPR spectra were detected for these 'pure' samples. Single crystals of KTaO_3 to which 0.5 to 2.0% by weight of CuO had been added prior to growth, were characterized by a light-green color, and the intense electron paramagnetic resonance spectrum shown in Figure 1 was observed at $T = 77 \text{ K}$ when the applied magnetic field was oriented parallel to a $[100]$ crystallographic direction. This spectrum was obtained using a sample doped with CuO which had been isotopically enriched to contain 99.7% ^{65}Cu . The full EPR spectrum due to Cu^{2+} consisted of three axially-symmetric spectra whose principal symmetry axes lay along the cubic fourfold axes of the KTaO_3 host crystal. As shown in Figure 1, for $\mathbf{H} \parallel [100]$, the spectrum consists of four major transitions which exhibit a strong resolved superhyperfine structure. These four groups of superhyperfine lines are due to the $I = 3/2$ nuclear spin of ^{65}Cu and they represent the spectrum associated with an orientation of the applied magnetic field parallel to the principal symmetry axis of one of the three tetragonally-symmetric spectra. The ^{65}Cu hyperfine structure is not resolved on the transitions associated with the orientation of the applied field perpendicular to the principal axes of the two remaining spectra, and these superimposed spectra interfere with the superhyperfine structure of the fourth (counting from low field) hyperfine transition of the spectrum at the parallel orientation. Divalent copper has the $3d^9$ electronic configuration and if the ion were located at the sixfold oxygen-coordinated site, the resulting 2E orbital electronic state could exhibit a static Jahn-Teller effect [3] resulting in the observation of three tetragonally-symmetric spectra with principal axis systems identical to those found here. If the observed spectra were associated with a static Jahn-Teller effect, however, one should observe an averaging of the tetragonal components with increasing temperature to produce an isotropic spectrum. For Cu^{2+} in KTaO_3 , no averaging was found and,

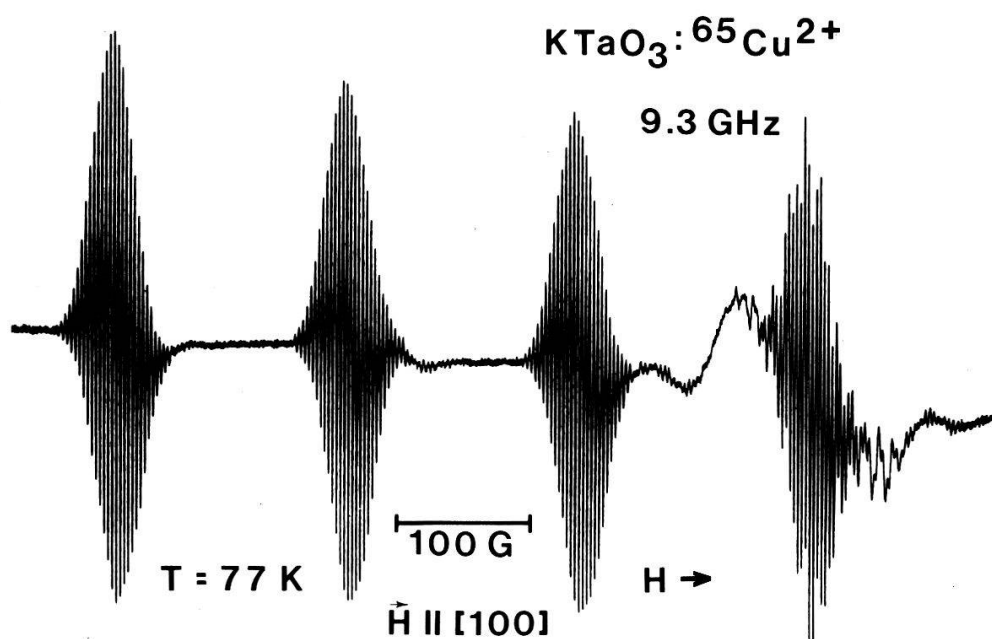


Figure 1
EPR spectrum of Cu^{2+} in KTaO_3 with $\mathbf{H} \parallel [100]$.

in fact, the tetragonally-symmetric spectra could still be observed at room temperature. The tetragonal symmetry is, therefore, apparently due either to a defect associated with the Cu^{2+} or to the incorporation of the ion in an interstitial site where the effective symmetry of the crystal field is tetragonal and is not due to the Jahn-Teller effect. The observed Cu^{2+} spectra are described by the usual axially-symmetric spin-Hamiltonian with $g_{\parallel} = 2.228 \pm 0.002$, $g_{\perp} = 2.056 \pm 0.005$, $A_{\parallel} = (173 \pm 2) \times 10^{-4} \text{ cm}^{-1}$, and $A_{\perp} = (45 \pm 3) \times 10^{-4} \text{ cm}^{-1}$. The superhyperfine splitting with $\text{H} \parallel [100]$ was 3.0 Gauss.

The dielectric constant of slow-cooled transparent samples was similar to that measured by Wemple [4], i.e., its maximum value was 4500 below 4 K. In all samples, the dielectric constant was reduced [5] upon transient application of a dc field (poling); in addition, such samples may be excited into flexure resonance [6] by means of an ac electric field. These observations are indicative of space charge introduced irreversibly by a dc electric field. The maximum charge density which can be introduced in this way is of the order of 50 C/m^3 corresponding to 3×10^{14} unit charges per cm^3 . This low space-charge density plus the fact that it depends on field treatment but not on intentional doping excludes the possibility of Cu^{2+} being the source of space charge.

The blue samples which were not intentionally doped had a high dielectric constant $\epsilon \sim 20,000$ from 0 to 400 K with a shallow maximum at 350 K. We presume that these samples have a number of oxygen vacancies sufficiently large to impair their growth habit. They are semiconducting, brittle, small and unsuitable for flexure resonances.

The cumulative results based on the currently available experimental evidence may be summarized as follows: single crystals of KTaO_3 doped with copper and other iron-group impurities are characterized by similar electronic properties. These results are in agreement with those reported previously [4]. Attempts to grow KTaO_3 single crystals without intentionally doping the material resulted in samples whose electronic properties were significantly different relative to the doped crystals. The present results are consistent with the presence of oxygen vacancies in both the doped and pure materials, however an interaction involving the impurities is apparently responsible for the observed difference.

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ANGEWANDTE PHYSIK UND ANDERE DOMÄNEN

Coefficient de dissociation d'oxygène dans les plasmas a basse pression

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La concentration des atomes dans les plasmas d'oxygène a été examinée en dépendance des paramètres de la décharge notamment du courant de décharge et de la pression du gaz. On a utilisé la méthode calorimétrique avec une sonde mobile. Les mesures ont été réalisées dans deux formes *T* et *H* des plasmas d'oxygène. Conformément aux résultats sont discutées les réactions dans les deux formes.

Mesures de l'effet de peau dans un plasma¹⁾

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On présente ici quelques résultats de mesures d'atténuation d'une onde électromagnétique dans un plasma d'argon. L'amplitude du champ décroît de façon monotone du bord du plasma jusqu'au centre lorsque la relation entre la densité de courant et le champ électrique est locale, c'est-à-dire [1]:

$$\lambda \ll \delta_0 \quad \text{ou} \quad d \ll \delta_0 \quad (1)$$

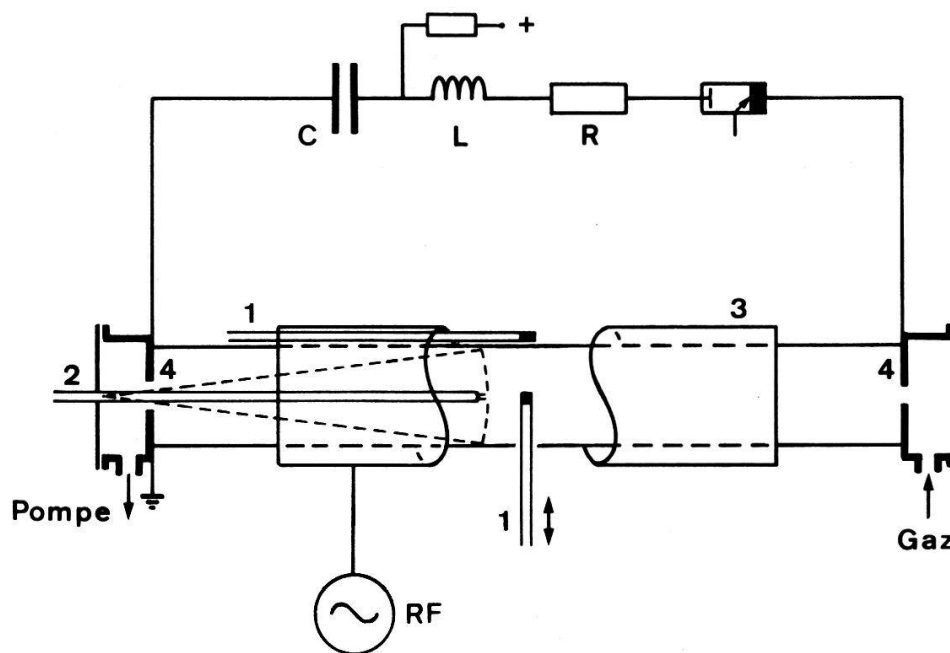


Figure 1

Schéma du dispositif expérimental:

- | | |
|-----------------------------|----------------|
| 1 = sondes magnétiques | 3 = bobine |
| 2 = sonde électrique double | 4 = électrodes |

¹⁾ Travail subventionné par le Fonds National Suisse de la Recherche Scientifique.

avec $\delta_0 = \delta \cos \alpha/2$, $\text{tg } \alpha = \nu/\omega$, $d = u/\omega$, $u = (2 \text{ kT/m})^{1/2}$, $\delta^{-1} = \text{Re}(i\omega\mu_0\sigma)^{1/2}$: profondeur de pénétration classique.²⁾ ν est la fréquence de collisions des électrons, et ω la fréquence angulaire de l'onde; d représente la distance parcourue durant une période divisée par 2π , et λ le libre parcours moyen. Nos mesures ont été effectuées dans le cas où aucune des deux conditions (1) n'est satisfaite (effet de peau anormal).

La Figure 1 montre le dispositif expérimental: le plasma est produit dans un cylindre de verre de 1 m de longueur et 9,5 cm de diamètre intérieur, par décharge de condensateurs (77,5 μF , 5kV), sur une bobine d'induction (2,2 mH) et une résistance (20,1 Ω). La pression de remplissage d'argon est 10 mTorr. Les paramètres du plasma ainsi formé ont été mesurés au moyen d'une sonde électrique double [2, 3]. Un émetteur de fréquence radio, relié à une bobine de cuivre (1 spire, 60 cm de longueur) produit un champ magnétique axial de faible amplitude, mesuré à l'aide de deux sondes magnétiques. L'une d'elles, mobile perpendiculairement à l'axe du cylindre, nous a permis de déterminer l'amplitude relative du champ en fonction de la distance à l'axe.

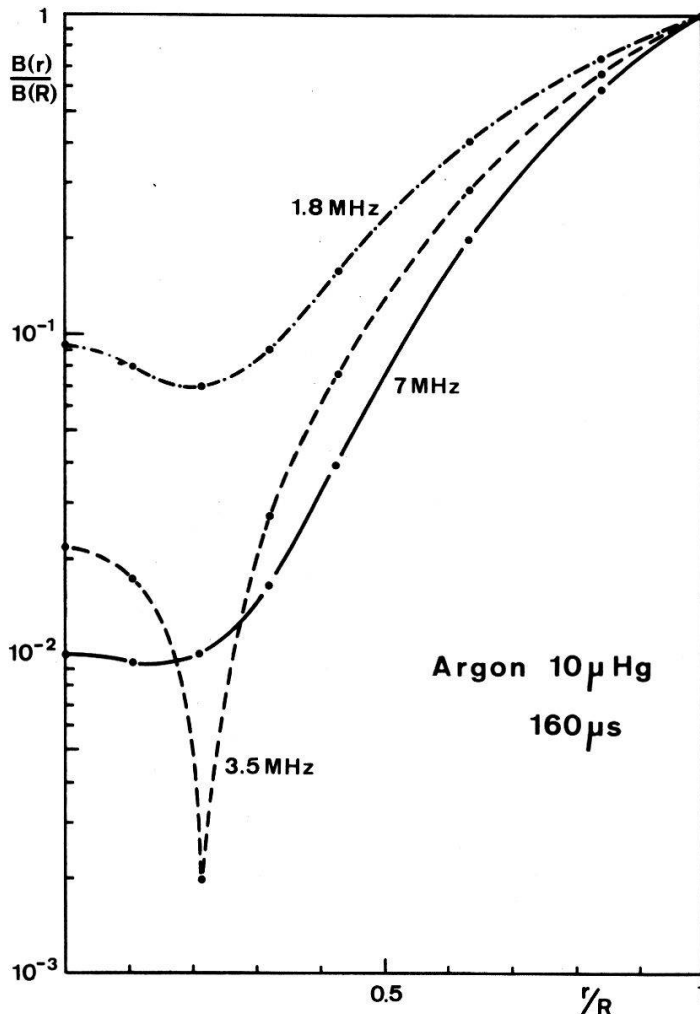


Figure 2

Amplitude du champ magnétique en fonction de la position radiale de la sonde pour différentes fréquences $\omega/2\pi$.

R = rayon du cylindre

160 μs après le début de la décharge

$n_e = 5 \times 10^{12} \text{ cm}^{-3}$, $T_e = 2 \text{ eV}$.

²⁾ Unités MKS.

Sur la Figure 2, sont représentés des résultats pour un plasma de densité $n_e = 5 \times 10^{12} \text{ cm}^{-3}$, et de température $T_e = 2 \text{ eV}$, résultats qui mettent en évidence l'existence d'un minimum d'amplitude situé environ à 1 cm de l'axe. Pour un plasma donné, il existe un domaine de fréquences pour lequel l'effet anormal est le plus prononcé (environ 3,5 MHz pour nos conditions), car soit en augmentant, soit en diminuant la fréquence, l'une ou l'autre des inégalités (1) tend à être satisfaite et la pénétration redevient normale.

Des mesures plus complètes d'amplitude et de déphasage, ainsi que des comparaisons avec les résultats théoriques paraîtront prochainement.

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Alfvén wave heating experiment of a high beta plasma

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Abstract. A helical $m = 1$ magnetic coil sets in motion the global kink mode. The frequency is chosen so that the discrete kink mode should be inside the continuous Alfvén spectrum. The source term of the excitation is proportional to the density gradient. This fact allows the kink to transmit energy to the torsional Alfvén waves, thereby building up a localized resonance of the velocity and current fields, with associated resonant absorption. Evidence of strong power coupling follows from energy balance. Temperature increase is measured by diamagnetic probe.

1. Introduction

During the last years, a growing interest has been devoted to a new type of additional heating by the excitation of low frequency MHD Alfvén waves. Theoretical investigations on the growth of the resonance surface and on the absorption mechanism have been made [1, 2, 3, 4]. In 1973 an experiment showed already that the damping of the kink mode can lead to plasma heating [5]. More recently experiments on stellarators have been undertaken [6, 7, 10]. The aim of our experiment is to clearly define the excitation and to vary the frequency and wavelength. Heating rates are measured for constant initial plasma conditions.

The measurements reported here concern primarily a theta pinch but some results on screw pinch discharges are also given.

2. Experimental apparatus

The basic parameters of the screw pinch are given below

Length of coil	142 cm
Inner diameter of coil	9 cm
Inner diameter of quartz discharge tube	5.2 cm
Maximum magnetic field	1.6 T
Quarter period (followed by crow-bar)	3.8 μs

Typical plasma parameters obtained at 16 mTorr D_2 filling pressure are:

Mean density \bar{n}_e	$1.3 \times 10^{16} \text{ cm}^{-3}$
Temperature $(KT_e + KT_i)/2$	48 eV
Mean beta	0.25

Mass collection after the implosion and following oscillations is 90%. Compression ratio is 13.8 (with plasma radius $r_p = 7 \text{ mm}$).

The excitation of the kink mode is produced by helical windings which give a perturbation of the form

$$\zeta = \zeta_0 e^{i(\theta + h_f z)} \cdot \cos \omega t \quad h_f \equiv -2\pi/\lambda_f$$

The windings are placed on the quartz tube at $r = 3 \text{ cm}$, covering sectors of 90° . Each coil is made of 8 turns connected in parallel. They are wound in the opposite sense to the screw pinch field, with a wavelength $\lambda_f = -57 \text{ cm}$. They are interrupted in the midplane to allow for diagnostics.

The $m = 1$ bank is composed of a 96 nF fast capacitor charged to 45 kV which oscillates at 0.9 MHz. The frequency can be varied by inserting capacitors or inductors in series with the $m = 1$ windings. A shunt spark gap provides for the possibility of current interruption after the desired number of periods.

To design the experiment one has to consider the low-frequency MHD spectrum [8]. The aim is to choose parameters (ω, h_f) so that the kink should be inside the Alfvén continuum. In the case of theta pinch, this condition is always fulfilled. The kink is described by

$$\omega_k^2 = v_A^2 h_f^2 (2 - \bar{\beta})$$

with $v_A^2 = B_z^2 / \mu_0 \bar{\rho}$, $\bar{\rho}$ and $\bar{\beta}$ stand for the mean values. The Alfvén spectrum is

$$\omega_A^2(r) = \frac{B_z^2(r)}{\mu_0 \rho(r)} \cdot h_f^2$$

The resonant layer is found where $\omega_k = \omega_A(r)$. For a Gaussian profile, as in the case of our pinch, the resonant layer is located at $r = 5.7 \text{ mm}$. It is possible to place the resonant surface closer to the center by lowering the applied frequency. The radius is independent on the excitation wavelength.

3. Diagnostics and results

The $m = 1$ current is started 1.5 μs after the main field, when the $m = 0$ bumps following the implosion have damped out. The stereoscopic streak photographs give the general evolution of the discharge. They show that kink-motion occurs only in the excited degree of freedom. The photographs reported here are made below, at and above the resonance absorption peak found at 0.92 MHz.

They are made in two different operation modes:

- continuous wave operation: below the resonance the plasma motion is important as it is less damped (Fig. 1). Above the resonance the damping quenches the kink-motion (Fig. 3).
- two-periods excitation: at resonance and before the excitation stops (Fig. 2), kink-motion is more strongly absorbed than at lower frequency, allowing for smaller displacement. After the excitation stops, the unforced kink shows nearly critical damping.

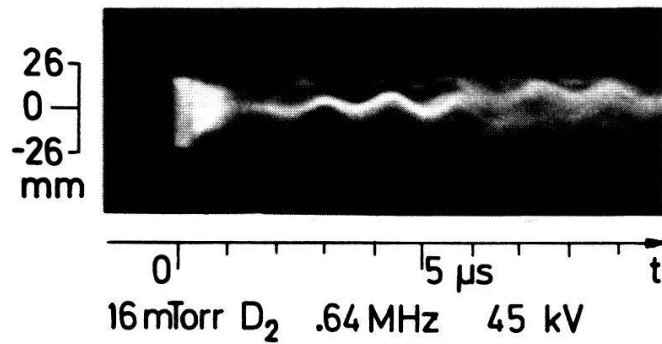


Figure 1

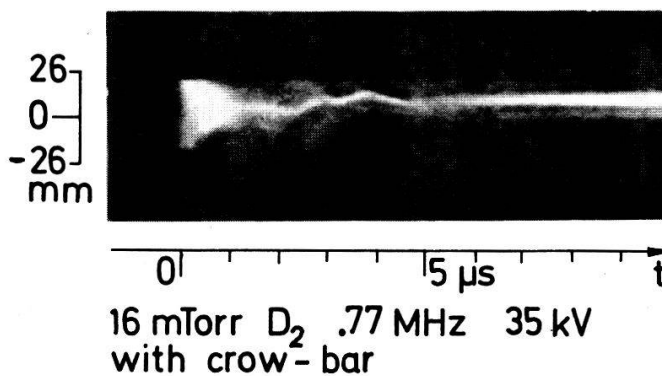


Figure 2

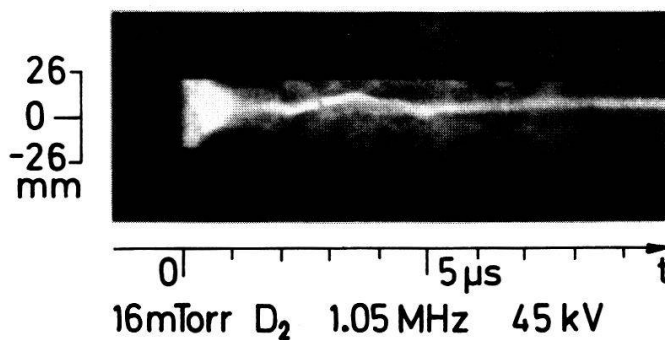


Figure 3

The diamagnetic probe (excluded flux measurements) gives the mean value of beta

$$\bar{\beta} = \frac{4\mu_0 \bar{n}KT}{B_{zv}^2}$$

The evolution in time is shown in Figure 4. The mean value of the temperature KT/e follows from the pressure balance (Fig. 5). Knowledge of the plasma radius is necessary and affects the value of $\bar{\beta}$, whilst KT/e is unnoticeably affected by the radius in the range of parameters used. The plasma radius $r_p = 7$ mm is determined from the luminosity profile and the relation $I(0)/I(r_p) = e^2$ in the case of a Gaussian density

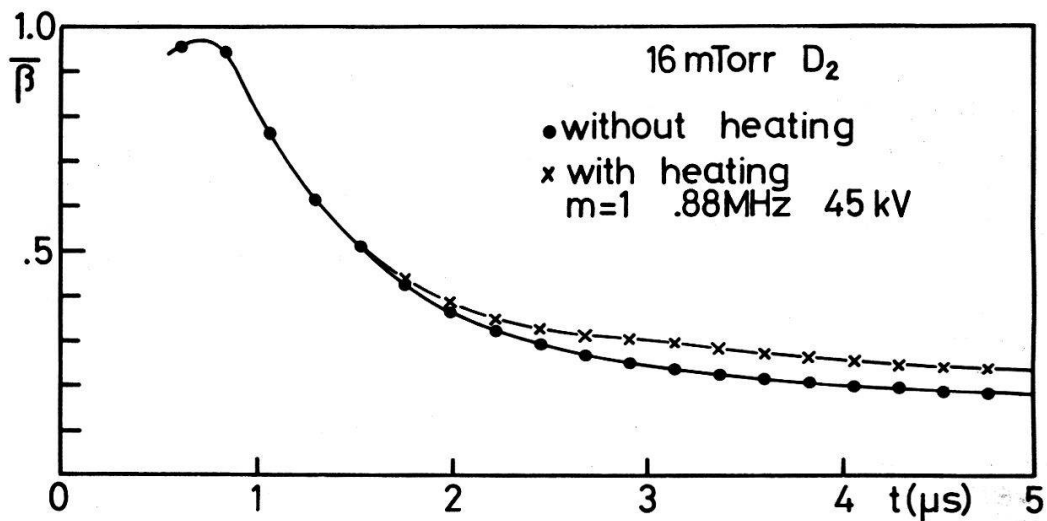


Figure 4

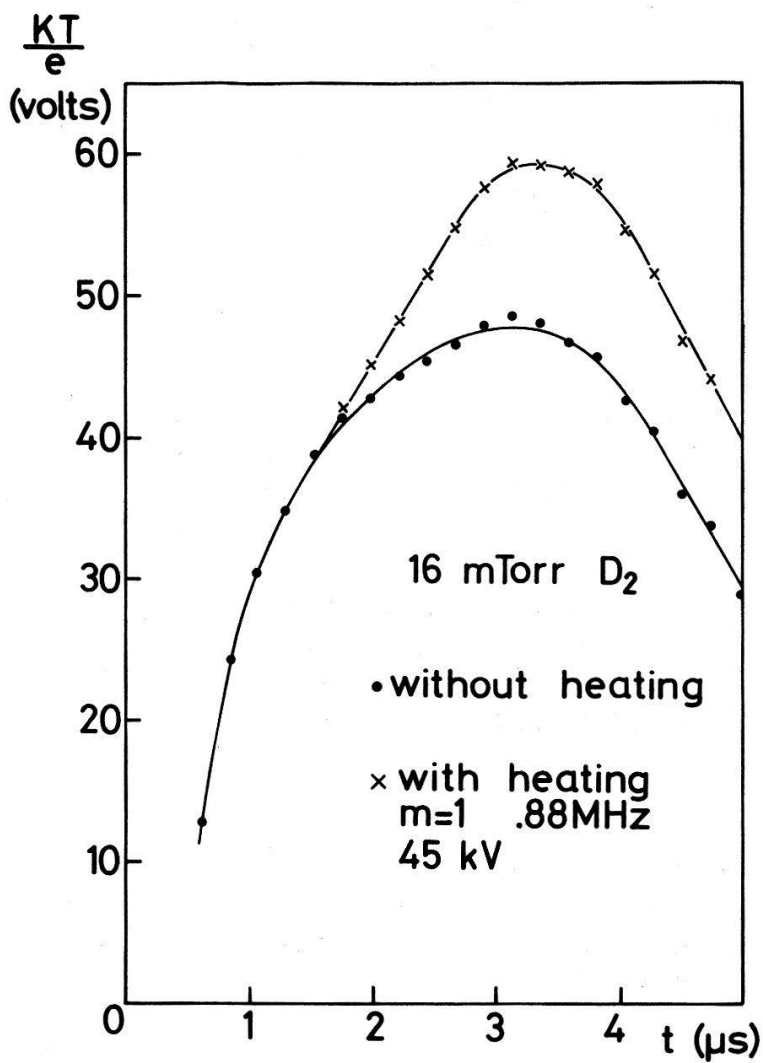


Figure 5

profile. This radius is consistent with the onset of the Kruskal-Shafranov instability limit at $I_z = 1800$ A.

The $\bar{\beta}$ is shown in Figure 4 with and without heating. After $1.5 \mu\text{s}$ the heating has increased the $\bar{\beta}$ measured at the same time by a factor of 1.24. Nevertheless the $\bar{\beta}$ is decreasing, even when heated, due to thermal conduction to the ends and to radiation.

Taking into account these loss processes, the thermal power W transmitted to the plasma is found by introducing a source term proportional to the applied power in the differential equation governing the decrease of $\bar{\beta}$. The resulting thermal power W in Figure 6 shows a resonance peak. The error bars represent the mean standard deviation.

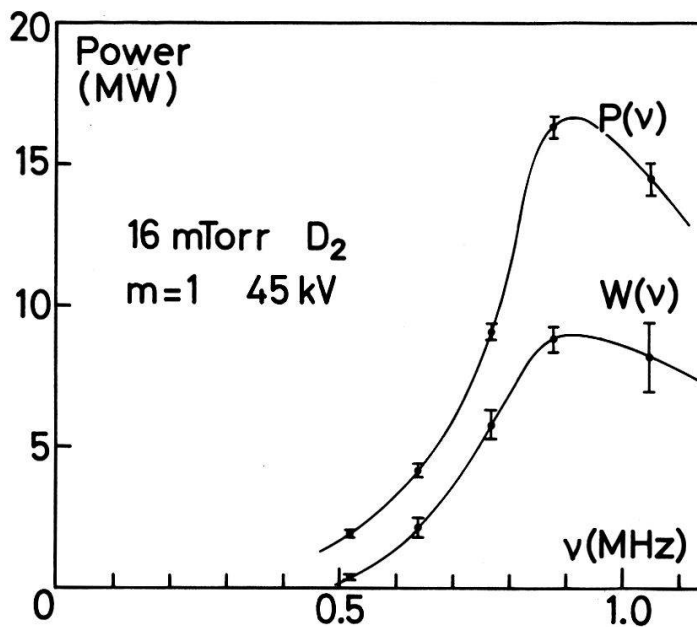


Figure 6

The power P delivered to the plasma by the $m = 1$ field is deduced from the damping of the current in the excitation bank, after subtracting the ohmic power lost in the circuit. P is found by comparing signals with and without plasma in the discharge tube. $P(\nu)$, shown on Figure 6 confirms the resonance peak obtained with the absorbed thermal power $W(\nu)$.

At resonance, the bank power is 30.5 MW, the power actually delivered to the plasma is 17 MW. The thermal power measured diamagnetically is 9 MW. That means that 53% of the excitation power is thermalized. In our experiment, 30% of the capacitor power contributed to heating the plasma.

4. Discussion

The heating rate may be increased by ohmic dissipation of the mirror currents set up by the $m = 1$ conductors. An estimate of this contribution is made here. If the current is dissipated in a skin layer given by the Spitzer resistivity at the plasma radius, the ohmic thermal power contribution is

$$W_{\Omega} = 10^{-5} r_p \nu^{5/2} T^{-3/4} \quad \text{mks, } T \text{ in eV}$$

At the resonance, this gives 2.8 MW. This should nevertheless be an overestimation,

as the resonance peak is not screened by this $v^{5/2}$ law, and as this contribution would be bigger than the thermal power measured by the diamagnetic coil at low frequency. It follows that the ohmic contribution represents at most 20% of the thermal power.

One part of the difference between $P(v)$ and $W(v)$ in Figure 6 is explained by the following: at 16 mTorr D_2 , the mass collection represents 90% of the total mass [9]. The ionisation energy per ion pair is 25 eV, which represents 2 Joules. The additional part may be explained by enhanced thermal conduction due to the crossing of field lines with the quartz tube.

It has been noticed that the heating rate is increased when the $m = 1$ bank starts before the end of the $m = 0$ implosion bumps. The phenomenon is complex as the plasma experiences strong acceleration in this phase and the radial density profile is quite different from the equilibrium, allowing for dips in the core, which modifies the Alfvén spectrum.

All the measurements reported here are made in the theta pinch mode, but measurements for a screw pinch below the Kruskal-Shafranov instability limit indicate that the absorption is slightly higher.

It has been noticed that the $m = 1$ oscillating field has an important stabilizing effect on the unstable screw pinch discharge.

The nature of the phenomenon of resonant absorption has been shown on a pinch discharge. The power coupling of 53% makes this method very attractive for larger scale heating of a plasma.

This work was supported by the Swiss National Science Foundation.

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Conductivity sum rules in two component plasmas

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Abstract. Conductivity sum rules are established for the linear and quadratic polarizabilities in two component non-relativistic plasma.

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The conductivity and long wavelength compressibility sum rules of plasma physics are model-independent and are therefore useful for checking the accuracy of response functions calculated in various approximations from kinetic equations. By now key sum rules have been formulated for a host of linear [1] and nonlinear [2] response functions for the one component plasma (ocp). Sum rules for tcp (electron-ion two component) systems are, however, still lacking. In this paper we shall partially fill this gap by reporting new tcp conductivity rules for the linear and quadratic polarizabilities, ${}_1\alpha_A$ and ${}_2\alpha_A$ ($A = \text{electrons, ions}$) defined by the relations,

$$\tilde{E}_A^{(1)}(\mathbf{k}\omega) = -{}_1\alpha_A(\mathbf{k}\omega)E^{(1)}(\mathbf{k}\omega), \quad (1)$$

$$\begin{aligned} \tilde{E}_A^{(2)}(\mathbf{k}\omega) = & -(1/V)\Sigma_{\mathbf{q}} \int_{-\infty}^{\infty} (d\mu/2\pi) \frac{{}_2\alpha_A(\mathbf{q}\mu; \mathbf{k} - \mathbf{q}, \omega - \mu)}{1 + {}_1\alpha_A(\mathbf{k}\omega)} \\ & \cdot E^{(1)}(\mathbf{q}\mu)E^{(1)}(\mathbf{k} - \mathbf{q}, \omega - \mu), \end{aligned} \quad (2)$$

connecting the n th order induced electric field $\tilde{E}_A^{(n)}$ and first order (of smallness in the excitation field \hat{E}) total electric fields $E^{(1)} = \hat{E} + \Sigma_A \tilde{E}_A^{(1)}$; note that $E^{(n)} = \Sigma_A \tilde{E}_A^{(n)}$. A derivation of the tcp compressibility rules is deferred to the later paper [3].

Consider a plasma free of externally applied magnetic fields; longitudinal and transverse²⁾ elements of the polarizability matrix are therefore uncoupled. The field quantities in equations (1) and (2) are taken to be along \mathbf{k} so that the polarizabilities there are purely longitudinal. We begin with the derivation of the sum rules for ${}_1\alpha_A$. Upon expanding the r.h.s. denominator of the Kramers-Kronig formula,

$$\text{Re } {}_1\alpha_A(\mathbf{k}\omega) = -(1/\pi)P \int_{-\infty}^{\infty} \frac{d\omega'}{\omega - \omega'} \text{Im } {}_1\alpha_A(\mathbf{k}\omega'), \quad (3)$$

in the $\omega \rightarrow \infty$ limit, one obtains

$$\text{Re } {}_1\alpha_A(\mathbf{k}\omega \rightarrow \infty) = -(1/\pi\omega^2) \int_{-\infty}^{\infty} d\omega' \omega' \text{Im } {}_1\alpha_A(\mathbf{k}\omega') \dots \quad (4)$$

The l.h.s. can be determined from the *macroscopic* equation of motion,

$$(\partial/\partial t)\mathbf{V}_A \cdot (\partial/\partial \mathbf{X}_A)\mathbf{V}_A = (e_A/m_A)\mathbf{E}, \quad (A = \text{electrons, ions}) \quad (5)$$

valid at high frequencies where the electrons and ions behave like collections of non-interacting particles.³⁾ From (5), one obtains

$${}_1\alpha_A(\omega \rightarrow \infty) \approx -(\omega_{pA}^2/\omega^2), \quad (6)$$

where $\omega_{pA} = (4\pi n_A e_A^2/m_A)^{1/2}$. Equations (4) and (6) then combine to give the desired conductivity sum rule,

$$\int_{-\infty}^{\infty} d\omega \omega \text{Im } {}_1\alpha_A(\mathbf{k}\omega) = \pi\omega_{pA}^2 \quad (A = \text{electron, ion}) \quad (7)$$

Equation (7) shows that the sum rule always involves the *dissipative* part of ${}_1\alpha_A$. We

²⁾ With respect to the wave vector k .

³⁾ For the derivation of the compressibility sum rules, it is the static ($\omega = 0$) value of the polarizability which is relevant. In this case, equation (5) must be modified by including (i) the pressure gradient term on its l.h.s. to take account of isothermal processes and (ii) non-local collision terms on its r.h.s. [3].

note that the tcp result [7], while necessary to report, is, nevertheless, a trivial generalization of the well-known ocp conductivity rule [1]. This will be the case for all the conductivity rules reported in this paper and is not surprising in view of the fact that at high frequencies, the non-local collision terms are negligible in the electron and ion macroscopic momentum equations leaving them uncoupled and therefore identical to the ocp $\omega \rightarrow \infty$ momentum equation.

If the tcp is in a constant external magnetic field \mathbf{B}_0 , then for the symmetric elements of the linear polarizability tensor ${}_1\alpha^A$ (denoted by ${}_1\alpha_{[\mu\nu]}^A$)⁴⁾ dissipation is reflected by the imaginary part, whereas for the antisymmetric elements (denoted by ${}_1\alpha_{[\mu\nu]}^A$)⁴⁾ it is reflected by the real part. Thus from $\omega \rightarrow \infty$ denominator expansions of the Kramers-Kronig formulae [2],

$$\text{Re } {}_1\alpha_{[\mu\nu]}^A(\mathbf{k}\omega) = -(1/\pi)P \int_{-\infty}^{\infty} \frac{d\omega'}{\omega - \omega'} \text{Im } {}_1\alpha_{[\mu\nu]}^A(\mathbf{k}\omega'), \tag{8}$$

$$\begin{aligned} \text{Im } {}_1\alpha_{[\mu\nu]}^A(\mathbf{k}\omega) &= (1/\pi)P \int_{-\infty}^{\infty} \frac{d\omega'}{\omega - \omega'} \text{Re } {}_1\alpha_{[\mu\nu]}^A(\mathbf{k}\omega') \\ &- (1/\pi\omega) \int_{-\infty}^{\infty} d\omega' \text{Re } {}_1\alpha_{[\mu\nu]}^A(\mathbf{k}\omega'), \end{aligned} \tag{9}$$

and the high frequency equations of motion,

$$(\partial/\partial t)\mathbf{v}_A + \mathbf{v}_A \cdot (\partial/\partial \mathbf{x}_A)\mathbf{v}_A \approx (e_A/m_A)[\mathbf{E} + (\mathbf{v}_A/c) \times (\mathbf{B} + \mathbf{B}_0)], \tag{10}$$

one obtains the tcp linear sum rules:

$$\int_{-\infty}^{\infty} d\omega \omega \text{Im } {}_1\alpha_{[\mu\nu]}^A(\mathbf{k}\omega) = \pi\omega_{pA}^2 \delta_{(\mu\nu)}, \tag{11}$$

$$\int_{-\infty}^{\infty} d\omega \omega^2 \text{Re } {}_1\alpha_{[\mu\nu]}^A(\mathbf{k}\omega) = -\pi\omega_{pA}^2 \omega_{cA} \varepsilon_{[\mu\nu]\alpha} \hat{b}_{0\alpha}, \tag{12}$$

where $\omega_{cA} = e_A B_0/m_A c$, $\varepsilon_{[\mu\nu]\alpha}$ is the unit permutation pseudotensor and \hat{b}_0 is the unit vector along \mathbf{B}_0 .

We next derive sum rules for the quadratic polarizability in the magnetic field-free tcp. Here the relevant Kramers-Kronig formula is

$$\text{Im } {}_2\alpha_A(\mathbf{k}_1\omega_1; \mathbf{k}_2\omega_2) = (1/\pi)P \int_{-\infty}^{\infty} \frac{d\omega_1'}{\omega_1 - \omega_1'} \text{Re } {}_2\alpha_A(\mathbf{k}_1\omega_1'; \mathbf{k}_2\omega_2). \tag{13}$$

As was pointed out in Ref. [2], useful conductivity rules can be derived only for the limiting case $\omega_2 \rightarrow \infty$ (or $\omega_1 \rightarrow \infty$). In this limit, $\text{Re } {}_2\alpha_A(\mathbf{k}_1\omega_1; \mathbf{k}_2\omega_2 \rightarrow \infty) = -\text{Re } {}_2\alpha_A(\mathbf{k}_1 - \omega_1; \mathbf{k}_2, \omega_2 \rightarrow \infty)$ so that even frequency moments are eliminated from further consideration. Then from the asymptotic expression,

$$\begin{aligned} &\text{Im } {}_2\alpha_A(\mathbf{k}_1\omega_1 \rightarrow \infty; \mathbf{k}_2\omega_2 \rightarrow \infty) \\ &\approx \frac{e_A\omega_{pA}^2}{2m_A\omega\omega_1\omega_2} \left[\frac{k(\mathbf{k}_1 \cdot \mathbf{k}_2)}{\omega k_1 k_2} + \frac{k_1(\mathbf{k} \cdot \mathbf{k}_2)}{\omega_1 k k_2} + \frac{k_2(\mathbf{k} \cdot \mathbf{k}_1)}{\omega_2 k k_1} \right], \end{aligned} \tag{14}$$

$(\mathbf{k} = \mathbf{k}_1 + \mathbf{k}_2, \omega = \omega_1 + \omega_2)$

⁴⁾ In the \mathbf{B}_0 -system [$\mathbf{k} = (k_x, 0, k_z)$, $\mathbf{B}_0 = (0, 0, B_0)$] or in the \mathbf{k} -system [$\mathbf{k} = (0, 0, k)$, $\mathbf{B}_0 = (B_{0x}, 0, B_{0z})$], the 11, 22, 33, 13, 31 elements are the symmetric ones and the 12, 21, 23, 32 elements are the antisymmetric ones.

derived from equation (5), and the high frequency denominator expansion of (13),

$$\begin{aligned} & \lim_{\omega^1 \rightarrow \infty} (1/\pi)P \int_{-\infty}^{\infty} \frac{d\omega_1'}{\omega_1 - \omega_1'} \operatorname{Re} {}_2\alpha_A(\mathbf{k}_1\omega_1'; \mathbf{k}_2\omega_2 \rightarrow \infty) \\ & \approx \lim_{\omega^1 \rightarrow \infty} \left(\frac{1}{\pi\omega_1^2} \int_{-\infty}^{\infty} d\omega_1' \omega_1' \operatorname{Re} {}_2\alpha_A(\mathbf{k}_1\omega_1'; \mathbf{k}_2\omega_2 \rightarrow \infty) + \dots \right), \end{aligned} \quad (15)$$

one obtains the desired conductivity sum rule,

$$\int_{-\infty}^{\infty} d\omega_1 \omega_1 \operatorname{Re} {}_2\alpha_A(\mathbf{k}_1\omega_1; \mathbf{k}_2\omega_2 \rightarrow \infty) = \frac{\pi e_A}{2m_A} \left(\frac{\omega_{pA}}{\omega_2} \right)^2 \frac{k_1(\mathbf{k} \cdot \mathbf{k}_2)}{kk_2}. \quad (16)$$

Similarly, one obtains

$$\int_{-\infty}^{\infty} d\omega_2 \omega_2 \operatorname{Re} {}_2\alpha_A(\mathbf{k}_1\omega_1 \rightarrow \infty; \mathbf{k}_2\omega_2) = \frac{\pi e_A}{2m_A} \left(\frac{\omega_{pA}^2}{\omega_1} \right) \frac{k_2(\mathbf{k} \cdot \mathbf{k}_1)}{kk_1}. \quad (17)$$

Unlike the compressibility sum rules, all the conductivity rules – both linear and quadratic – are completely exhausted by the Vlasov polarizabilities. Finally, we expect that the sum rules derived here will be equally valid for quantum mechanical plasmas.

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On the dielectric response function of two-dimensional classical plasmas

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Abstract. A numerical investigation of the static properties of the 2-*d* electron plasma is carried out on the basis of a recent improved theory of the dielectric response function. The static form factor is computed over the thermodynamically stable domain of the plasma parameter γ . The onset of instability is found to be located in a range not far from the theoretical value.

The longitudinal dielectric response function $\varepsilon(\mathbf{k}, \omega)$ plays a central part in describing various plasma properties [1]; the fluctuations and correlations are readily calculated with a knowledge of such a dielectric function.

In this note we present preliminary numerical results concerning the static properties of the 2-*d* one component classical plasma; a more complete analytical as well as numerical work on the subject will be published elsewhere [19].

In this dimensionality, the model contains various problems of fundamental interest and we refer the reader to the paper of D. Montgomery [2]; especially the

model arising in connection with practical problems in real plasmas, such as diffusion; it has also been recognized that a remarkable model of strong turbulence exists in the highly magnetized 2- d Coulomb plasma [3]; analogous phenomena to those which have been emerged in recent theoretical studies on this model [2], are shown to occur in other plasma domains, such as high beta turbulent regime in magnetohydrodynamics, as reported recently [4].

Within a theoretical approach, not all basic properties of a classical plasma emerge from the study of the model in the small g limit; for a temperature greater than the Fermi energy, there still exist regions in the density temperature plane where the plasma parameter g is greater than unity and for such plasmas, the classical treatment is still applicable.

Lately, much attention has in fact been given to the study of the static properties of a plasma in the region where the dimensionless plasma parameter g is not necessarily small [5, 6].

In the 3- d case, a careful truncation scheme of the BBGKY set of equations, well suited to the purpose (and briefly discussed below, in order to situate our 2- d study), has been advanced by S. Ichimaru [7] and the approach is basically similar to that employed by the same author in his approach to a theory of strong plasma turbulence [8]. Moreover, the validity of such an approximation (which we call the I approximation) has been checked numerically in the 3- d case only [6].

A part of a remarkable improvement with respect to conventional calculations of the random phase approximations for short ranged correlations and other approaches [9], a 3- d numerical study of the I approximation reproduced with good accuracy the range for the onset of instability, which was also predicted and located by a statistical mechanics treatment of the model [10, 11]; in fact, in [6], the pronounced peaks in the static form factor $S(k)$ emerges, as

$$\varepsilon = (4\pi\rho)^{1/2} \cdot e^3 \cdot \beta^{3/2} = \frac{g^3}{4\pi} = \frac{1}{4\pi\rho\lambda_D^3}$$

lie in the range $\varepsilon \sim 9.3 - 10.6$, i.e., $\Gamma = \varepsilon^{2/3} \cdot 3^{-1/3} \sim 3.0-3.4$ while in [11], $\Gamma = 2.5 - 3.3$ as was shown using inequality as well as H -stability property of statistical mechanics [12]. Moreover, these values agree also with that recently found in [13] i.e., $\Gamma \sim 3$, where a change in the slope of the dispersion curve has been emerged.

To remember, the basic assumption involved in such a truncation scheme of the BBGKY set of equations is essentially that the ternary correlation function can be expressed as a functional of the binary correlation; and a physical guideline in selecting the choice of the functional was the correct long range behavior of the correlation in a weakly non ideal 3- d plasma [14].

No, as was remarked in [15, 16], interesting effects are also expected to occur in the 2- d case, where an instability should occur at $\gamma = \beta \cdot e^2 = 1/2\pi \cdot 1/\rho\lambda_D^2 = g^2/2\pi = 4$ [11]. There the pressure as well as the sound speed vanish; notice that the value $\gamma = 2$ occurs in the 2-component case, and is a temperature below which some thermodynamic quantity cannot be proved to exist [17]. (The partition function diverges there even for a $+-$ pair of particles.)

It turns out, that even in the 2- d case the I approximation can be considered as well (this follows from an analysis of the screening factors which occur in the O'Neil-Rostoker equations); in considering a system of electrons in a 2- d box Λ with positive charge background, the 2 nonlinear self-consistent equations relating the structure

factor $S(x)$ and the screening function $1 + u$, written in the dimensionless variable k/k_D , with $k_D^2 = 2\pi\rho\beta \cdot e^2$, $\gamma = g^2/2\pi = \beta \cdot e^2$, $\beta = 1/kT$ are then given by:

$$S(x) = \frac{x^2}{x^2 + 1 + u(x)}$$

$$u(x) = \frac{\gamma}{2\pi} \cdot x \cdot \int_0^\infty dy \int_0^{2\pi} \cos\theta S(y) S(|\mathbf{x} - \mathbf{y}|) \quad (1)$$

where S is related to the radial correlation function by: $\rho g(r) = \sum_{\mathbf{k}} (S(k) - 1) \cdot e^{i\mathbf{k}r}$ and to the static dielectric function by the fluctuation dissipation theorem.

Notice that for $\gamma = 0$, S take on his Debye value: and give rise to an unphysical short range behavior of the correlations; for each γ , the correct long wavelength limit can be found and is given by:

$$u(x) \sim \alpha x^2 + \beta x^4$$

with

$$\beta = \frac{\gamma}{16} \cdot \int_0^\infty \left(\frac{dS}{dy} \right)^2 \cdot \frac{dy}{y} \quad (2)$$

and where α is connected with the sound speed; it can be shown, that the compressibility sum rules are satisfied for each γ , contrary to the 3- d case.

A numerical solution in the range up to $\gamma = 4$ and $x = 6$ has been found using digital computer; a short description of the method used is now briefly presented together with the new results.

The set of equations (1) to be solved for $S(x) - 1 = (\rho \cdot g)(x)$ and $u(x)$, after the variable change $e^{-x} = X$, $e^{-Y} = Y$ take the form:

$$U(X) = -\frac{\gamma}{\pi} \cdot \ln X \cdot \int_0^1 dY \frac{1 + G(Y)}{Y} \int_0^\pi d\theta \cos \theta \cdot G(Z)$$

$$G(X) = -\frac{1 + U(X)}{(\ln X)^2 + 1 + U(X)} \quad (3)$$

where $z = \text{esp} - [(\ln X)^2 + (\ln Y)^2 - 2 \ln X \cdot \ln Y \cos \theta]^{1/2}$

$$u'(-\ln x) = U(X)$$

$$(\rho g)'(-\ln x) = G(X)$$

The integral occurring in (3) is discretized by means of Gauss method, i.e.:

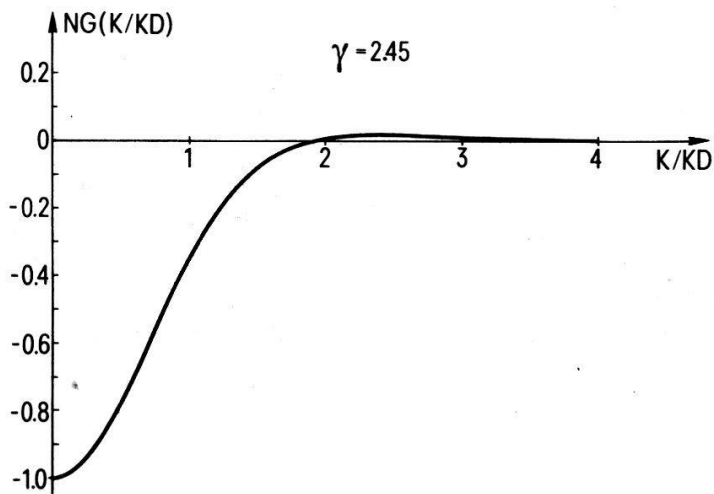
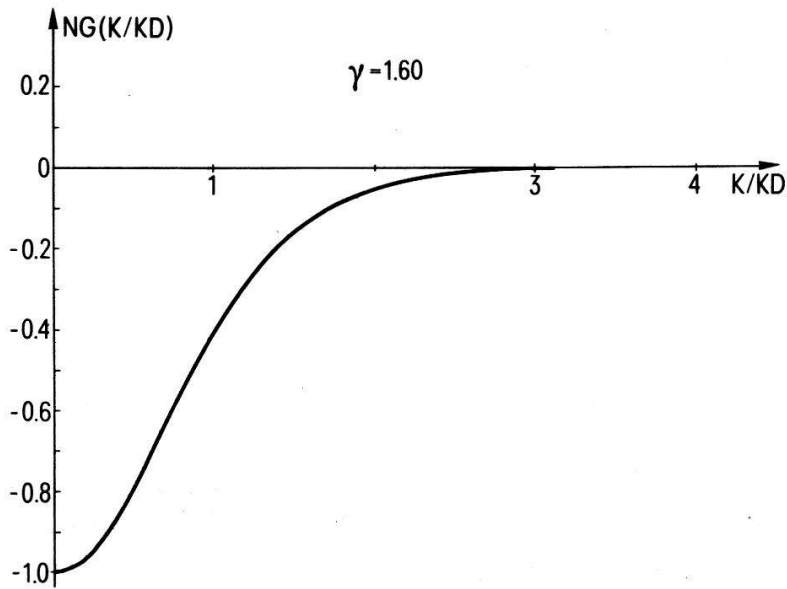
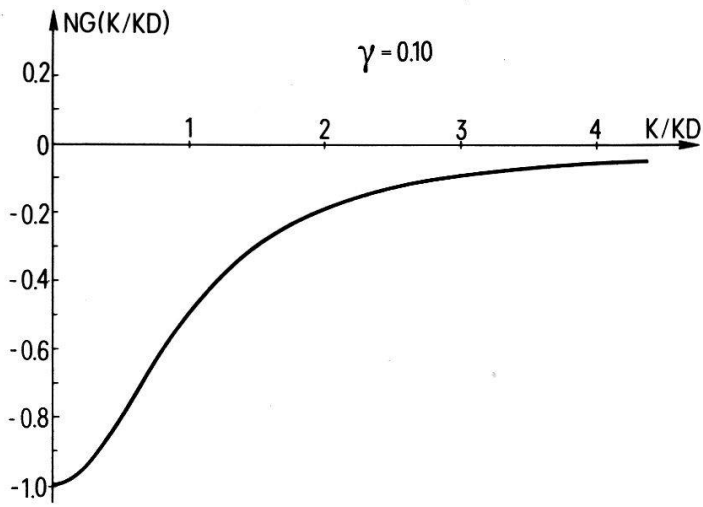
$$U(X_j) = -\frac{\gamma}{4} \cdot \ln X_j \sum_{k=1}^n w_k^n \frac{1 + G(X_k)}{X_k} \sum_{l=1}^m w_l^m \cos \theta_e G(Z_{j,k}^l)$$

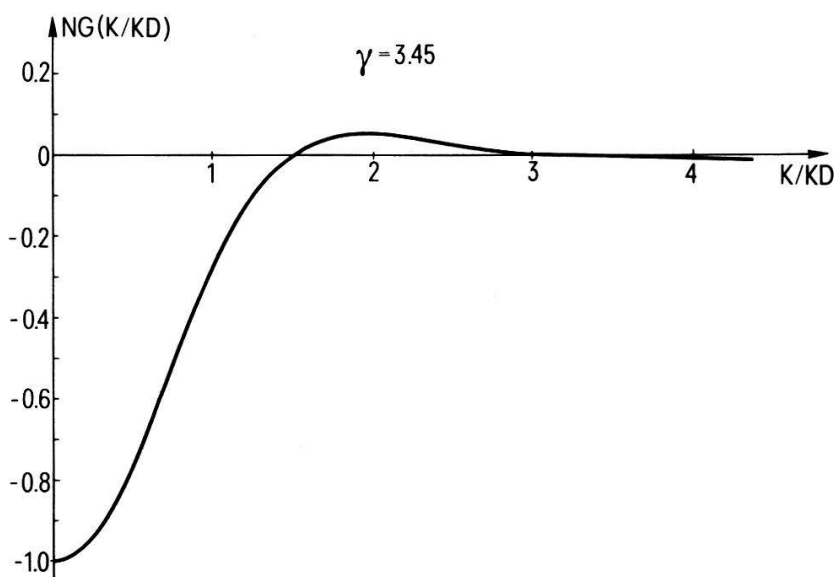
$$G(X_j) = -\frac{1 + U(X_j)}{(\ln X_j)^2 + 1 + U(X_j)}$$

and where $j \in [1; n]$; $\{x_j\}$, $\{\theta_j\}$, $\{W_j^m\}$, $\{W_j^m\}$ are given by the method.

Having first computed the set $\{Z_{j,k}^l\}$, we have then deduced $G(Z_{j,k}^l)$ interpolating $\{G(X_j)\}$ by means of cubic spline method [18]; the relation $G(0) = 0$ and $G(1) = -1.0$ has also been employed. We thus have calculated the set $\{U(X_j)\}$ and deduced the new set $\{G(X_j)\}$. The iteration has been stopped if

$$\text{Max}_{j \in [1, n]} (|G^{i+1}(X_j) - G^i(X_j)|) < 10^{-8}$$





It will be shown in [19] that this method can be improved in considering the decomposition:

$$U(X) = -\frac{\gamma}{\pi} \ln X \left\{ \int_0^A dY \frac{1 + G(Y)}{Y} \int_0^\pi d\theta \cos \theta G(Z) + \int_A^1 dY \frac{1 + G(Y)}{Y} \int_0^\pi d\theta \cos \theta G(Z) \right\}.$$

it will then be possible to work with a denser set of points in the interval $[0, A]$ to precise the structure of the peak or oscillations which have emerged in the static form factor.

Moreover, the initial value problem has been solved, using for any value of γ the corresponding result of the step $\gamma - \epsilon$. In the first step ($\gamma = 0.05$) the initial value for $g(x)$ has been chosen equal to his Debye limit ($u = 0$).

These results obtained for the different set of γ will then be used in [19] as initial values to further approach the correlation function $g(k)$ solution of (1) and this in a much bigger domain, containing smaller wavelengths. Results on $\rho \cdot g$ for different values of γ are plotted in Figures 1, 2, 3, 4. Above $\gamma = 1.6$, $g(k)$ begin to overshoot the zero level and at $\gamma = 3.45$ a peak emerges, and will be pronounced for γ approaching the value 4; a careful examination of the screening function in this region is in progress. Moreover it is expected that in the short range domain, $g(r)$ violate the physical condition that $g(r) \geq -1$ above some $\gamma_0 < 4$, and this only on a reduced scale, as in the 3-dimensional case; this may be related to the change of $\beta(\gamma)$ at γ_0 and thus to a change in the structure of the long wavelength part of the fluctuations, of Ornstein-Zernike type, occurring in the system [19].

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Hochstabiler Farbstofflaser und spektroskopische Anwendung

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Zur hochauflösenden Spektroskopie im Bereich von 420–800 nm an Atom- und Molekularstrahlen wurde ein frequenzstabilisiertes Farbstofflaser-System mit einer Bandbreite < 10 MHz entwickelt. Der Freistrahllaser eigentümliche Frequenzjitter konnte durch Verwendung einer mit Polyvinylalkohol verdickten Wasserlösung wesentlich reduziert werden. Thermisch bedingtes Langzeitdriften des Resonatorsystems wird mit Hilfe elektronisch abstimmbarer Piezokristalle kompensiert, mit denen ebenfalls die Frequenz des Lasers synchron abgestimmt werden kann. Mit dem Lasersystem konnte ein Na-Atomstrahl abgelenkt und die Hyperfeinstrukturaufspaltung der *D*-Linie gemessen werden. Weiterhin wurden molekulare Systeme spektroskopisch untersucht.

Farbreproduktion in der nullten Beugungsordnung

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Wir beschreiben und demonstrieren ein neuartiges Verfahren zur Speicherung von Farbbildern. Die Information ist als reine Phasenstruktur in einem Oberflächenrelief eines transparenten Medium gespeichert. Die Bilder können wie konventioneller Farbfilm auf Dia- oder Mikrofilmprojektoren ausgelesen werden. Zur Farberzeugung dienen drei überlagerte Rechteck-Phasengitter, die in der nullten Beugungsordnung geeignete spektrale Eigenschaften besitzen. Die Bildinformation wird als binäre Amplitudenmodulation dem Gitter überlagert.

Die Reliefstrukturen werden zuerst in Photolack hergestellt und können dann preisgünstig und einfach durch Prägen in Plastik vervielfältigt werden. Im Vergleich

mit konventionellem Farbfilm zeigt das neue Verfahren etwa den gleichen Farbumfang, jedoch bessere Auflösung und Stabilität, und bietet einen Kostenvorteil von einem Faktor ≥ 10 bei genügend hoher Auflage.

Size, shape and electrostatic interaction of human fibrinogen in solution

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Summary. Rotational and translational diffusion constants were determined on human fibrinogen solutions with a clottability greater than 95% by means of dynamical light scattering. The sedimentation coefficient was measured by ultracentrifugation. The results can be made to agree qualitatively with the predictions based on the model by Hall and Slayter [1] if one assumes that in addition to the single molecules associates in end-to-end configuration are present (pairs, triplets etc.). The association seems to be reversible and is favoured by increasing fibrinogen concentration. The sphere-like model of Köppel [2] and the rod-like model of Bachmann [3] are hydrodynamically incompatible with our measurements. When the ionic strength is reduced below physiological values a higher diffusion constant appears. There is some evidence that it might be due to a very strong electrostatic interaction between the fibrinogen molecules causing local order.

1. Introduction

In spite of numerous electron microscopic studies and a large body of knowledge from physicochemical investigations there is no agreement on the size and shape of the fibrinogen molecules [4].

The intriguing diversity of published diffusion constants can hardly be explained unless interactions between the molecules are postulated. In the present study we have investigated the possible effect of molecular association and electrostatic intermolecular interactions on the Brownian motion.

2. Experimental results

2.1. Dynamic light scattering [5]

The correlation function corresponding to *translational diffusion* deviates from a single exponential. The initial slope is found to depend upon the concentration of fibrinogen and on the ionic strength of the solvent. For example, at a concentration of 10 mg/ml in 0.15M Tris-NaCl the initial slope yields a translational diffusion constant $D_{T_{20,W}} = (1.7 \pm 0.3) 10^{-7} \text{ cm}^2/\text{s}$. For concentrations above 2 mg/ml and the same solvent the experimental data can be fitted perfectly to two exponentials yielding

$D_{T_{20,w}} = (2.0 \pm 0.1) \times 10^{-7} \text{ cm}^2/\text{s}$ for the first component and $D_{T_{20,w}}$ between 1.4 and $0.8 \times 10^{-7} \text{ cm}^2/\text{s}$ for the second. With decreasing concentration below 2 mg/ml the fit to two exponentials becomes increasingly poor and a third, still larger diffusion constant emerges. These results are compatible with a polydisperse solution, in which single molecules are in equilibrium with associates.

A hint about the nature of the associates has been found in the measurement of the correlation function corresponding to *rotational diffusion*, in which at least two exponentials can be distinguished in high concentration samples. A fast rotational diffusion with $D_{R_{20,w}} = 40\,000 \pm 20\% \text{ s}^{-1}$ and a slow one with $D_{R_{20,w}} = 10\,000 \pm 30\% \text{ s}^{-1}$ are observed. The large difference between these two constants suggests that the molecules associate end-to-end.

Using this picture and taking the Hall-Slayter [6] model for the single molecule we made an attempt to fit the observed correlation function for translational diffusion. The diffusion of the associates can be expressed in terms of the diffusion of the single molecules without the introduction of new hydrodynamic parameters. The adjustable parameters were the length of the Hall-Slayter molecule and the relative concentrations of the different associates. A minimum of three exponentials is necessary for an acceptable fit. The resulting translational diffusion constants $D_{T_{20,w}}$ attributed to single molecules, pairs and triplets are $(2.8 \pm 0.4) \times 10^{-7} \text{ cm}^2/\text{s}$, $(1.9 \pm 0.4) \times 10^{-7} \text{ cm}^2/\text{s}$ and $(1.2 \pm 0.4) \times 10^{-7} \text{ cm}^2/\text{s}$, respectively. The relative concentrations of the different associates is found to depend upon the total fibrinogen concentration, the single molecules being predominant in the dilute solutions.

2.2 Centrifugation

The mean sedimentation coefficient for a dilute solution of fibrinogen (0.1 mg/ml) was found to be $7.9 \pm 0.1 \text{ S}$. The concentration gradient at the boundary as determined by the Schlieren method deviates from a Gaussian, confirming polydispersity. Thus the light scattering results are qualitatively compatible with the sedimentation parameters although the latter are not sensitive with respect to the length of the chain of end-to-end associated molecules. On the basis of the Hall-Slayter model and the assumption of end-to-end association one calculates the following ratio of the sedimentation velocities of the different associates

$$v_{\text{single}} : v_{\text{pair}} : v_{\text{triplets}} = 1 : 1.14 : 1.20$$

2.3 Molecular charge and electrostatic intermolecular interaction

On the basis of the accepted molecular weight of 340 000 and of the known total number of titrable amino acids we have estimated the number of negative and positive charges and the net charge for different pH values. Assuming the absence of counterions in the solution we obtain:

pH	net charge in units of $1.6 \times 10^{-19} \text{ clb}$
7.4	-165
9.5	-250

The electrophoretic mobility measured by means of dynamical light scattering at pH 9.5 and 0.007 molar salt concentration is $2.4 \pm 0.2 \text{ (microns/s)/(volt/cm)}$. This corresponds to a charge of 100 electrons in reasonable agreement with the above calculations.

In a solution of 4 mg/ml the mean distance between two fibrinogen molecules is about 1000 Å. Two point charges consisting each of 165 electrons immersed in a medium with the dielectric constant of water have a potential energy of about 250 kT. This suggests that at zero ionic strength the diffusion might be strongly influenced by electrostatic interactions. The dynamic light scattering shows indeed that for salt concentrations below 0.05 molar a diffusion coefficient appears which is several times larger than for single fibrinogen molecules at physiological ionic strength. It is conceivable that this fast component is due to local order in the fibrinogen solution, resembling the lattice formation in suspensions of latex particles in solvents of zero ionic strength [7].

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Druckinduzierte Mikrowellen-Absorption in Gasgemischen

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Es wurde bei 94 Ghz in einem Brückenspektrometer mit einer Empfindlichkeit von $5 \cdot 10^{-8} \text{ cm}^{-1}$ die Restabsorption in Gasgemischen von H_2O , SO_2 , CO_2 , N_2 im Druckbereich 0.1–200 Torr gemessen. Die Restabsorption kann mit der Bildung von Molekülassoziationen erklärt werden. Im Falle von SO_2 – H_2O – N_2 -Gemischen wird die durch Van-der-Waals'sche Kräfte hervorgerufene Assoziationsbildung diskutiert.

Darstellung der mit Hilfe der Elektronenmikroskopie ermittelten räumlichen Anordnung von Kristallbaufehlern¹⁾

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Mit Hilfe der Elektronenmikroskopie [1] lassen sich Aussagen darüber machen wie die zu Bildkontrasten führenden Gitterbereiche (z.B. Kristallbaufehler) innerhalb

¹⁾ Diese Arbeit wurde teilweise vom Schweizerischen Nationalfonds zur Förderung der wissenschaftlichen Forschung unterstützt, Kredit Nr. 2.486–0.75.

der Probe angeordnet sind. Handelt es sich um reine Tiefenverteilungsmessungen von kleinen begrenzten Gebieten (z.B. Ausscheidungen, Zwischengitteratom-oder Leerstellenagglomerate), so lassen sich die Ergebnisse in Histogrammen wiedergeben (z.B. [2]). Bei ausgedehnten Kontrasten (z.B. Versetzungen, Stapelfehler oder Korngrenzen) bereitet die figürliche Darstellung oft Schwierigkeiten.

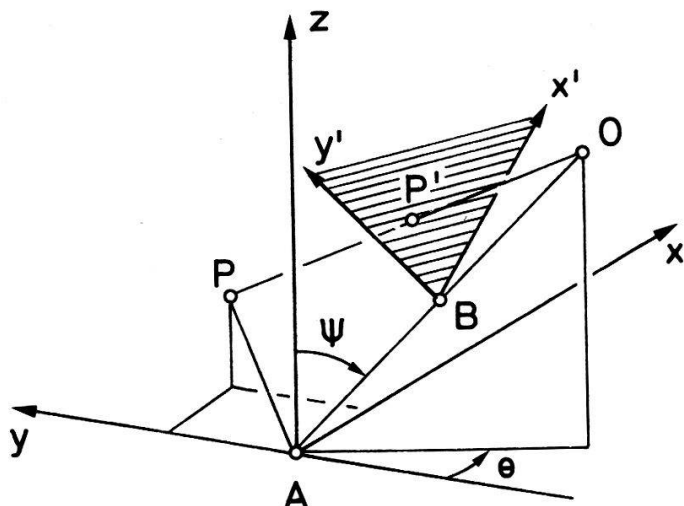


Abbildung 1

Die dem Rechenprogramm zu Grunde gelegte Projektion.

- P : Messpunkt (x, y, z) , O : Beobachtungspunkt,
- P' : Bildpunkt in der Projektionsebene $Bx'y' \perp AO$

Für eine zweidimensionale Wiedergabe der räumlichen Anordnung solcher Kristallbaufehler wurde ein einfaches Rechenprogramm entwickelt [3]. Die dabei verwendete Projektion zeigt Abbildung 1. Das Programm erlaubt, ausgehend von x, y, z -Messwerten, die Anordnung unter beliebig wählbaren Beobachtungsbedingungen (Beobachtungsrichtung und Abbildungsmaßstab) zu betrachten und direkt von einem X-Y-Schreiber aufzeichnen zu lassen. Eine einmal gemessene Anordnung kann somit auch nachträglich aus verschiedenen Richtungen beobachtet werden. Es

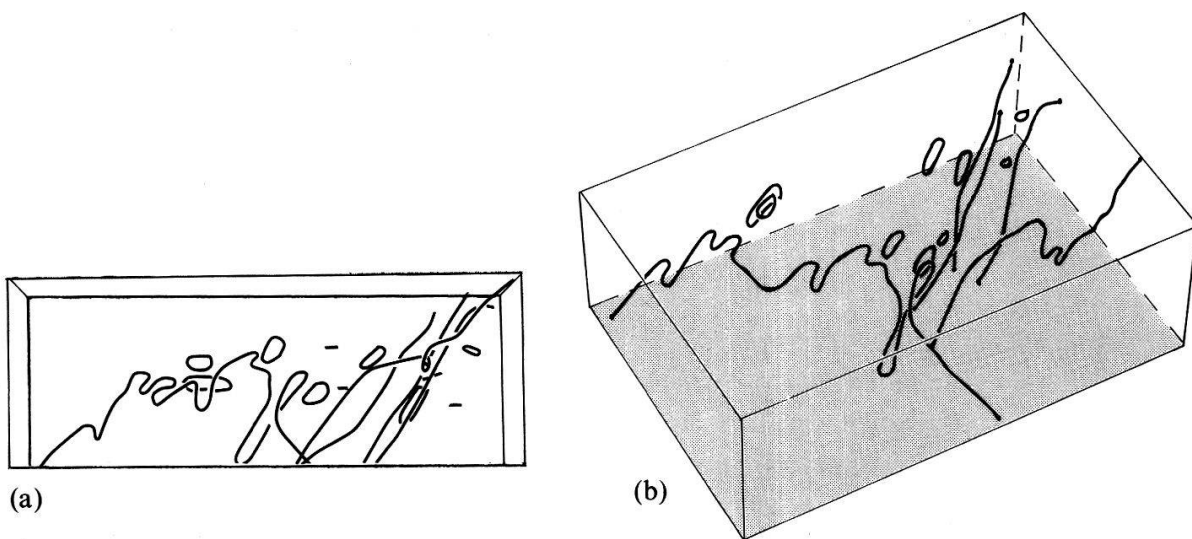


Abbildung 2

Räumliche Darstellungen derselben Versetzungsanordnung für zwei verschiedene Beobachtungsbedingungen (Bedeutung von θ und ψ siehe Abbildung 1)

- (a) $\theta = -90^\circ; \psi = 90^\circ$
- (b) $\theta = -120^\circ; \psi = 40^\circ$

ist nicht notwendig, zusätzliche Stereopaare im Mikroskop aufzunehmen, bei denen bei Änderung der Abbildungsbedingungen Kontraste von Versetzungen mit bestimmten Burgersvektoren unsichtbar werden können. Ausserdem wird mit dieser Methode die Darstellung dreidimensionaler kristallographischer Zusammenhänge wesentlich erleichtert. Als einfaches Beispiel zeigt Abbildung 2 eine Versetzungsanordnung aus zwei verschiedenen Blickwinkeln.

Der berechnete, wenig realistische Verlauf der Versetzungslinien in der Nähe der nur *ca.* 100 Messpunkte wurde leicht abgerundet wiedergegeben. Dieser Eingriff in die berechneten Perspektiven lässt sich umgehen, wenn man die Zahl der Messpunkte erhöht. Automatische Stereo-Auswertgeräte, wie sie in der Photogrammetrie üblich sind, könnten hierbei von grossem Nutzen sein. Solche Geräte müssten allerdings den speziellen Abbildungsbedingungen der Elektronenmikroskopie angepasst werden.

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Quantitative Elementenanalysen mit einer Protonen-Mikrosonde

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Werden die Atome mit Protonen angeregt und die entstehende Röntgenstrahlung mit einem hochauflösenden Halbleiterdetektor registriert, so können die meisten Elemente des periodischen Systems in festen Proben mit einer Empfindlichkeit von *ca.* 0.1–1 ppm nachgewiesen werden. Es wird eine Protonen-Mikrosonde mit einem digitalen Abtastsystem beschrieben, die es erlaubt die Verteilung der Spurenelemente mit einer Auflösung von *ca.* 20 μm quantitativ zu bestimmen. Anhand von Beispielen werden einige Anwendungsmöglichkeiten einer solchen Apparatur gezeigt.

Volatile and non-volatile elements in grain-size separated samples of Apollo 17 lunar soils

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Three samples of Apollo 17 lunar soils (75081, 72501 and 72461) were separated into 9 grain-size fractions between 540 and 1 μm mean diameter. In order to reduce the risk of contamination the samples were irradiated with thermal neutrons prior to the mechanical handling. A semi-automatic sieving procedure was used for the separation of the highly radioactive material. Seven size fractions (540–11 μm) were obtained by sieving, two additional fractions of mean grain-sizes of 3 and 1 μm diameter resulted from sedimentation in acetone. Microscopic and electron-microscopic examination show a uniform size distribution within each of the fractions obtained by these techniques. In order to detect mineral fractionations caused during the separation procedures major elements were determined by instrumental neutron activation analyses performed on small aliquots of the separated samples.

Twenty elements were measured in each size fraction using instrumental and radiochemical neutron activation techniques. The samples were dissolved with HF-HClO_4 mixtures in teflon lined pressure bombs for radiochemical determina-

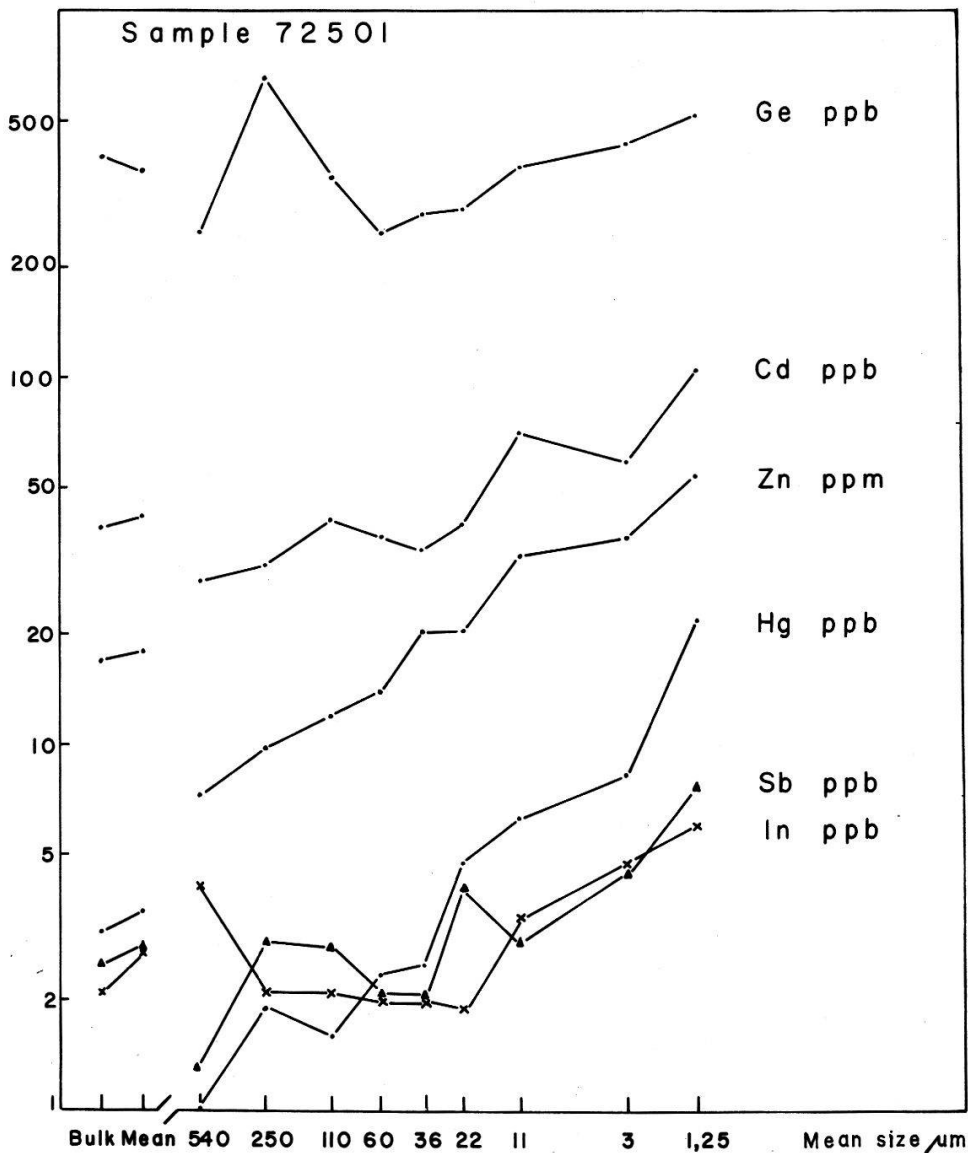


Figure 1

Distribution of volatile elements in grain-size separated samples of Apollo 17 soil 72501. The 'Bulk'-value resulted from small aliquots of the unseparated sample. The 'Mean'-value is the weighted average of the individual size fractions. Errors (1σ) are of the order of 10–20%.

tions. Sulfide precipitations and methods described by Keays *et al.* [1] were used to separate and clean the elements of interest.

The concentration of the main elements in sample 75081 does not change with the grain-size. Exceptions are Fe and Ti which decrease slightly and Al which increases slightly with the decrease in the grain-size. These changes in the composition in main elements suggest a decrease in Ilmenite and an increase in Anorthite with decreasing grain-size. However, it can be concluded that the mineral composition of the fractions changes less than a factor of 2. Samples 72501 and 72461 are not yet analyzed for the main elements.

Figure 1 shows the results of the determination of volatile trace elements in sample 72501. Similar distribution patterns were obtained for 75081 and 72461. The concentrations of the investigated volatile elements are anticorrelated to the grain-size, suggesting a surface correlation. The irregularities found in the distribution of Ge and In are not fully understood and need further investigation. Non-volatile elements such as uranium and rare earths show no dependency on the size of the grains. Our results are in qualitative agreement with investigations of Boynton *et al.* [2, 3] on lunar soils from other landing sites.

Acknowledgment

The help of Miss E. Rössler and Mr. D. Jost in sample preparation and analysis is highly appreciated. The authors thank Prof. J. Geiss for helpful discussions. Part of the work was supported by the Swiss National Science Foundation.

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Effektivwertmessungen der 'lokalen' Komponente des Neutronenflussrauschens in einem Siedewasserreaktor

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In verschiedenen in der letzten Zeit erschienenen Arbeiten wird versucht, einen Zusammenhang zwischen dem RMS – Wert der 'lokalen' Komponente der Neutronenflussschwankungen im Core eines Siedewasserreaktors und dem volumetrischen Dampfgehalt herzustellen. Zu diesem Zweck wurde eine Serie von Messungen am BWR des Kernkraftwerkes Mühleberg durchgeführt. Der sich heutzutage weitgehend durchgesetzten Ansicht folgend, dass zwischen 'Globaleffekt' und 'Lokaleffekt' eine Trennung längs der Frequenzskala besteht, haben wir den 'lokalen' Anteil mit Hilfe eines Bandpasses (7–40) Hz ausgesiebt.

Whistler and electrostatic instabilities in a plasma trapped in a magnetic field (loss-cone instabilities in the solar corona)

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The initial growth rates of whistler and electrostatic perturbations are computed numerically in a situation representative of the solar corona during a type IV burst. The non-relativistic case is considered. The plasma is completely ionized. The electron velocity distribution is made of two parts: a thermal one, isotropic, and a small fraction of fast electrons trapped in an arched magnetic field. The fast electrons distribution is a loss-cone type (the plasma is homogenous in space).

The computation of the whistler instability is made for propagation parallel to the magnetic field B_0 with the formulas of Kennel *et al.* (1966). The plasma is whistler unstable for large ranges of the parameters (some curves are presented, showing the dependance of the instability in the following parameters: distribution, loss-cone angle, anti-loss-cone angle, dispersions of the distributions, plasma frequency).

For electrostatic perturbations with propagation perpendicular to B_0 , the magnetic field effects are very important. The general dispersion relation of Rosenbluth *et al.* (1965) is considered. The numerical solution gives instability where the 'Berstein mode' at the harmonic of the cyclotron frequency under the upper hybrid frequency is coupled with the upper hybrid frequency mode. The plasma is unstable with relatively severe conditions over the parameters: e.g. the instability area lies in a 'band' of the plane defined by the plasma frequency and the fraction of fast particles. An important result is that the plasma frequency may not be harmonic of the cyclotron frequency to be unstable. The necessary condition is that the fraction of fast particles may be large enough.

Furthermore the instability disappears when the temperature of the fast electrons approaches the one of the thermal electrons.

Electrostatic perturbations parallel to B_0 are unstable only in the presence of an anti-loss-cone in the velocity distribution of fast electrons.

The maximas of electrostatic initial instability are an order of magnitude faster than the whistler initial instability, for the same parameter values.

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Excitation function of the reaction $^{85}\text{Rb}(p, 5n)^{81}\text{Sr}$

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The excitation function of the reaction $^{85}\text{Rb}(p, 5n)^{81}\text{Sr}$ was determined by means of stacked foil technique for proton energies between 54–69 MeV. By comparing the results to the theoretical predictions it was found that the shape of both $\sigma(E)$ curves are similar. However, the experimental cross sections are lower and perhaps shifted towards higher energies.

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KERN- UND TEILCHENPHYSIK

Experience complète sur la diffusion p - p aux moyennes énergies

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La description de la diffusion p - p au-dessous de 1 GeV est réalisée par l'analyse en déphasage qui appuie ses prédictions sur 3000 points expérimentaux comprenant environ 60% de mesures de la section efficace différentielle ou totale, 35% de mesures du paramètre de polarisation, et 5% seulement de mesures des tenseurs de dépolarisation ou de transfert. Nous proposons ici un ensemble de mesures sur la diffusion élastique de protons polarisés entre 400 et 600 MeV, qui permet la détermination directe et complète des amplitudes de la matrice de diffusion, indépendamment pour chaque angle et chaque énergie. Nous discutons la précision des expériences nécessaire à une bonne détermination des amplitudes et nous montrons que ces mesures sont accessibles aux appareillages actuels pour la diffusion élastique p - p à grand angle.

Diffusion quasi-libre neutron-neutron à 14,1 MeV

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Deux mesures cinématiquement complètes de la réaction $D(n, nn)p$ ont été effectuées pour des configurations favorisant la diffusion quasi-libre n - n . Ces mesures

montrent que la section efficace différentielle à $\theta_1 = \theta_2 = 40^\circ$, $\phi_{12} = 180^\circ$ ($E_p^{\min} = 0$) est environ deux fois plus petite qu'à $\theta_2 = 30^\circ$ ($E_p^{\min} = 180$ keV), alors que les calculs 'exacts' basés sur les équations de Faddeev prédisent un facteur 1.3 seulement. De plus la forme de la section efficace différentielle à 30° présente une structure marquée qui n'est prédite par aucun calcul théorique existant.

Protonentransferreaktionen in ungebundene Zustände

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Die Protonentransferreaktion $^{12}\text{C}(d, n)^{13}\text{N}$, die zu den 3 niedrigsten Resonanzen in ^{13}N führt, wird mit Hilfe der allgemeinen DWBA-Theorie für Strippingreaktionen in Kontinuumszustände [1] analysiert. Dabei werden als 'input' die experimentell bestimmten Phasen für die elastische p - ^{12}C -Streuung verwendet. Die modellabhängigen Beiträge, die vom Kerninneren kommen, werden abgeschätzt, indem verschiedene Resonanzwellenfunktionen für das ungebundene Proton benutzt werden und die damit berechneten Resonanzlinienformen mit dem Experiment verglichen werden.

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Elastische Transferprozesse

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Anhand verschiedener elastischer Transferprozesse wie $^{12}\text{C}(^{13}\text{C}, ^{13}\text{C})^{12}\text{C}$, $^6\text{Li}(\alpha, \alpha)^6\text{Li}$, $^{16}\text{O}(\alpha, \alpha)^{16}\text{O}$ wird die Bestimmung von spektroskopischen Faktoren aus gemessenen Winkelverteilungen mittels DWBA-Analysen untersucht. Rückstoss, endliche Reichweite und Coulomb-Wechselwirkungsterme werden mitberücksichtigt. Der Einfluss des optischen Potentials und der Form der gebundenen Wellenfunktion auf den ermittelten spektroskopischen Faktor werden diskutiert. Es zeigt sich, dass deren Einfluss nur entscheidend unterdrückt werden kann bei Einschussenergien im Subcoulombbereich und bei gleichzeitig genügend grossem Wert des Coulombparameters der reagierenden Teilchen.

Ein optisches Potential für schwere Ionen mit partieller Berücksichtigung des Pauli-Prinzips

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Das Nukleus-Nukleus-Potential, das durch Faltung der Dichte des einen Kernes und des Nukleon-Nukleus-Potentials des andern erzeugt wird, verletzt das Pauli-

Prinzip. Durch geeignete Ueberlagerung aller möglichen Nukleonen-Austausch-Amplituden kann jedoch die Streuwellenfunktion antisymmetrisiert werden. Es ist deshalb im Prinzip möglich, mit mehreren Streukanälen das Pauli-Prinzip zu berücksichtigen. Wir berechnen nun mit Hilfe der Theorie der gekoppelten Kanäle ein optisches Potential, das einen Teil dieser Austauschkanäle mitberücksichtigt.

Kernreaktionen mit ${}^9\text{Be}$

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Die elastische Streuung von ${}^9\text{Be}$ an Mg, Al, Si, Ca, Ni, Ag und Au wurde bei 14, 20 und 26 MeV gemessen. Als Wechselwirkungsdistanz (Abstand, bei dem die Kernwechselwirkung einzusetzen beginnt) wurde $d = 1.82 \cdot (A_t^{1/3} + A_b^{1/3})$ gefunden. Die Beschreibung der Winkelverteilungen mit einem optischen Potential der Form

$$V(r) = (V_0 + V_1 \cdot E + iW_0 + iW_1 \cdot E) \cdot f(r) + V_{\text{coul}}(r)$$

mit

$$f(r) = \frac{1}{1 + e^{(r - r_0 \cdot A_t^{1/3} + r_b)/a_0}}$$

ergab beste Uebereinstimmung für folgende Parameter:

$$\begin{array}{ll} V_0 = 5.17 \pm 0.74 \text{ MeV}; & V_1 = 0.15 \pm 0.04 \text{ MeV} \\ W_0 = 16.10 \pm 1.39 \text{ MeV}; & W_1 = 0.07 \pm 0.10 \text{ MeV} \\ r_0 = 1.22 \pm 0.03 \text{ fm}; & a_0 = 0.825 \pm 0.014 \text{ fm} \\ r_b = 2.46 \pm 0.08 \text{ fm} & \end{array}$$

Ausserdem wurden an Si Anregungsfunktionen von 14 bis 20 MeV unter vier Winkeln gemessen. Sie zeigten keine Struktur und können mit den gleichen optischen Parametern beschrieben werden.

Messungen von Neutrontransferreaktionen mit ${}^9\text{Be}$ von 20 MeV an Si und Ca zeigten, dass der Spin des Endzustandes auf die Form der Winkelverteilung der ${}^8\text{Be}$ wenig Einfluss hat. Aus dem Vergleich von gemessenen und mit einem 'finite range DWBA' - Programm [1] berechneten Wirkungsquerschnitten konnte aus neun bekannten Niveaux in ${}^{29}\text{Si}$, bzw. deren acht in ${}^{41}\text{Ca}$ der spektroskopische Faktor von ${}^9\text{Be}$ bestimmt werden. Wir fanden 0.44 bei ${}^{29}\text{Si}$ und 0.41 bei ${}^{41}\text{Ca}$. Für die Berechnung der gestörten Wellen im Ausgangskanal wurden für ${}^8\text{Be}$ die optischen Potentiale von ${}^9\text{Be}$ eingesetzt.

Der direkte Aufbruch von ${}^9\text{Be}$ im Coulombfeld kann mit einem semiklassischen Modell berechnet werden [2]. Wir haben diese Wirkungsquerschnitte mit einem Goldtarget bei Energien von 18 bis 26 MeV unter verschiedenen Winkeln gemessen. Das Verhältnis von experimentellen und theoretischen Wirkungsquerschnitten sollte den spektroskopischen Faktor von ${}^9\text{Be}$ ergeben. Wir erhielten jedoch 2.4.

${}^9\text{Be}$ bildet mit vielen Targets hochangeregte Compoundkerne, die durch Verdampfen von Teilchen zerfallen. Dabei werden vor allem Protonen, Neutronen und Alphas emittiert. Die Spektren der Protonen und Alphas wurden unter verschiedenen Winkeln gemessen und werden mit einem statistischen Modell [3] analysiert.

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Radioactive pion capture in light nuclei

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(SIN)

γ -ray spectra from pion capture by different light nuclei (${}^6\text{Li}$, ${}^7\text{Li}$, ${}^9\text{Be}$, ${}^{12}\text{C}$, ${}^{16}\text{O}$, ${}^{18}\text{O}$, ${}^{28}\text{Si}$, ${}^{32}\text{S}$, . . .) obtained with the SIN pair spectrometer will be presented.

The 1 MeV resolution permits separation of the contributions of different giant resonance final states corresponding to analog states seen either in electron inelastic scattering or in photoexcitation which are due to M1 or E1 transitions.

 π^+ and π^- scattering from ${}^{12}\text{C}$ with the SIN pion spectrometer

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High resolution elastic and inelastic data has been obtained for π^+ and π^- scattering from ${}^{12}\text{C}$. Spectra were taken with the SIN pion spectrometer at 150, 165 and 230 MeV and a wide range of angles. Figure 1 shows a typical spectrum with π^- at 165 MeV and 85 degrees. The elastic peak and the inelastic 4.4 MeV(2^+), 7.6 MeV(0^+) and 9.6 MeV(3^-) states are clearly separated. In addition, levels up to excitation energies of 20 MeV can be seen. The analysis is still in progress. Our measurement should complement well the $\pi^- + {}^{12}\text{C}$ data by Binon *et al.* [1].

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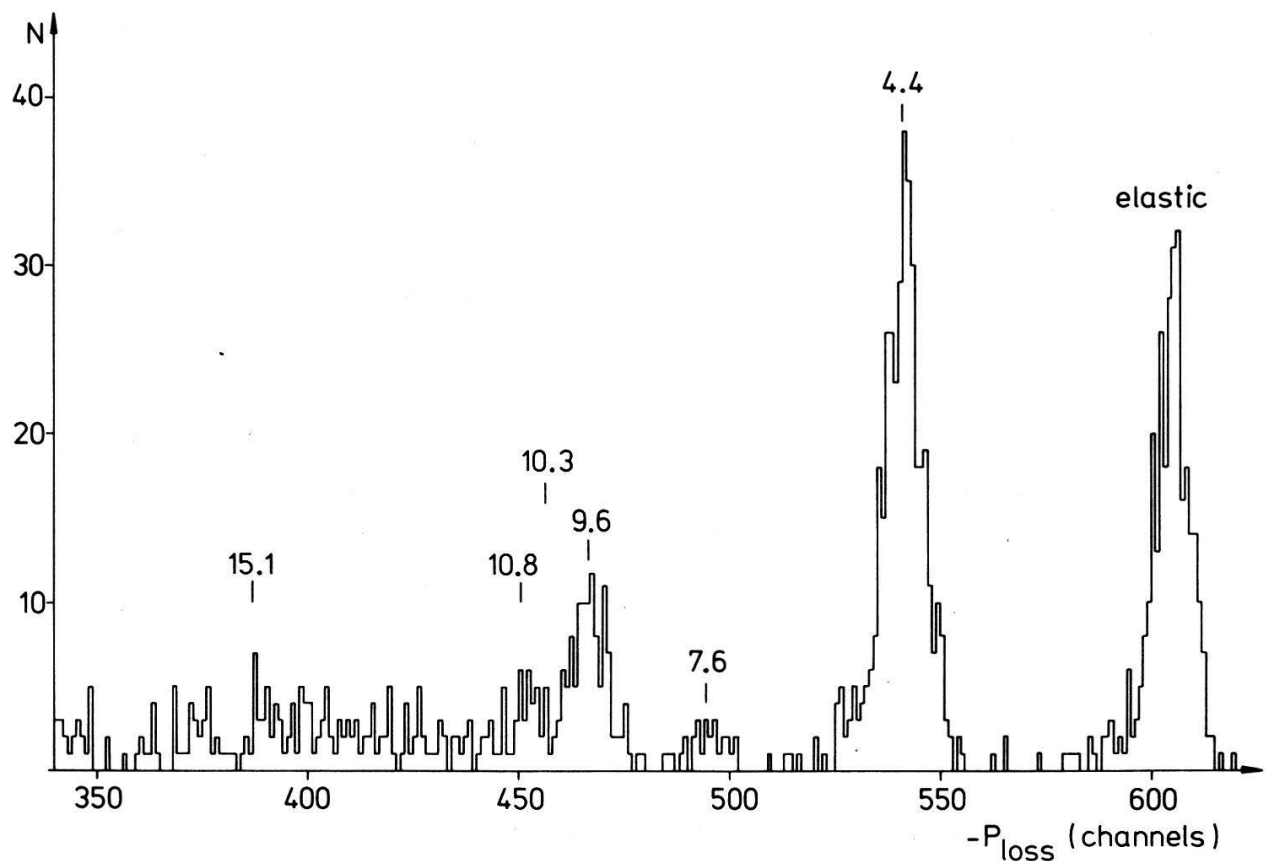


Figure 1
 $\pi^- + {}^{12}\text{C}$ scattering. $T_0 = 165$ MeV, $\theta_{\text{lab}} = 85^\circ$. Off-line spectrum including kinematic correction. Spurious events (with trajectories incompatible with optics) were rejected.

Kernspaltung nach Einfang negativer Pionen

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Die Spaltwahrscheinlichkeiten schwerer bis mittelschwerer Kerne nach Einfang negativer Pionen wurden am SIN mit der Rotationskammer [1] gemessen. Die Resultate[2] stehen in guter Übereinstimmung mit einem einfachen Modell nach Nix *et al.* [3].

Literatur

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X-ray spectra were taken with a Ge(Li)-detector in coincidence with the π -stop signal of a conventional counter telescope (Fig. 1). The γ -spectra of the produced β -unstable nuclei were measured after defined exposure in the π -beam with a well shielded Ge(Li) detector (Fig. 2). The absolute intensity of the observed nuclear γ -rays in the prompt spectrum was calibrated via the π -X-rays calculated from a π -cascade program [2].

Isotopes found in both the prompt and activation spectra were used to normalize the activation spectra. The number of emitted neutrons and protons is given by the yield of the product nucleus, which is identified by its characteristic γ -rays.

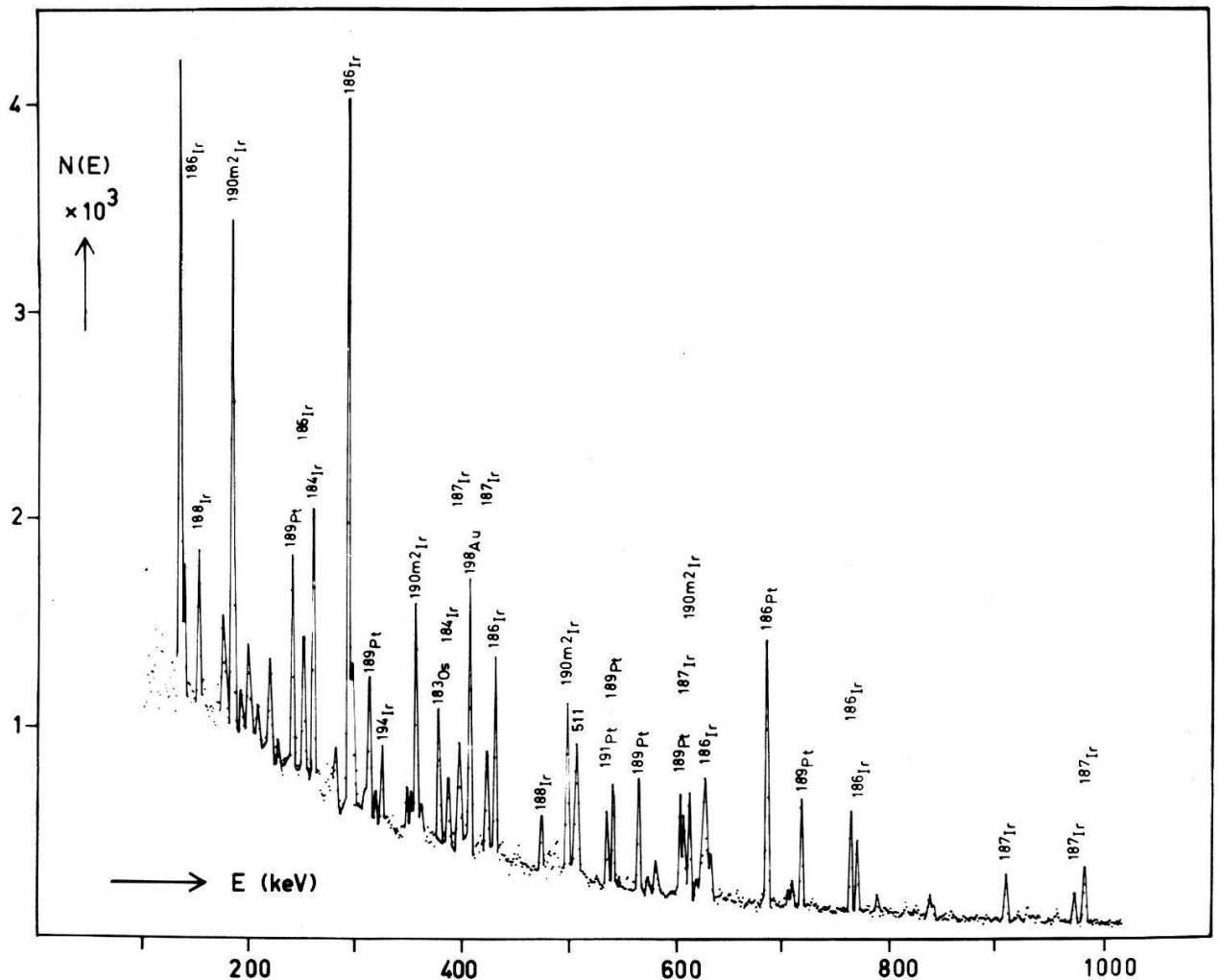


Figure 2

Activation spectrum of π^- -absorption in ^{197}Au . The isotopes identified by the energy of the γ -lines are given.

3. π^- —Absorption in ^{75}As

In the prompt and activation γ -spectra Ge, Ga, and Zn isotopes have been identified, indicating an emission of neutrons together with zero, one or two protons. In the activation γ -spectra also Cu and Ni isotopes have been found. In Figure 3 the yields of the final isotopes are depicted as a function of the number of emitted neutrons and protons. In total the neutron emission is dominant. In the mean $\bar{x} = 4.5$ neutrons and $\bar{y} = 1.2$ protons are emitted. For calcium $\bar{x} = 3.4$ and $\bar{y} = 2.1$ and for niobium $\bar{x} = 4.0$ and $\bar{y} = 1.9$ have been found by Engelhardt *et al.* [3]. The large

proton emission in niobium ($Z = 41$) is in contradiction to our low value for arsen ($Z = 33$), which should be reduced by the Coulomb barrier compared to calcium ($Z = 20$). The extrapolated total yields are $\Sigma(p, xn) = 40\%$, and $\Sigma(xn) = 20\%$.

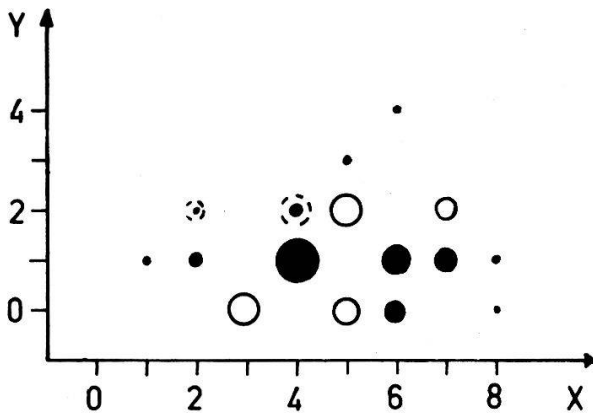


Figure 3

Yields of final isotopes X from the reaction $^{75}\text{As}(\pi^-; xn, \gamma p)X$. The yields are proportional to the area of the points. Open circles are even-even isotopes determined from the prompt γ -spectrum, the black circles are β -unstable isotopes determined from the activation spectrum. For the dotted points (isomeric states), only a lower limit can be given.

4. π^- —Absorption in ^{197}Au

In π^- -absorption in ^{197}Au Pt- and Ir-isotopes have been observed, but isotopes corresponding to the emission of neutrons and two or more protons could not be found. Rotational levels of the even Pt-isotopes easily could be identified (Fig. 1) up to the levels with $I = 10^+$. The relative intensities of the corresponding E2 transitions are plotted in Figure 4 in comparison to results of the $\text{Os}(\alpha, 4n)\text{Pt}^*$ reaction [4]. The excitation of these high-spin states can be explained by the π -absorption on a quasi deuteron and the emission of one fast neutron and absorption of its 180° counter

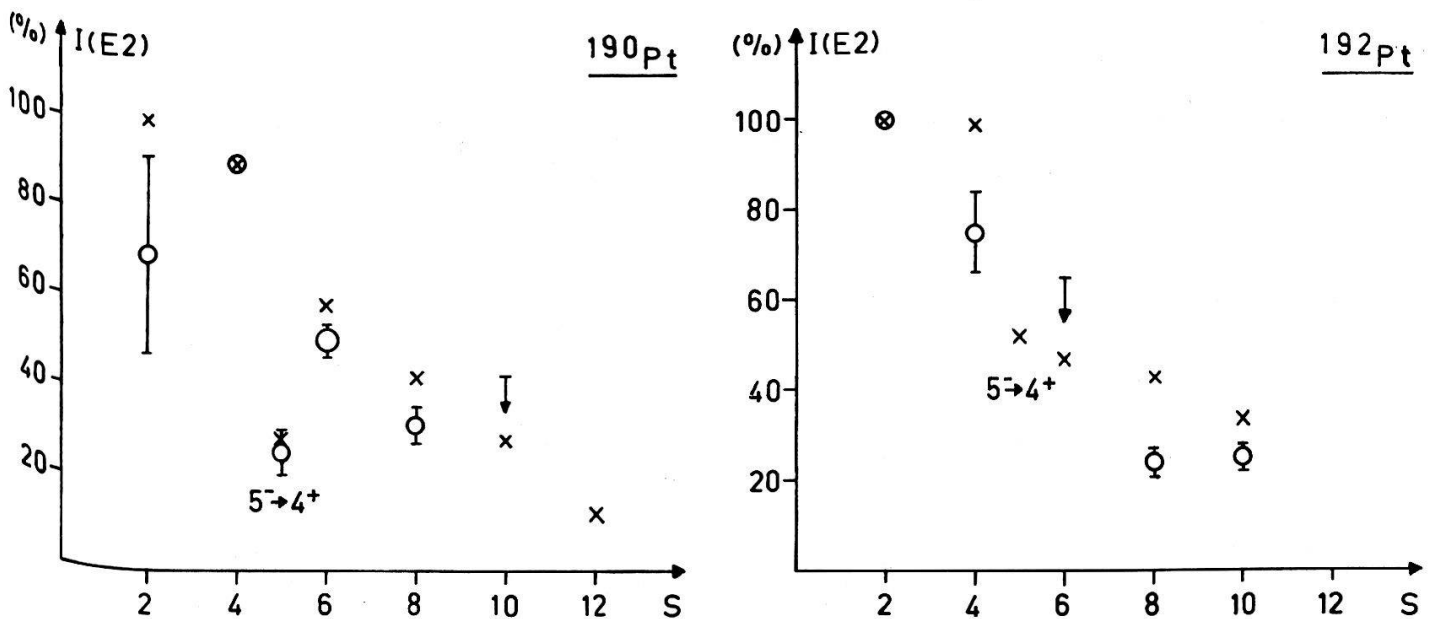


Figure 4

Relative intensities of E2 transitions in ^{190}Pt and ^{192}Pt excited by the reaction $^{197}\text{Au}(\pi^-, xn)\text{Pt}^*$ (circles). The spin of the initial state is given on the abscissa. For comparison the data of the reaction $\text{Os}(\alpha, 4n)\text{Pt}^*$ [4] are given (crosses).

part [5]. From this fact we estimate that about 20% of the neutrons are emitted with an energy of > 15 MeV in agreement with the observed number of fast neutrons after π^- -absorption in ^{165}Ho [6].

The yields of the (π^-, xn) and (π^-, pxn) reactions are plotted in Figure 5. The total yield per stopped π^- for the (π^-, xn) reaction is 87% with a mean neutron emission of $\bar{x} = 6.5$ in general agreement with ^{165}Ho [7].

For the (π^-, pxn) reaction only ^{190}Ir and ^{194}Ir could be measured without ambiguities due to the β -decay of Pt-isotopes. The ratio of the two reactions for 7 nucleons emission is about 6 corresponding to 2% proton emission per π^- stop.

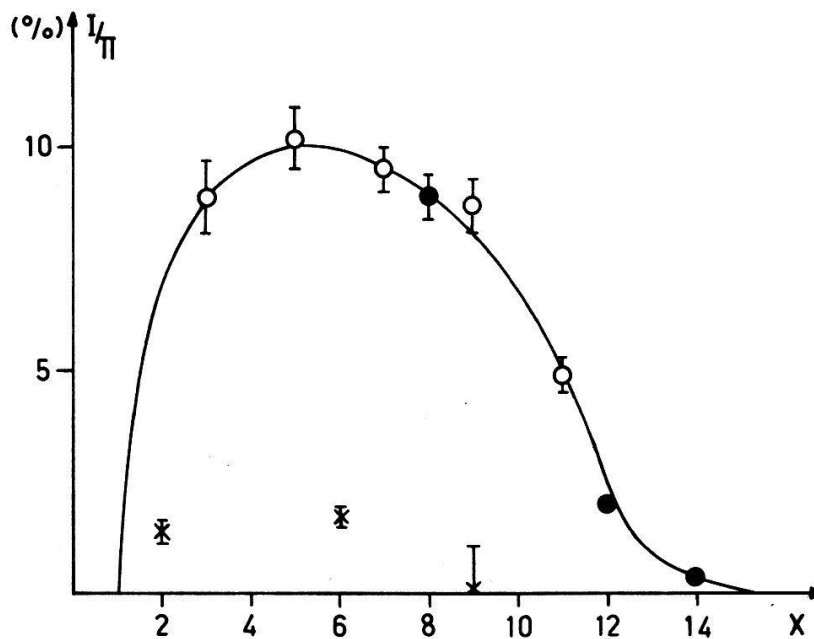


Figure 5

Yields of Pt and Ir isotopes from the $^{197}\text{Au} (\pi^-, xn)^{197-x}\text{Pt}$ (open circles from prompt spectrum, black circles from activation) and $^{197}\text{Au} (\pi^-, pxn)^{196-x}\text{Ir}$ (crosses, from activation) reactions.

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Description du spectromètre de l'Université de Genève et résultats préliminaires de l'analyse de la réaction $\bar{p}p \rightarrow \bar{\Delta}(1236)\Delta(1236)$

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Nous présentons, dans ce rapport une description d'un spectromètre, consistant essentiellement de chambres proportionnelles multifils et d'une système de temps de

vol, qui a été utilisé à l'accélérateur PS du CERN. Dans une deuxième partie, nous présentons quelques résultats préliminaires de l'analyse de la réaction $\bar{p}p \rightarrow \bar{\Delta}(1236)\Delta(1236)$ à 10 GeV/c.

Résultats partiels sur la production de résonances mésoniques dans la réaction $\pi^- p \rightarrow K^0 K^- p$ à 10 GeV/c

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Parallèlement à la recherche de mésons étranges, des événements décrits par la réaction ci-dessus ont été enregistrés à l'aide du spectromètre de l'Université de Genève au cours de l'expérience effectuée durant les années 1974/76 au PS du CERN. On présente les premiers résultats de l'analyse des résonances du système $K^0 K^-$ à masses et spins élevés par la méthode des moments.

Suche nach einem seltenen Myonzerfall

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Am SIN wurde nach dem hypothetischen Zerfall $\mu^+ \rightarrow x^+ \nu$ gesucht. Im Massenbereich von $60 \text{ MeV}/c^2 < m_x < 90 \text{ MeV}/c^2$ ergab das Experiment eine obere Grenze von $\sim 10^{-6}$ für das Verzweigungsverhältnis $\Gamma_{\mu-x\nu}/\Gamma_{\mu}^{\text{tot}}$.

Exchange current effects in pion double charge exchange from ^4He

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The cross section for the reaction $^4\text{He}(\pi^+, \pi^-)4p$ is estimated in a model where the double charge exchange takes place on one of the virtual pions responsible for the nuclear binding [1]. Our calculation takes fully into account the five-body final phase space with the effects of Pauli principle between the four outgoing protons. Angular and momentum distributions as well as total rate are computed from threshold up to 2 GeV/c using Monte Carlo techniques. They will be compared with existing data especially the new ones which are presently being analysed at Neuchâtel [2].

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Double échange de charge π^+ (${}^4\text{He}$, $4p$) π^- entre 1.5 et 2 GeV/c

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10' 165 clichés de la chambre à bulles à hélium d'Oxford/RHEL exposée au faisceau de π^+ de Nimrod ont été explorés et 1' 099 événements candidats (étoiles à 5 branches) ont été mesurés et analysés. Il en est résulté 24 événements du type double échange de charge augmentant la statistique mondiale à 89 événements de ce type. Les sections efficaces de double échange de charge trouvées sont, à chacune des 3 impulsions incidentes :

$$\begin{array}{ll} p_i = 1.46 \pm 0.12 \text{ GeV/c} & \sigma = 0.41 \pm 0.14 \text{ mb} \\ p_i = 1.72 \pm 0.11 \text{ GeV/c} & \sigma = 0.24 \pm 0.08 \text{ mb} \\ p_i = 1.86 \pm 0.13 \text{ GeV/c} & \sigma = 0.14 \pm 0.07 \text{ mb} \end{array}$$

Dans ce domaine d'énergie, le double échange de charge est caractérisé par :

- Une faible proportion d'événements avec deux protons spectateurs (5/24) ou sans spectateur du tout (8/24), un proton étant ici considéré comme spectateur si son impulsion dans le système du laboratoire est inférieure à 300 MeV/c.
- Un transfert d'énergie relativement faible au π^- (environ 30% de l'énergie cinétique à disposition).
- Une production abondante de résonances $\Delta(1'232)$ et $\Delta(1'650)$ (environ une résonance par événement).

Un article plus détaillé est soumis à Nuovo Cimento.

Ionisation der K-Schale bei Schwerionenstößen

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Ueblicherweise werden Berechnungen des Wirkungsquerschnittes für K-Schalenionisation im Rahmen der Bornschen Näherung, d.h. unter Vernachlässigung der Wechselwirkung zwischen Projektil und Targetkern, durchgeführt. Man erhält dabei befriedigende Uebereinstimmung mit den experimentellen Ergebnissen der protoneninduzierten Ionisationsprozesse. In der vorliegenden Arbeit untersuchen wir, inwieweit bei Ionisationsprozessen, die durch schwere Ionen ($Z=4$) induziert werden, die Ablenkung des Projektils durch die Coulombwechselwirkung berücksichtigt werden muss, indem wir den Ionisationsquerschnitt in der halbklassischen Näherung (SCA) berechnen und die Resultate mit der Bornschen Näherung vergleichen.

Anisotropie der quasimolekularen K-Röntgenstrahlung bei Schwerionenkollisionen

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Die Eigenschaften der quasimolekularen K-Röntgenstrahlung wurden an den Systemen Al-Al, Ca-Ca, Cl-Ni, Fe-Fe, Fe-Ni, Ni-Fe, Ni-Ni, Ge-Ge, Br-Nb,

Br-Br und Ag-Ag bei Strahlenergien zwischen 24 und 64.5 MeV untersucht. In allen Fällen wurde eine Anisotropie gefunden, die von der Röntgenenergie abhängt und insbesondere in der Nähe der *K*-Uebergangsenergien im vereinigten Atom ein Maximum durchläuft. Für diesen Effekt gibt es zur Zeit noch keine befriedigende Erklärung.

Rayons de charge nucléaire de Na et S déterminés à partir de rayons X muoniques

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Pour des noyaux légers, on peut déterminer un paramètre de la répartition de charge à partir des transitions atomiques du muon aboutissant à l'état $1s$. Les énergies des transitions ($2p-1s$) ont été mesurées au SIN. Nous avons obtenu 516.344(45) keV pour le soufre et 250.238(30) keV pour le sodium. A l'aide d'une répartition de charge du type de Fermi et en tenant compte des corrections d'ordres supérieurs, on calcule un rayon quadratique moyen $\langle r^2 \rangle^{1/2} = 3.261(5)$ fm pour S et 2.99(1) fm pour Na. Pour le rayon équivalent R_k déterminé par l'intermédiaire d'un moment généralisé on obtient $R_{2,11}(S) = 4.186(4)$ fm et $R_{2,11}(Na) = 3.842(11)$ fm.

Stopped negative muon-induced fission of ^{238}U

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The mass-distribution from the stopped negative muon-induced fission of ^{238}U is being measured by radiochemical methods. The following fission products were determined: ^{91}Sr , ^{93}Y , ^{99}Mo , ^{111}Ag , ^{112}Pd , ^{113}Ag , ^{115}Cd , ^{132}Te , ^{140}Ba , ^{141}Ce , ^{143}Ce , ^{147}Nd , ^{149}Pm , ^{151}Pm and ^{153}Sm . The mass-distribution is asymmetric; the position and the shape of the heavy-mass peak are very similar to those in 14 MeV neutron-induced fission of ^{238}U , whereas a slight shift towards lighter masses is indicated for the light-mass peak. The peak-to-valley ratio is about 3 compared to 9 in 14 MeV neutron-induced fission of ^{238}U . The fission probability per stopped muon is 0.14 ± 0.03 assuming a fission yield of 0.038 ± 0.004 for ^{143}Ce .

New investigations on the empirical mass formula $m \approx N \times 3m_e$ for 'stable' particles

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1. Introduction

The rest masses of particles stable against strong decays and heavier than the electron (μ^\pm , π^0 , π^\pm , K^\pm , K^0 , η , p , n , Λ , Σ^+ , Σ^0 , Σ^- , Ξ^0 , Ξ^- , Ω^-) tend to be close to

integral multiples of $3m_e$, the threefold electron rest mass. This empirical mass formula, which was first discussed in 1965 [1], has now been tested for the most recent mass values of the Particle Data Group [2]. The tests have been improved, mainly by choosing a more satisfactory definition of the weight factors introduced in order to reduce the influence of inaccurate experimental mass values.

In the following, a detailed discussion will be given for the new results only; a more complete treatment can be found in ref. [1].

2. The evidence for a 'mass quantum' $3m_e$

The most recent mass values m_i for the particles under consideration are listed in Table 1. The number N_i in the third column is the integer closest to $m_i/(3m_e)$, and x_i is the deviation from an exact 'quantum law':

$$x_i = \frac{m_i}{(3m_e)} - N_i. \quad (1)$$

As discussed in [1], the expectation value of the absolute magnitude $|x_i|$ for *randomly distributed* mass values m_i is very close to 0.25. The experimental deviations $|x_i|$ show a clear tendency to lie below that expectation value.

Table 1
Experimental values of N_i , x_i and w_i for the 15 stable particle masses m_i [2] greater than the electron mass

Particle	m_i (MeV/c ²)	N_i	x_i	w_i for $\Delta m_i = 0.1$ MeV/c ²	w_i for $\Delta m_i = 0.2$ MeV/c ²
μ^\pm	105.659 48 \pm 0.000 35	69	-0.077 12 \pm 0.000 30	1.000	1.000
π^0	134.964 5 \pm 0.007 4	88	0.038 9 \pm 0.004 8	0.997	0.999
π^\pm	139.568 8 \pm 0.006 4	91	0.042 3 \pm 0.004 2	0.998	0.999
K^\pm	493.707 \pm 0.037	322	0.051 \pm 0.024	0.935	0.983
K^0	497.70 \pm 0.13	325	-0.345 \pm 0.085	0.483	0.816
η	548.8 \pm 0.6	358	-0.01 \pm 0.39	0.028	0.111
p	938.279 6 \pm 0.002 7	612	0.0505 \pm 0.0025	1.000	1.000
n	939.573 1 \pm 0.002 7	613	-0.105 8 \pm 0.002 5	1.000	1.000
Λ	1115.60 \pm 0.05	728	-0.281 \pm 0.033	0.885	0.969
Σ^+	1189.37 \pm 0.06	776	-0.160 \pm 0.039	0.840	0.956
Σ^0	1192.47 \pm 0.08	778	-0.138 \pm 0.052	0.739	0.924
Σ^-	1197.35 \pm 0.06	781	+0.045 \pm 0.039	0.840	0.956
Ξ^0	1314.9 \pm 0.6	858	-0.28 \pm 0.39	0.028	0.111
Ξ^-	1321.29 \pm 0.14	862	-0.107 \pm 0.091	0.438	0.791
Ω^-	1672.2 \pm 0.4	1091	-0.20 \pm 0.26	0.062	0.245

N_i = integer closest to $m_i/(3m_e)$.

x_i = deviation from exact quantum law; $x_i = m_i/(3m_e) - N_i$; $m_e = (0.511\,003\,4 \pm 0.000\,001\,4)$ MeV/c².

w_i = new weight factors for significance tests.

Δm_i = uncertainty limit used to calculate w_i ; see equation (3).

3. The weight factors w_i

The significance of the experimental evidence is discussed in terms of the *mean deviation* D of the 15 mass values from an exact quantum law:

$$D = \frac{\sum_{i=1}^{15} w_i \times |x_i|}{\sum_{i=1}^{15} w_i}. \quad (2)$$

The weight factors w_i as used in the earlier studies [1] are, for large uncertainties Δm , inversely proportional to Δm , and this long 'tail' lead to criticism: Inaccurate mass values which accidentally have a small deviation $|x_i|$ may have too much weight. Indeed for the 1973 mass values [1] both the η and the Ω^- had deviations of 0.01 with uncertainties of 0.39 and 0.33, respectively.

The main purpose of the present note is to show that the conclusions are not altered significantly by improvements in the definition of the weight factors.

If one would propose an *exact* quantum law, one would have to use weight factors $(1/\Delta m_i)^2$; however it is evident from Table 1 that the law can be at best approximate: The deviations $|x_i|$ are of the order of 0.05 or 0.1. Therefore an uncertainty Δx_i of that magnitude, i.e. a mass uncertainty around 0.1 MeV/c², is good enough for our purposes. For mass uncertainties large compared to 0.1 MeV/c², the weight factors should thus be proportional to $(1/\Delta m_i)^2$; for smaller uncertainties, w may be allowed to level off; with such a definition one prevents the tests from being dominated by very few particles as it would happen with the choice $(1/\Delta m_i)^2$ for all values of Δm_i .

The new definition is

$$w = (1 - e^{-\xi^2})/\xi^2, \quad (3)$$

where $\xi = \Delta m/\Delta m_i$; here, Δm is the particle mass uncertainty, and Δm_i is a fixed 'uncertainty limit'; the tests were done for $\Delta m_i = 0.1$ and 0.2 MeV/c². The corresponding new weight factors for the individual particles are presented in Table 1.

4. Results of the significance tests

The new weight factors have been used to redo the significance tests discussed in ref. [1], for the most recent mass values [2].

(4a) Test with $M_0 = 3.0 m_e$

For $\Delta m_i = 0.1$ (0.2) MeV/c², the 'experimental' mean deviation D , as defined by equation (2) above, has the value

$$D_{\text{exp.}} = 0.109 \text{ (0.119)}. \quad (4)$$

Among 200.000 sets of 15 *fictitious* mass values randomly distributed in a suitable range [1], only 55 (75) had a mean deviation smaller than that of equation (4). So, the probability for random mass values to fit our formula better than the experimental mass values, is very small ($\sim 3 \times 10^{-4}$).

(4b) Search for alternative 'mass quanta'

By replacing, in equation (1), the 'quantum' $3m_e$ by an arbitrary value M_0 , a general definition of the mean deviation D is obtained. The weight factors w_i were still taken from Table 1; this may be objected to, because the size of the uncertainty

limit Δm_i in equation (3) is related to the mass quantum (Δm_i should be small compared to M_0); however weight factors changing with M_0 were felt to be even less satisfactory.

For $\Delta m_i = 0.1 \text{ MeV}/c^2$, the result of the search in the neighborhood of $3m_e$ is shown in Figure 1. There is a deep minimum of D very close to $3.0m_e$:

$$D = 0.089 \text{ at } M_0 = 2.999\,48m_e. \quad (5)$$

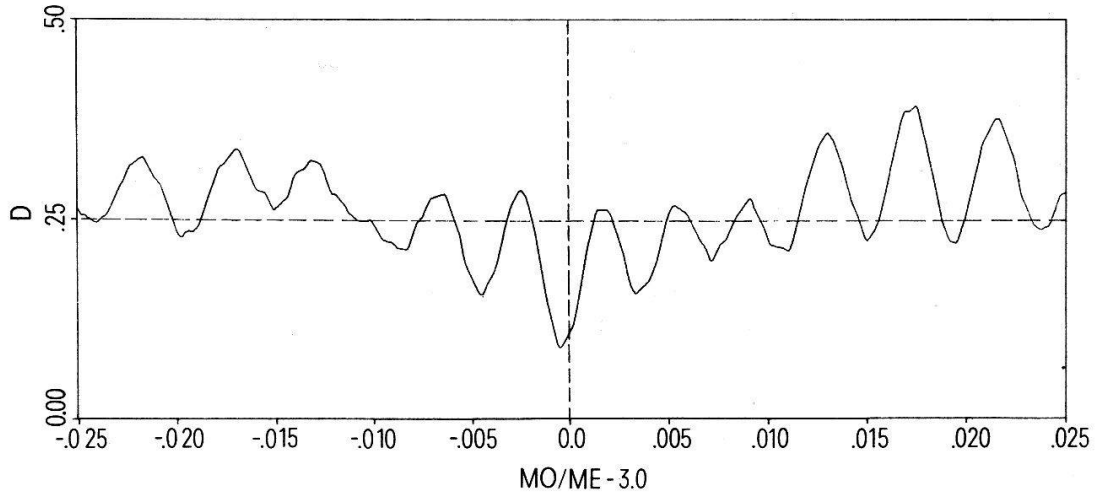


Figure 1

Result of the search for a mass quantum M_0 in the neighborhood of $3.0m_e$. Abscissa: $(M_0/m_e) - 3$; ordinate: Mean deviation D of the 15 stable particle masses from nearest multiple of M_0 ; see equation (2). There is a deep minimum of D very close to $3.0m_e$: $D = 0.089$ at $M_0 = 2.999\,48m_e$. Weight factors w_i for this plot calculated with 'uncertainty limit' $\Delta m_i = 0.1 \text{ MeV}/c^2$; see equation (3) and Table 1.

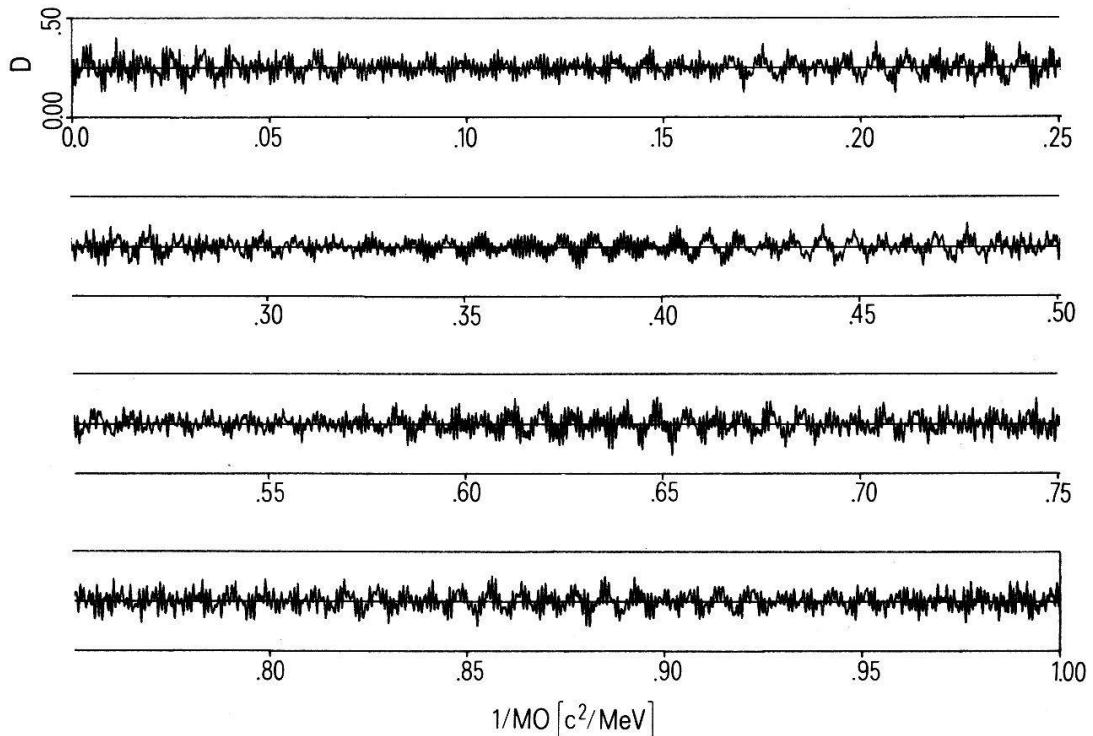


Figure 2

Same as Figure 1, but for a much larger range of M_0 -values. Abscissa here is $(1/M_0)$, rather than M_0 , so that period of oscillation is constant. There is a 'trivial' zero of D at $(1/M_0) = 0$ (see text); apart from that, the deepest minimum is still $D = 0.089$ at $(1/M_0) = 0.652\,42 \text{ c}^2/\text{MeV}$, i.e. at $M_0 = 1.532\,74 \text{ MeV}/c^2 = 2.999\,48m_e$. The *second-deepest* minimum of this plot is $D = 0.114$ at $(1/M_0) = 0.881\,09 \text{ c}^2/\text{MeV}$, $M_0 = 1.134\,96 \text{ MeV}/c^2 = 2.220\,91 m_e$. The search was continued for $(1/M_0) > 1.0 \text{ c}^2/\text{MeV}$; see text.

A plot with a much larger M_0 -range has also been obtained; it is reproduced in Figure 2. Here D is plotted as a function of $(1/M_0)$, so that the period of oscillation is constant. For small values of $(1/M_0)$, i.e. for $3m_e < M_0 < \infty$, there is no minimum of D deeper than that of equation (5), except the trivial zero at $(1/M_0) = 0$, where the integer N_i (see Table 1) is zero for all particles.

The search was also extended in the other direction, i.e. to $M_0 < 3m_e$, or $(1/M_0) \rightarrow \infty$. Mass quanta which are not large compared to most uncertainties Δm_i should not be considered; therefore the search has to be stopped at $M_0 \approx 1 \text{ MeV}/c^2$. In the region where $1.0 \text{ MeV}/c^2 < M_0 < 2000 \text{ MeV}/c^2$ there is no minimum of D deeper than that of equation (5). The *second-deepest* minimum in that region, still for $\Delta m_i = 0.1 \text{ MeV}/c^2$, is $D = 0.114$ at $M_0 = 2.22091m_e$.

If one does continue to search below $1.0 \text{ MeV}/c^2$, one finds a minimum slightly deeper than that of equation (5), $D = 0.088$, at $M_0 = 1.30030m_e$. Below that, one has to go down to $M_0 = 0.318664m_e$ to find the next such minimum ($D = 0.087$).

For weight factors with $\Delta m_i = 0.2 \text{ MeV}/c^2$ the results are even more convincing: For $1.0 \text{ MeV}/c^2 < M_0 < 2000 \text{ MeV}/c^2$, the deepest minimum is $D = 0.091$ at $2.99948m_e$; second-deepest minimum: $D = 0.117$ at $3.07054m_e$. Below $1.0 \text{ MeV}/c^2$ there is no minimum deeper than 0.091 down to very small M_0 -values: $D = 0.080$ at $M_0 = 0.1950655m_e$. As stated before, such a small mass quantum should not be considered because it is not large compared to most mass uncertainties.

(4c) *Is the minimum near $3.0m_e$ strikingly deep?*

The best candidate for a mass quantum in a large range (1.0 to $2000 \text{ MeV}/c^2$) is alarmingly close to $3.0m_e$. Would this minimum be unusual even if it were not close to any special mass value? As in ref. [1] we generated 100 sets of 15 fictitious mass values and for each set we redid the above search in the range from 1.0 to $2000 \text{ MeV}/c^2$, again for $\Delta m_i = 0.1$ (0.2) MeV/c^2 . For 24 (14) of the 100 sets there was at least one minimum of D deeper than the "experimental" minimum at $2.99948m_e$. So the interesting feature of the experimental minimum is its M_0 -value, while its D -value, taken by itself, is rather low, but not extraordinarily so.

The three significance tests for the 1976 mass values were compared with those for the 1973 values, both with the old and the new definition of the weight factors; from these comparisons it follows that the overall evidence has stayed about the same in the past three years.

5. Conclusions

The conclusions of ref. [1] have not been altered significantly by recent changes of the experimental particle mass values, nor by the introduction of the new weight factors. The best candidate for a mass quantum in a large range (1.0 to $2000 \text{ MeV}/c^2$) is still very close to $3m_e$, namely at $2.99948m_e$.

More accurate mass measurements for K_0 , η , Ξ^0 , Ξ^- , and Ω^- are still very desirable.

References

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Charm exists—what next¹⁾

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Abstract. The most recent data on D -meson production in e^+e^- annihilation are analyzed by a simple model. The experimental and theoretical implications of the discovery of charmed particles are explained.

1. A 'charmed' new world

The recent discovery of narrow states in the invariant mass spectra of the $K\pi$, $K2\pi$, and $K3\pi$ decay channels around 2 GeV in electron-positron initiated reactions [1] is commonly considered as conclusive evidence for the existence of a new quark which has been anticipated for a long time [2]. This piece of evidence does not leave much doubt about the existence of quarks and in particular of this new charmed quark [3]. It is therefore natural to ask what consequences one may expect on the experimental as well as on the theoretical side from this discovery.

In this note we present a short 'tour d'horizon' on the recent data [1] on the theoretical questions arising from it [4] and from the presently emerging picture of particle dynamics [5], and on the experiments which are expected to give further information on the characteristics of the charmed quark [6]. Our aim is to expose the immediate impact of the discovery of a charmed quark on the theoretical and the experimental side and to point to possible further discoveries one may expect as we go to higher energies.

In Section 2 we sketch the anticipated framework for a description of interacting particles and review the arguments leading to the postulate of a charmed quark. Some earlier and the most recent experimental results giving evidence for such a new quark are discussed in Section 3. The models and questions arising from the most recent data on charmed particle production are touched on in Section 4. What experimental implications are expected from the existence of a charmed quark? We attempt an answer to this question in Section 5. In Section 6 we discuss the theoretical implications of the charmed quark and point to missing pieces in the overall picture of particle dynamics. Section 7 presents our conclusions.

2. Theoretical motivation for a new quark

Before presenting the motivations for introducing a new quark, we first take a brief look at the present picture of particle dynamics.

The weak and electromagnetic interactions have been successfully unified in a gauge theory [7], admitting charged as well as neutral currents, which does not violate the unitarity bound at large energies. It involves leptons, quarks, and gauge bosons which group into doublets and triplets, respectively.

The structure of non-Abelian gauge theories has many attractive features believed to be relevant in strong interaction dynamics, such as renormalizability [8], asymptotic freedom [9], possibly quark confinement [10], such that strong inter-

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action dynamics is believed to be dominated by a field theory of this type too [5]. There thus appears a strong similarity between strong and weak interactions in the sense:

$$\begin{aligned} \text{weak + e.m. interactions} &\leftrightarrow \text{strong interactions} \\ \text{leptons} &\leftrightarrow \text{quarks} \\ \text{gauge bosons} &\leftrightarrow \text{gauge gluons} \\ \text{?} &\leftrightarrow \text{Pomeron} \\ \text{Higgs bosons} &\leftrightarrow \text{?} \end{aligned}$$

The above-sketched weak interaction framework does permit strangeness changing neutral currents $\Delta s = 1$ which, however, are not compatible with the experimental results. First clues on a charmed quark came from the study of the strangeness changing second-order weak processes, such as $K_L \rightarrow \mu^+ \mu^-$, $K^+ \rightarrow \pi^+ \nu \bar{\nu}$ and the $K_L K_S$ mass difference. In a unified gauge theory of weak and electromagnetic interactions, the magnitude of a second order weak amplitude is in general $G_F \cdot \alpha$, so, in order to explain the observed magnitude of the $K_L \rightarrow \mu \mu$ amplitude which experimentally is of the order $G_F \cdot \alpha^2$, a suppression mechanism is needed. In the Weinberg-Salam model, the charmed quark indeed does remove the strangeness changing neutral current effects in first order and in higher orders as well; this was the main reason for introducing a new quark [3].

3. Experimental evidence for a charmed quark

In this section we present briefly the early experimental evidence for the existence of a charmed quark and subsequently discuss the most recent discoveries of charmed mesons and baryons.

(A) Earlier indications

Further evidence on the existence of a new quark came from the rising ratio $R \equiv \sigma(e^+ e^- \rightarrow \text{hadrons})/\sigma(\mu^+ \mu^-)$ in electron-positron initiated reactions [11] although at the time of the first experimental evidence of this fact many alternative explanations were offered such as color, statistical models, new Pomeron-like interactions, and others [12]. So far, the strongest piece of evidence in favor of charmed quarks came from the new resonances ψ and ψ' and their radiation transitions, although many alternative explanations seemed to be possible in the time just after their discovery [12]. Experimental information from electron-positron annihilation reactions has meanwhile accumulated to such an extent [13] that the existence of a charmed quark was almost undeniable and it was considered a matter of time until charmed mesons, bound states composed of a charmed quark and a conventional (u, d, s) quark, would be discovered.

(B) Recent discoveries

Indeed experimental evidence for narrow bumps in the $K\pi$, $K3\pi$, and $K2\pi$ channels of electron-positron initiated reactions now exists [1] and there is little doubt that the expected charmed mesons D^0, D^\pm , as well as $D^{*0}, D^{*\pm}$ [3], have been discovered. There is further evidence of charmed baryons $C_0 \equiv (cud)$ and $C_1 \equiv (cuu)$ in photoproduction [14].

In the following we present the characteristics of the data which seem to emerge.

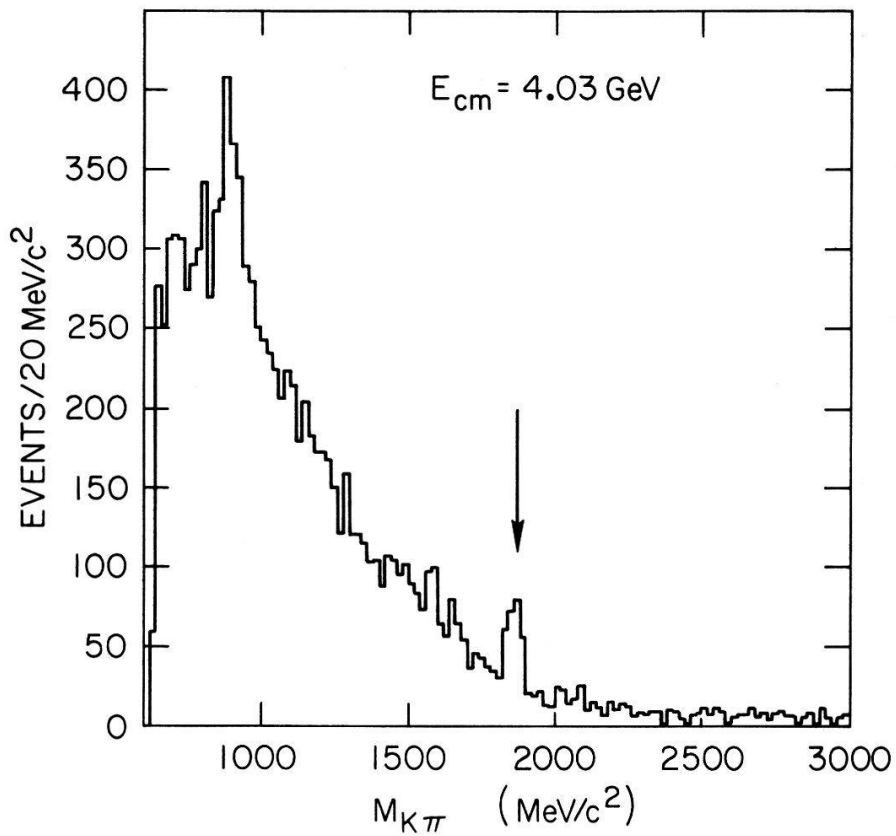


Figure 1
Invariant mass spectrum of the $K\pi$ channel in e^+e^- annihilation at $E_{c.m.} = 4.03$ GeV.

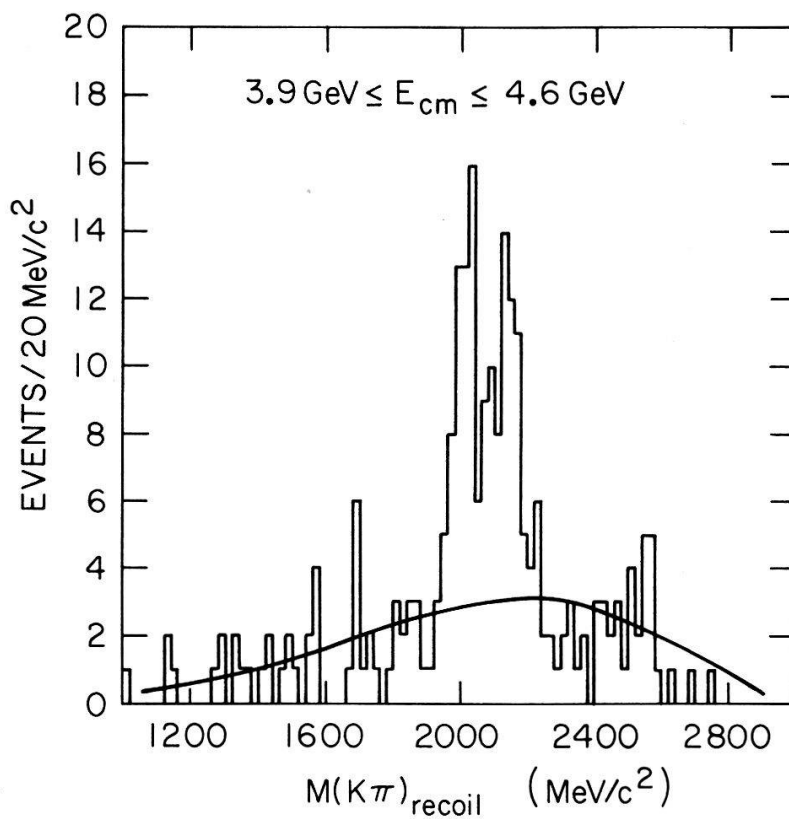


Figure 2
 $D^0(K\pi)$ -recoil spectrum in e^+e^- annihilation in the energy range $3.9 \text{ GeV} \leq E_{c.m.} \leq 4.6 \text{ GeV}$.

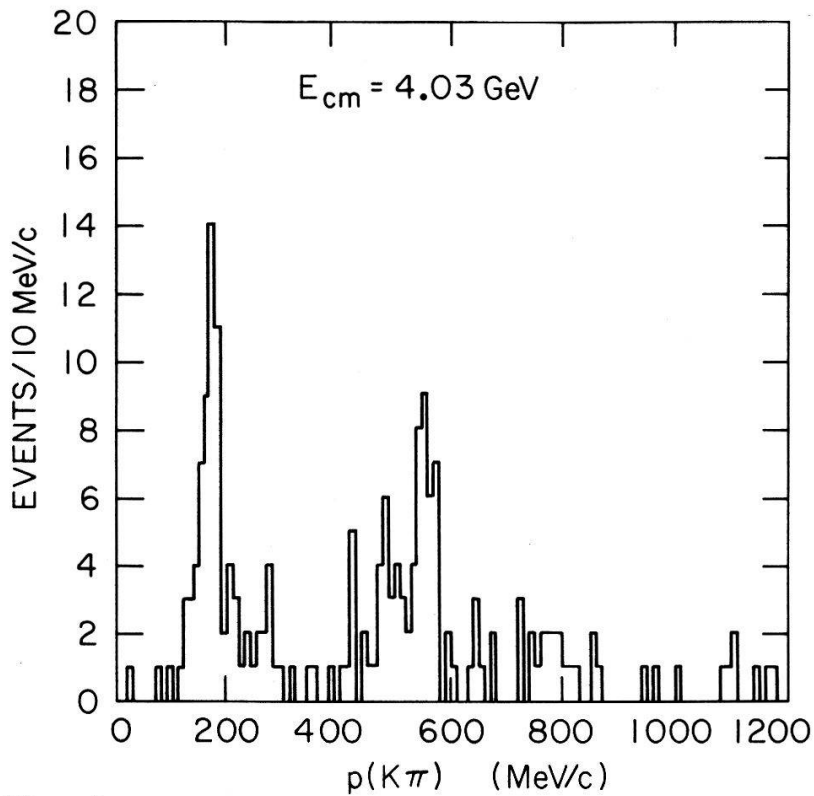


Figure 3
 $D^0(K\pi)$ - momentum spectrum at $E_{c.m.} = 4.03$ GeV.

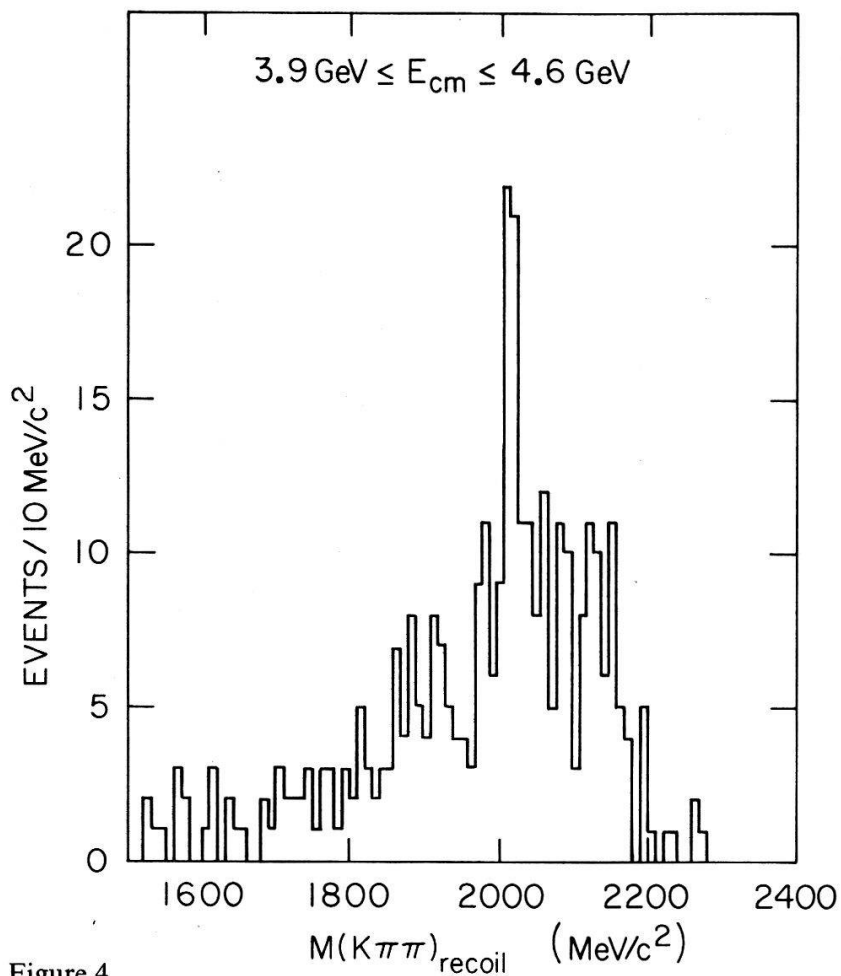


Figure 4
 $D^\pm(K\pi\pi)$ - recoil spectrum in e^+e^- annihilation in the energy range $3.9 \text{ GeV} \leq E_{c.m.} \leq 4.6 \text{ GeV}$.

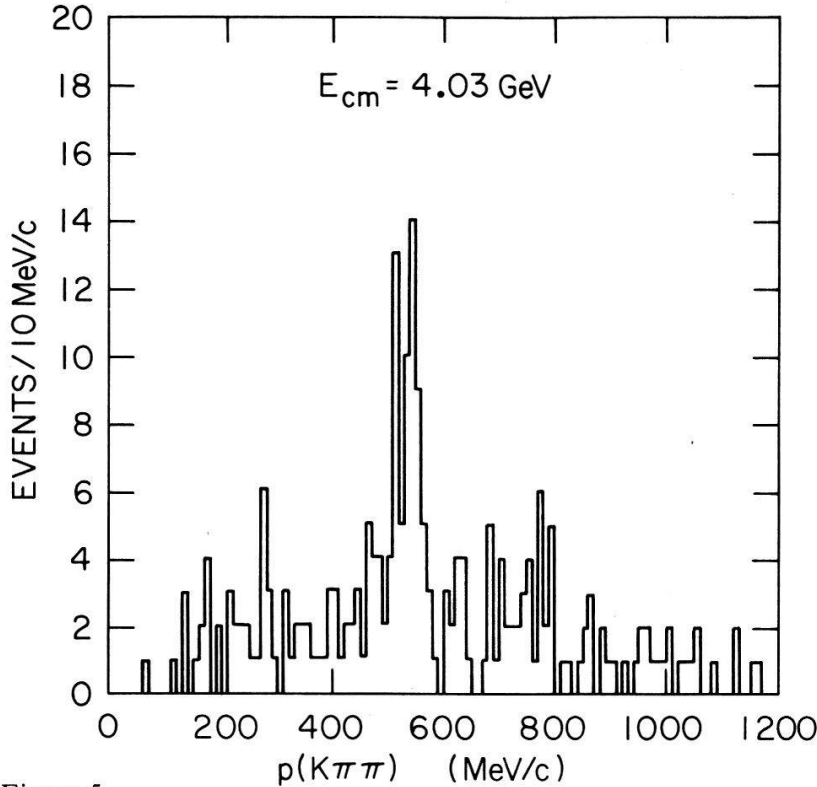


Figure 5
 $D^\pm(K\pi\pi)$ - momentum spectrum at $E_{c.m.} = 4.03$ GeV.

The $K^\pm\pi^\mp$ and $K^\pm\pi^\mp\pi^\pm\pi^\mp$ invariant mass spectra show a peak at 1.87 GeV/ c^2 which is assigned to the neutral charmed meson $D^0 \equiv (c\bar{u})$; its mass is 1865 ± 15 MeV/ c^2 (Fig. 1). Similarly a further narrow state was reported in the invariant mass distribution of the $K^\mp\pi^\pm\pi^\pm$ exotic decay channel with mass 1876 ± 15 MeV/ c^2 . In charm spectroscopy this state is assigned to be the charged partner $D^+ \equiv (c\bar{d})$ of the above neutral charmed meson D^0 . Both of the above states appear in association with a system having a mass of approximately 2 GeV. The $K\pi$ recoil spectrum shows two pronounced peaks at ~ 2010 MeV/ c^2 and ~ 2150 MeV/ c^2 of roughly equal height with an almost equal number of events under the two peaks (Fig. 2). Below and above this region events are scattered, which is partially attributed to background; one might speculate on the existence of a small peak at ~ 2600 GeV/ c^2 . The $K3\pi$ recoil spectrum shows the same characteristics with the 2150 MeV/ c^2 peak more pronounced. The D^0 -momentum spectrum shows two narrow peaks around ~ 180 MeV/ c and ~ 150 MeV/ c which are interpreted as reflections of the $D^{0*}\bar{D}^{0*}$ and $D^0\bar{D}^{0*}$ production channels respectively (Fig. 3). The $D^{0*} \rightarrow D^{\pm 0}\pi^{\mp 0}$ decay modes certainly exist and if $D^{0*} \rightarrow D^0\gamma$ is important it will lead to a broadening of the momentum spectrum. The analogous recoil-mass and momentum spectra for charged D^\pm -production are shown in Figures 4 and 5. One notices that the $D^\pm D^{*\mp}$ reflection is more pronounced in comparison to the $D^{*\pm}D^{*\mp}$ peak and a clear $D^{*\pm}D^{*\mp}$ signal in the momentum spectrum is missing.

4. Theoretical modeling and problems

The experimental results have recently been analyzed by De Rujula, Georgi, and Glashow [15] and by Eichten and Lane [16] by model calculations. Both groups qualitatively account for the striking predominance of the associated D^* production over D (at $\sqrt{s} = 4.03$ GeV) by a sequential production of the quark pair in which

the more massive $c\bar{c}$ contribution is produced initially through the virtual photon and subsequently an uncorrelated pair of lighter quarks is produced having no direct interaction with the photon. The ratios $\sigma_{D\bar{D}}:\sigma_{D\bar{D}^*+\bar{D}D^*}:\sigma_{D^*\bar{D}^*} = 1:4:7$ for both the neutral and charged cases are then obtained using the above assumptions, angular momentum conservation, and the traditional method of counting statistical weights of the allowed final angular momentum states with the implied assumption that the electromagnetic coupling to each allowed spin state is equivalent.

The model in Ref. [15] parametrizes the integrated cross sections or the resulting ratios, $R \equiv \sigma_h/\sigma_\mu$, by their threshold rise and a subsequent falloff due to the form factor which accounts for binding effects between the quarks. Their form is:

$$R_{D\bar{D}} \propto 1 \cdot p^{3/2} \cdot e^{-\gamma p^2} \quad (4.1)$$

$$R_{D\bar{D}^*+\bar{D}D^*} \propto 4 \cdot p^{3/2} \cdot e^{-\gamma p^2} \quad (4.2)$$

$$R_{D^*\bar{D}^*} \propto 7 \cdot p^{3/2} \cdot e^{-\gamma p^2} \quad (4.3)$$

$$p = \sqrt{\frac{[s - (m_1 + m_2)^2][s - (m_1 - m_2)^2]}{4s}} \quad (4.4)$$

where p is the c.m. momentum of the produced charmed meson pair with masses m_1 and m_2 . In Figure 6 we have drawn the curves resulting from equations (4.1–4.3) for the above three channels with the mass values given by Goldhaber *et al.* [1]:

$$\begin{aligned} m_{D^0} &= 1865 \pm 15 \text{ MeV}/c^2, & m_{D^+} &= 1876 \pm 15 \text{ MeV}/c^2 \\ m_{D^{0*}} &= 2007 \pm 20 \text{ MeV}/c^2, & m_{D^{+*}} &= 2010 \pm 20 \text{ MeV}/c^2. \end{aligned} \quad (4.5)$$

In Figure 6(a) we show the cross-section shapes without any form factor damping ($\gamma = 0$, dash-dotted curves); subsequently we introduce the exponentially falling form factor with different values for the parameter γ . $\gamma = 4 \text{ GeV}^{-2}$ (short dashed curves) corresponds to the value $\Gamma \sim 1 \text{ GeV}^2$ used in Ref. [15]. The sensitive dependence of these curves on the charmed particle masses is exhibited by the shaded area. Its left boundary corresponds to $m_{D^0} = 1865 - 15 \text{ MeV}$, $m_{D^{0*}} = 2007 - 20 \text{ MeV}$, and its right boundary was determined using $m_{D^0} = 1865 + 15 \text{ MeV}$, $m_{D^{0*}} = 2007 + 15 \text{ MeV}$. One notices in particular that the intersection of the left boundary with the vertical line at $E_{\text{c.m.}} = 4.03 \text{ GeV}$ is at almost the same height as the DD^* cross section. $D\bar{D}$ production is small in comparison to $D\bar{D}^*$ or $D^*\bar{D}^*$ production; it peaks around $E_{\text{c.m.}} = 3.9 \text{ GeV}$ and gradually falls off with increasing energy. $D\bar{D}^*$ production dominates around 4.05 GeV where the $D^*\bar{D}^*$ mode has its threshold onset. Due to its steep rise $D^*\bar{D}^*$ production is very sensitive to the initial c.m. energy as well as the D^* mass and dominates D production around 4.2 GeV through $D^* \rightarrow D\pi$ decay. Introducing a stronger exponential damping with $\gamma = 16 \text{ GeV}^{-2}$ (long dashed curves) shifts the maxima of all three D -production cross sections substantially to the left. Figure 6(b) compares the cross-section shapes in equations (4.1–4.3) with the experimental $E_{\text{c.m.}}$ -dependence of $R \equiv \sigma_h/\sigma_\mu$. We notice structure in the region 3.8–3.9 GeV which cannot be explained by the above ansatz although the model predicts some $D\bar{D}$ -production in that area. The bump between 3.9 and 4.0 GeV can be explained by the threshold onset of DD^* -production. The point at 4.03 GeV is measured with high accuracy and there is little room left for an interpretation excluding a resonance (with width $\sim 20 \text{ MeV}$). There is a further bump with maximum at 4.1 GeV which, according to the model, must be identified with the

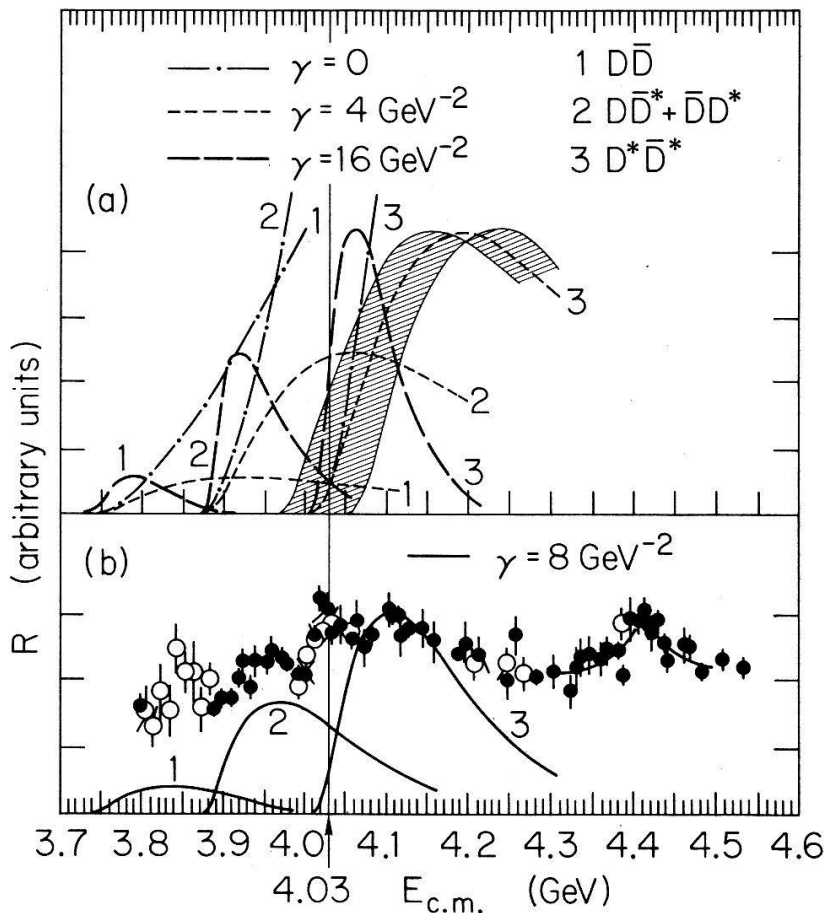


Figure 6

Cross section ratio $R \equiv \sigma_h/\sigma_\mu$ of the channels $D\bar{D}$, $D\bar{D}^*$, $D^*\bar{D}^*$ in the model of Ref. [15] near threshold.

$D^* \bar{D}^*$ threshold onset. The solid lines in Figure 6(b) correspond to $\gamma = 8 \text{ GeV}^{-2}$. At $E_{\text{c.m.}} = 4.028 \text{ GeV}$ (where most of the data have been taken) the production ratios are

$$R_{D^0\bar{D}^0} : R_{D^0\bar{D}^{0*} + \bar{D}^0 D^{0*}} : R_{D^{0*}\bar{D}^{0*}} = 1 : 4.5 : 0.7 \quad (4.6)$$

We conclude that the $D\bar{D}^*$ mode outweighs the $D^*\bar{D}^*$ mode by a factor 6 at this energy point which is clearly in disagreement with the D^0 recoil spectrum; it predicts equal amounts of $D^* \bar{D}^*$ and $D\bar{D}^*$ production, roughly. Variation of γ can shift the peaks of the curves and modify the relative contributions at $E_{\text{c.m.}} = 4.03 \text{ GeV}$; however, reasonable choices cannot explain the recoil spectrum since the cross section for $D^*\bar{D}^*$ is substantially smaller than the cross section for $D\bar{D}^* + \bar{D} D^*$. Ways out of this discrepancy are:

1. The peak at $E_{\text{c.m.}} = 4.03 \text{ GeV}$ in the experimental data is due to the $D^*\bar{D}^*$ threshold onset. Such interpretation, however, has to explain the large value of R and shifts the $D\bar{D}^*$ peak to $E_{\text{c.m.}} = 3.92 \text{ GeV}$, which is slightly too low (see Fig. 6(a)). In addition, the 4.1 GeV peak lacks interpretation.
2. One can assume that the $D\bar{D}^*$ channel has a strong influence and is mainly responsible for the second peak in the recoil spectrum [17]; however, the D^{**} mass appears then unexpectedly low.
3. A resonance at 4.03 GeV which favors substantially the $D^*\bar{D}^*$ mode might be another explanation; dynamical and/or higher symmetry motivations for such a point of view are lacking so far.

4. The reasoning that a shift of the D^{0*} mass does substantially change the threshold onset of the $D^*\bar{D}^*$ channel (see Fig. 6(a)) can be excluded by the following arguments. In Figure 7 we have drawn the s -dependence of the momenta using the mass values given in equation (4.5). We subsequently varied the masses of the D -mesons by adding and subtracting their error values (shaded areas). One notices that $p_{D^*D^*}$ varies very sensitively around $E_{c.m.} \sim 4.03$ GeV. The momentum spectrum (Fig. 3) of the D -meson shows a peak around ~ 180 GeV/c which fixes the momentum value and in turn the mass m_{D^*} . This method can also be applied to p_{DD^*} and leads to the same value $m_{D^*} = 2007$ MeV/c².
5. The most likely explanation is to assume that the form factors depend very sensitively on the charmed particle's momentum.

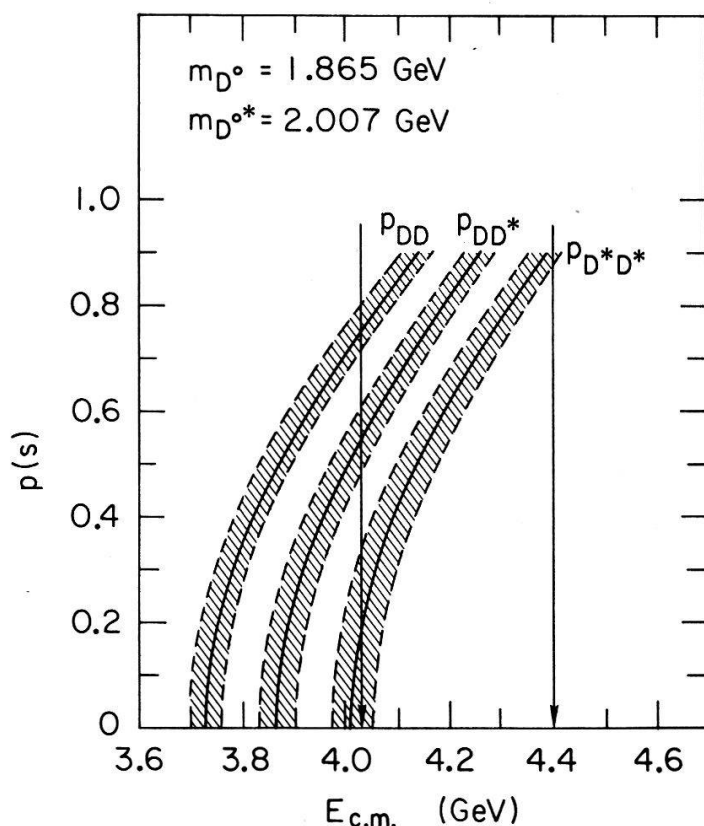


Figure 7

Energy and (D -meson) mass dependence of the c.m.-momenta corresponding to the channels $D\bar{D}$, $D\bar{D}^*$ and $D^*\bar{D}^*$.

We would like to add a brief comment on the widths of the bumps in the recoil spectrum (Figs. 2 and 4). As the c.m. energy increases the recoil bumps due to reflection become broader; their lower boundaries vary relatively little whereas their upper boundaries increase. The amount of broadening depends on the exact masses of D and D^* (see Fig. 8) in particular at larger values of $E_{c.m.}$. The widths of the recoil peaks thus provide further checks.

In the model presented above the photon coupling to the lighter $q\bar{q}$ -pair and subsequent c -quark association was ignored; this assumption was motivated by the phenomenological fact that the creation of a quark-antiquark pair out of the vacuum is less likely with increasing quark mass. Eichten and Lane [16] in their coupled channel model find a suppression of $\sim (m_q/m_c)^4$ for diagram (b) in Figure 9; this result is however only reliable for small charmed meson momenta where their theory

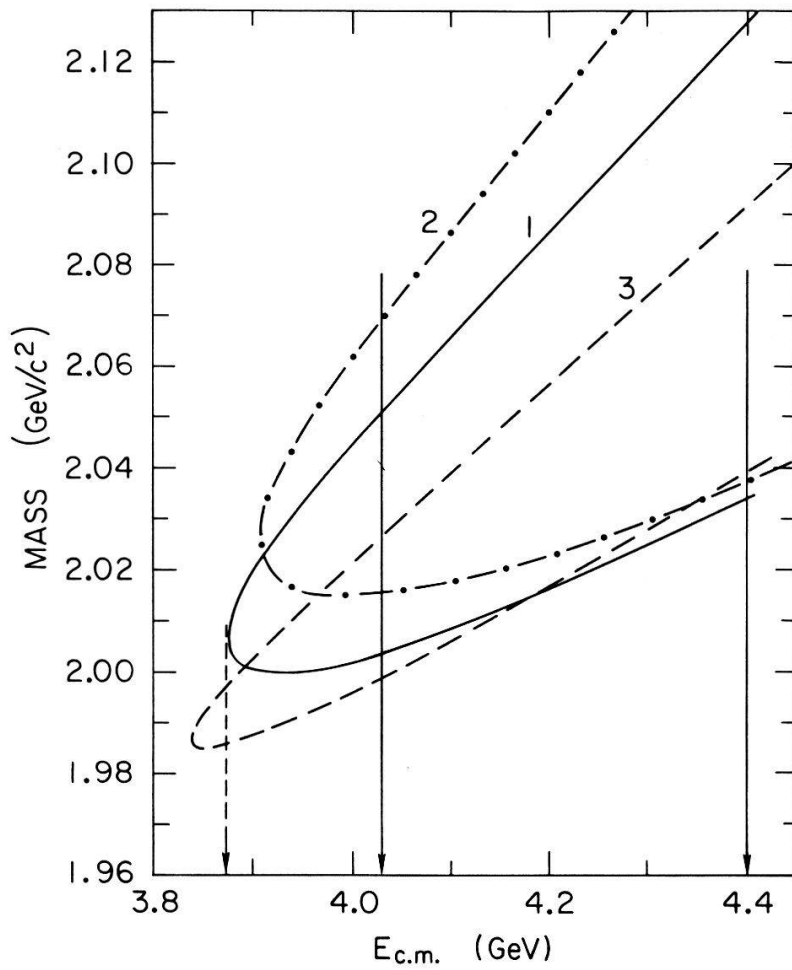


Figure 8
Kinematical upper and lower limits of the reflection bumps in the invariant mass spectrum as a function of the c.m.-energy and the D -meson masses.

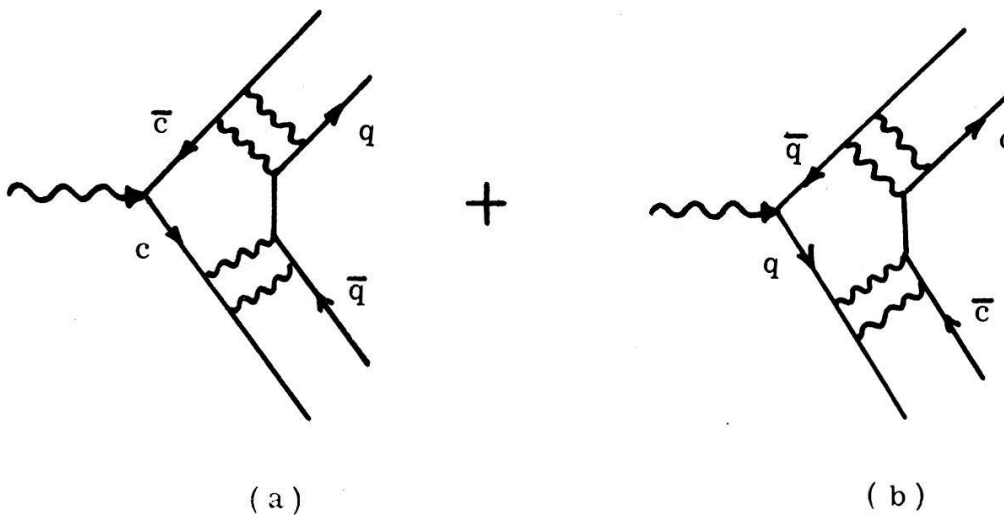


Figure 9
Quark diagrams for electromagnetic D -pair production. (a) Photon c -quark coupling and $q\bar{q}$ association. (b) Photon q -quark coupling and $c\bar{c}$ association.

applies. In order to find the structure of the cross sections we have calculated [18] the influence of both diagrams in Figure 9 in a free quark model assuming that the D -meson wave functions are linear combinations of spinors with the current operator

$$J^\mu = ie[Q_c \bar{\psi}_c \gamma^\mu \psi_c + Q_q \bar{\psi}_q \gamma^\mu \psi_q] \quad (4.7)$$

'sandwiched' in between. Q_c and Q_q are the charges of the charmed and lighter quarks in units of e . The resulting cross-section ratios then are

$$\begin{aligned} \sigma_{D\bar{D}} : \sigma_{D\bar{D}^* + \bar{D}D^*} : \sigma_{D^*\bar{D}^*} &= (Q_c F_c + Q_q F_q)^2 : 4 \left(Q_c F_c \frac{E_c}{m_c} - Q_q F_q \frac{E_q}{m_q} \right)^2 : \\ &: 3(Q_c F_c + Q_q F_q)^2 + 4 \left(Q_c F_c \frac{E_c}{m_c} + Q_q F_q \frac{E_q}{m_q} \right)^2 \end{aligned} \quad (4.8)$$

where kinematical factors due to phase space have identical forms in all three channels and have been ignored. The form factors F_c and F_q are momentum-dependent, thus $F_{c,q} \equiv F_{c,q}(p)$, and give the relative size of the contributions due to diagrams (a) and (b) in Figure 9. One notices that the DD -channel has an electric coupling whereas the form of $\sigma_{D\bar{D}^* + \bar{D}D^*}$ reflects a magnetic coupling and $\sigma_{D^*\bar{D}^*}$ has both. If $F_q \equiv 0$ we are back at the 1:4:7 ratio as was found earlier in Refs. [15] and [16]. The form of the cross sections (equation (4.8)) gives us the possibility to test experimentally whether $F_q \ll F_c$. The cross sections of the three channels are *measured at one and the same momentum value p which means at different energy points*:

$$\sigma_{D\bar{D}}(s_1), \sigma_{D\bar{D}^* + \bar{D}D^*}(s_2), \sigma_{D^*\bar{D}^*}(s_3) \quad (4.9)$$

We then form the combination

$$\sigma_{D^*\bar{D}^*} - 3\sigma_{D\bar{D}} - \sigma_{D\bar{D}^* + \bar{D}D^*} \propto F_c(p) \cdot F_q(p) \quad (4.10)$$

and thus we can experimentally test whether the above assumption is correct. This test, although quite general, is however only applicable in a region where no resonance is present. A similar technique can also be used to determine the ratio formed by the electric and magnetic form factors [18, 19, 20].

5. Experimental implications of charmed quarks

In this section we discuss other experimental tests which are being carried out at present in order to obtain more information on the nature of the charmed quark and we will point to missing pieces in the overall charm picture.

(A) Prompt leptons in e^+e^- annihilation

Experiments at DESY measure the inclusive production of $e^+e^- \rightarrow e^\pm X$ and $e^+e^- \rightarrow e^\pm KX$ at 3.6, 4.1, and 4.4 GeV with the results [21]:

- (i) Inclusive electron signals are found at 4.1 and 4.4 GeV, but not at 3.6 GeV, indicating a new particle production threshold.
- (ii) The multiplicity for these events peaks at $n \simeq 5-6$ at 4.1 GeV indicating that the source is not heavy lepton production, for which $n \simeq 2-2.5$ is predicted.
- (iii) An eK signal is found at 4.1 GeV; the signal is suppressed at 4.4 GeV. This suggests that D and F semileptonic decays are being observed at 4.1 GeV,

while charmed baryon production and decay may account for the inclusive e^\pm at 4.4 GeV.

- (iv) The electron energy distributions measured at 4.1 GeV peak at low E_e with no events observed for $E_e > 0.7$ GeV; this is quite compatible with charmed particle decays $D \rightarrow K e \bar{\nu}_e$, $K \pi e \bar{\nu}_e$, $K \pi \pi e \bar{\nu}_e$ (see Fig. 10).

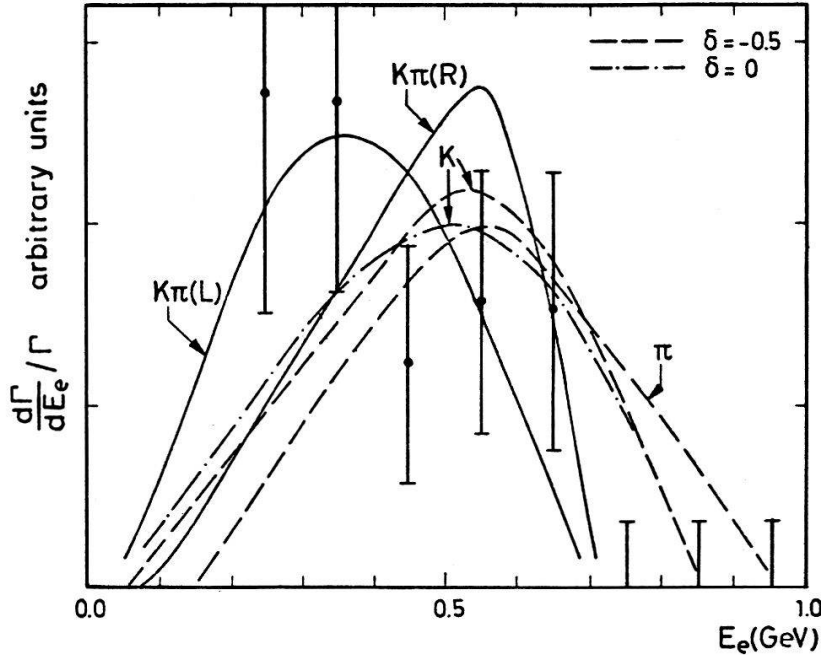


Figure 10

Energy spectrum of direct electrons coming from D -mesons at rest.

(B) Deep inelastic neutrino experiments

The dilepton events found in neutrino initiated deep inelastic experiments indicate the production and subsequent semileptonic decay of new particles which are compatible with being charmed mesons (see Fig. 11) [22].

(C) Photoproduction experiments

The energy dependence of the differential cross section of $\gamma N \rightarrow \psi + X$ near threshold shows an unusually flat onset but rises steeply from $E_\gamma \sim 12$ GeV (Fig. 12). This suggests that there is a pseudothreshold at this energy corresponding to the production of charmed particles. Prompt lepton measurements in this process show a significant excess of leptons at $E \sim 20$ GeV with a ratio $\mu/\pi = (1.4 \pm 2.5) \cdot 10^{-4}$ whereas no excess is observed at $E_\gamma = 8$ or 12 GeV [23]. This fact again is attributed to charmed mesons decaying semileptonically. Measurements in photon initiated reactions at Fermilab with energies $\langle E_\gamma \rangle \sim 120$ GeV have led to the discovery of charmed baryon states C_0 , C_1 , and C_1^* with masses near 2.26, 2.42, and 2.48 GeV [14].

The pseudoscalar state η_c , an important missing piece in the overall charm picture, will perhaps be discovered by the Primakoff production process (Fig. 13). The characteristic energy and Z -dependence of this strongly forward-peaked process open it to a very selective observation. η_c is expected to decay predominantly into two photons [24, 6].

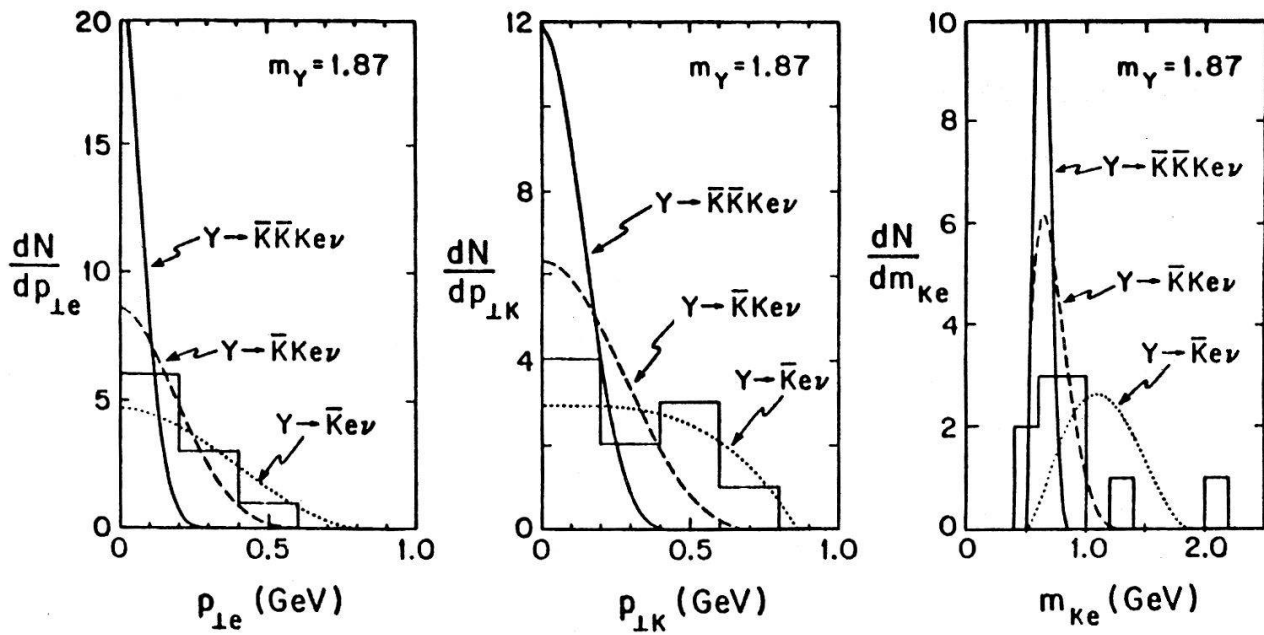


Figure 11
Momentum and invariant mass spectrum of the dilepton events observed in deep-inelastic neutrino reactions.

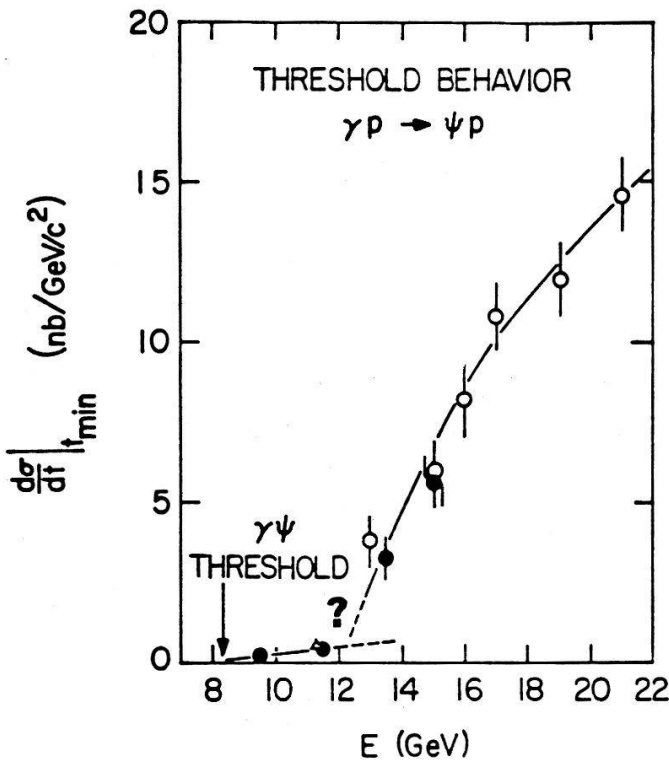


Figure 12
Threshold onset of the differential cross section for ψ -photoproduction.

(D) Hadron experiments

In purely hadronic processes, charmed particles have not yet been discovered but their production cross sections have been estimated by supposing a Drell-Yan type mechanism [25] (Fig. 14) [26]. One possible experiment is to look for dileptons coming from the charmed particle decays. Another type of experiment triggers on

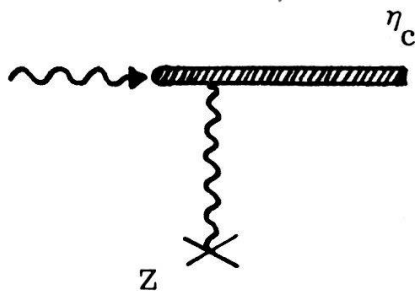


Figure 13
 η_c -production using the Primakoff effect.

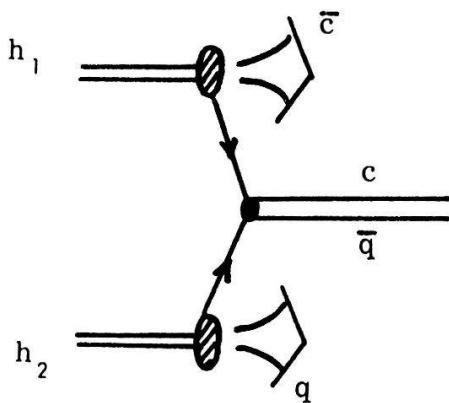


Figure 14
Drell-Yan production mechanism of charged particles in hadron initiated reactions.

leptons and looks for bumps in the reconstructed invariant mass of the final state hadrons which are expected to be mainly of strange type. A third type of experiment searches for the occurrence of sharp mass peaks associated with kaons, for two simultaneously occurring particle combinations [6]. Although not impossible, charmed particle detection in purely hadronic processes will be much more difficult than in weak and electro-magnetic production processes since the conventional hadronic processes are strongly competing.

What is missing in the charm picture? The state at 2.8 GeV in electron-positron annihilation, seen at DESY and conjectured as paracharmonium η_c does not fit too well in the overall picture; it is expected to be much closer to ψ . The $F \equiv (c\bar{s})$ mesons are still missing and a number of charmed baryons are waiting to be detected (Fig. 15) [27]. The anticipated higher $c\bar{c}$ resonances have to be experimentally confirmed as for instance the 3^3S_1 -resonance which is placed at 4.03 GeV by Eichten and Lane [16].

Since charmed particles now exist one wonders about their properties. Do they violate parity in their decays, or CP-invariance perhaps? We briefly indicate proposed tests:

(a) *P-violation*. Parity violation in the weak decays of the charmed D -mesons can be determined by the following three methods:

- (i) The simultaneous appearance of the decay modes $K^+\pi^-\pi^-$ and $K_s^0\pi^-$ or $K^+\pi^-$, $K^-\pi^+\pi^+\pi^-$ and $K_s^0\pi^+\pi^-$ indicates P -violation [28].
- (ii) If the $K^\pm\pi^\mp$ and $K^\pm\pi^\mp\pi^\mp$ peaks belong to the same isospin multiplet with

$J = 0$ then parity is violated. If instead $J^P(K\pi) = 0^+, 1^-, 2^+ \dots$ and $K\pi\pi$ does not vanish on the Dalitz plot boundary, then again parity is violated [29].

- (iii) Measurement of the momentum of two or three particles in an inclusive decay and formation of

$$\begin{aligned} A &\equiv \mathbf{p} \cdot (\mathbf{K}_1 \times \mathbf{K}_2) \cdot (\mathbf{K}_1 - \mathbf{K}_2) \cdot \mathbf{p} \\ B &\equiv \mathbf{K}_1 \cdot (\mathbf{K}_2 \times \mathbf{K}_3) \end{aligned} \tag{5.1}$$

indicates P -violation if the average value $\langle A \rangle$ or $\langle B \rangle$ is nonvanishing [30].

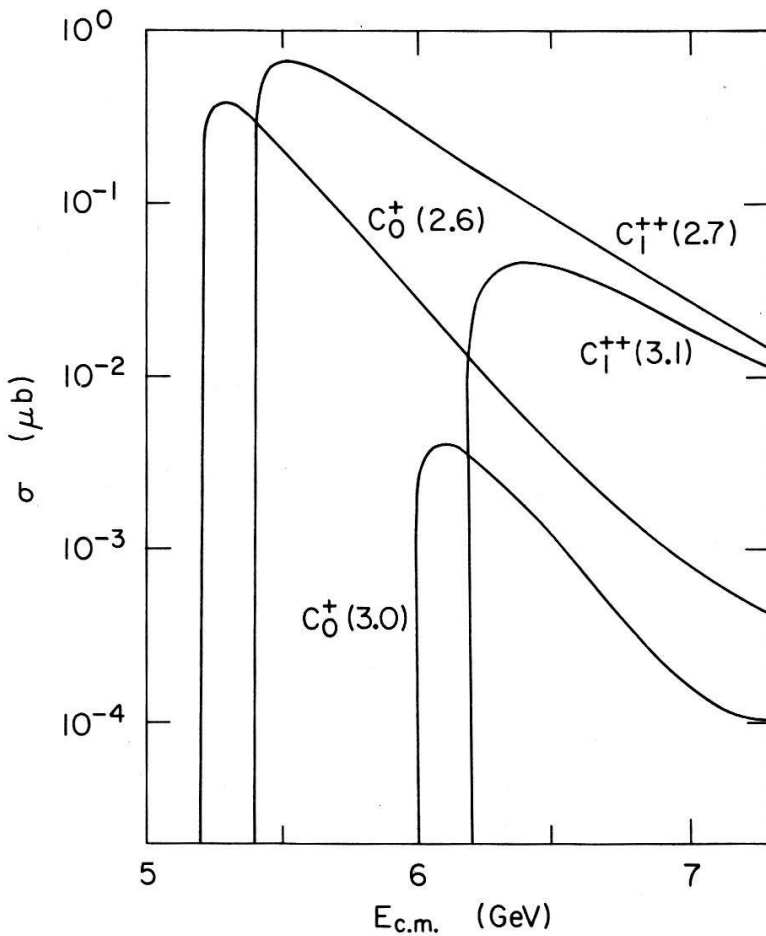


Figure 15

Threshold onset of charmed baryon-antibaryon pair produced in e^+e^- -annihilation.

(b) CP -violation [31]. CP -invariance could also be tested using for example the leptonic decay modes: $D^0 \rightarrow K^- \mu^+ \nu_\mu$. Let us denote by $N^{\pm\pm}$ decays like $(K^\pm \mu^\mp \bar{\nu}_\mu)$ and similarly N^{+-} for decays like $(K^+ \mu^- \bar{\nu}_\mu)(K^- \mu^+ \nu_\mu)$. CP -violation would give a charge asymmetry

$$\delta \equiv \frac{N^{++} - N^{--}}{N^{++} + N^{--}} \simeq 4 \operatorname{Re} \epsilon \tag{5.2}$$

where ϵ is the CP -violation parameter in the wave functions of the D_s and D_L mesons formed by linear combinations of D_1 and D_2 which are pure CP -eigen-states. Estimates, however, indicate that the effect will be almost unmeasurably small.

6. Charmed quarks—what next?

In the preceding sections we have limited ourselves to the immediate theoretical and experimental implications of the charmed quark discovery. In this section we would like to go beyond this limit by presenting the arising new questions and by pointing to missing pieces of the gauge theory approach to particle dynamics.

(A) How many quarks are there?

Since four elementary constituents (quarks) now seem to exist, the question naturally arises: why not more? The introduction of a fourth quark into the strong interaction scheme established a close resemblance between weak and strong interaction theory

$$\begin{pmatrix} \nu_e \\ e^- \end{pmatrix} \begin{pmatrix} \nu_\mu \\ \mu^- \end{pmatrix} \rightarrow \begin{pmatrix} u \\ d \end{pmatrix} \begin{pmatrix} c \\ s \end{pmatrix} \quad (6.1)$$

besides solving a number of deeper problems; for instance the cancellation of anomalies. Moreover the generalization to schemes with four quarks became evident. The extension of weak interaction theory to six quarks has been investigated by a number of authors [32] which could show that such a step, although not unique, does not lead to any obvious difficulties. Indeed, there are experimental indications and theoretical arguments which give hints at a six quark scheme:

- (i) The μe events at SPEAR indicate that charged heavy leptons L^\pm exist with decay modes: $L^- \rightarrow e^- + \bar{\nu}_e + \nu_L, \mu^- + \bar{\nu}_\mu + \nu_L$. If this interpretation is correct and if the leptons decay weakly via \bar{V} -A interaction, they are most likely grouped in three SU_2 doublets of left-handed leptons instead of two:

$$\begin{pmatrix} \nu_e \\ e^- \end{pmatrix} \begin{pmatrix} \nu_\mu \\ \mu \end{pmatrix} \begin{pmatrix} \nu_L \\ L \end{pmatrix} \rightarrow \begin{pmatrix} u \\ d \end{pmatrix} \begin{pmatrix} c \\ s \end{pmatrix} \begin{pmatrix} t \\ b \end{pmatrix} \quad (6.2)$$

which on the hadronic side, by analogy, makes it plausible to expect six quarks.

- (ii) In the four-quark scheme the asymptotic hadron to μ -pair production ratio in electron-positron initiated reactions is $R = 3\frac{1}{3}$ whereas the experimental value above $E_{c.m.} = 4.5$ GeV is between 5 and 5.5. This discrepancy cannot be explained without additional quarks.
- (iii) Deep inelastic neutrino processes also give indications for a six quark scheme. The ratio $\sigma(\nu N \rightarrow \mu^+ + X)/\sigma(\nu N \rightarrow \mu^- + X)$ seems to rise from 0.4 at low energies to 0.6–0.7 at higher energies whereas a four-quark model with the usual small component of $q\bar{q}$ -pairs in the large t nucleon, predicts a ratio 1/3.
- (iv) There exist further theoretical reasons [33].

It thus appears quite possible that, as we go to higher energies, more ψ -like $t\bar{t}$ or $b\bar{b}$ narrow resonances as well as $(b\bar{u}), (b\bar{d}), \dots (b\bar{c}), \dots (t\bar{c}), \dots (t\bar{b})$ bound states, will be found – a wealth of new particles!

(B) Where are the gluons?

What are the interacting forces between the quarks–gluons? If quantum chromodynamics is the underlying theory dominating the interaction between

quarks, as charmonium-type calculations motivated by asymptotic freedom arguments indicate, one should be able to find measurable consequences of gluons.

If gluons are confined, only indirect tests are possible. If they are unconfined, they might have escaped detection. In the latter case their masses (widths) have been estimated in the energy range 2–3 GeV (1–few MeV) and a number of measurable predictions such as decay channels, branching ratios, etc., exist [34]. In a theory with Han–Nambu quarks, the ratio R in e^+e^- annihilation is unaffected if the gluon mass is much bigger than the momenta under consideration. However if the gluon mass is small we have

$$R_{1/2} = 2 \left(1 + \frac{g'^2}{\bar{g}^2} \right)$$

$$R \equiv R_{1/2} + R_1 \quad \text{with}$$

$$R_1 = \frac{1}{4} \left(\frac{m_x}{M_G} \right)^4$$

where $m_x \geq 2.5$ GeV is the mass of a neutral field. M_G is the mass of the produced gluons. The structure function \bar{F}^2 describing $e^+e^- \rightarrow p + X$ is also proportional to $(m_x/M_G)^4$ whereas \bar{F}^1 is unaffected by the produced gluons [35].

Let us assume that gluons are confined and look for measurable consequences:

- (i) Heavy lepton pair production in hadron collisions, described by the Drell–Yan process, will show logarithmic deviations from scaling [36].
- (ii) Large- p_T hadrons in e^+e^- initiated reactions are expected to be mainly due to gluons [37].
- (iii) In deep-inelastic experiments color excitation which manifests itself by gluon terms is predicted to lead to a 15% change in the momentum conservation sum rule if color threshold is passed [38].
- (iv) Hadron multiplicities in hadron and lepton initiated reactions are the same. Furthermore, in the central region the jet structure and associated hadron multiplicities are the same in e^+e^- annihilation in deep inelastic scattering and in forward hadron collisions [39].
- (v) Spin measurements in ψ -photoproduction will reveal almost exact s -channel helicity conservation even in the threshold region if gluons are responsible for the diffractive characteristics of this process [40].

Gluons are a characteristic ingredient of gauge theories and therefore should be looked for in nature.

(C) What more?

Intermediate vector bosons are waiting to be discovered [41] and more leptons still might exist. Speculations ranging from new leptonic interactions up to leptonic quarks still exist [42]. This list of open questions could be continued for quite a while and shows that many of the fundamental problems still are unsolved even after the discovery of the charmed quark.

7. Conclusion

In this paper we have presented the consequences of the discovery of a charmed quark. We have presented the most recent experimental results on D -production in

e^+e^- annihilation and we have analysed the data by simple model calculations. The experimental search for more information on the characteristics of charmed particles has been indicated and some of the fundamental theoretical questions awaiting more information from the experiments were discussed.

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THEORETISCHE PHYSIK

Théories unifiées et groupes exceptionnels

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On présente une étude systématique de théories de jauge permettant d'unifier les interactions faibles, électromagnétiques et fortes, basées sur des groupes de Lie exceptionnels. Les conséquences phénoménologiques (couplage, structure des courants faibles, etc.) seront discutées.

Exact results for random planar conductance networks

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We consider planar (two-terminal) networks Z of conductances σ_i which are independently distributed according to a probability density $\rho(\sigma)$. The equivalent conductance $\bar{\sigma}$ of Z will then be distributed according to a probability density $R(\bar{\sigma})$ that is determined by the graph of Z and by ρ . We derive exact relations between the probability densities R and R^* of two dual random networks Z and Z^* , and discuss their application to some important special cases (self-dual networks, binary disorder, square lattice with anisotropic $\rho(\sigma)$, etc.). If, for infinitely extended regular networks, it is assumed that $\bar{\sigma}$ has a unique value, we can prove that the effective-medium theory satisfies all these exact relations.

Verallgemeinerung der adiabatischen Approximation der Werthamer Gleichung für einen punktförmigen Josephson Kontakt

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Unter der adiabatischen Näherung (AA_1) der Werthamer integro differentialgleichung WE versteht man diejenige Approximation, wo $\dot{\phi}$ (instantane Spannung $\dot{\phi}$)

in den Integraltermen Null gesetzt wird. AA_1 liefert eine Differentialgleichung die üblicherweise Josephsongleichung JE genannt wird. Es ist konsistent höhere Ableitungen $\ddot{\varphi}$, $\ddot{\ddot{\varphi}}$, \dots , $\varphi^{(N)}$ bei der Auswertung der WE-integrale linear mitzunehmen (AA_N), zumal dies für $N = 2$ eine phasenabhängige Korrektur der Kapazität ergibt. Erfahrungsgemäss liefert AA_1 gute Resultate nur dann, falls die Lösung eine langsam mit der Zeit veränderliche oder mit kleiner Amplitude um Null oszillierende Spannung hat. Dies trifft in Parameterbereichen zu, wo die WE und alle AA_N analytisch durch asymptotische Entwicklungen gelöst werden können. Damit kann gezeigt werden, dass die Charakteristiken von WE und AA_N nicht einmal für $N \rightarrow \infty$ in der niedersten nichttrivialen Ordnung übereinstimmen. Die Ursache der Diskrepanz wird genauer diskutiert.

Atomic-number dependent model potential for first row atoms

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A new atomic model potential based on the exact pseudopotential theory and a realistic description of core electrons is proposed. The potential is explicitly given for all atoms of the first row and has several interesting properties: (i) it scales with atomic number, (ii) it depends on a few parameters only, and (iii) it has a simple analytic expression. This potential, contrary to previous models, reproduces with high accuracy not only electronic energy levels but also the corresponding wave functions. It is therefore suitable for accurate atomic, molecular, and solid state calculations.

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Induzierte Röntgenübergänge in Schwerionenkollisionen

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Wir zeigen, dass die Bewegung der Kernladungen Strahlungsübergänge der Elektronen induziert. Das kommt daher, dass das Vektorpotential als Summe $A^{kl} + A^q$ auftritt, wobei A^{kl} das klassische Vektorpotential der kollidierenden Kerne ist und A^q das übliche, quantisierte Vektorpotential bedeutet. In der nicht relativistischen Formulierung wird die Strahlung durch den Operator $(e^2/mc^2)A^{kl}A^q$ verursacht, der im A^2 Term enthalten ist. Die induzierte Strahlung ist anisotrop. Sie ist dann relevant, wenn $(e/c)A^{kl}$ von der Grössenordnung des Elektronenimpulses \mathbf{p} ist. Dies ist der Fall für schwerere Kerne, hohe Projektilgeschwindigkeiten und für Übergänge in den äusseren (M, N) Schalen, wo \mathbf{p} klein ist. Der Effekt könnte somit die beobachteten Anisotropien in den molekularen M -Röntgenspektren bei Schwerionenkollisionen erklären.

Order-parameter relaxation and Josephson effects in short superconducting weak links

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The relaxation time τ for the order parameter ψ is typically longer than the characteristic Josephson time $\tau_j = \hbar/2eV_j$, where $V_j = R_N I_m$ for a dirty weak link of length $a < \xi$ (coherence length) and maximum super-current $I_m \propto \xi/a$. The resulting consequences were studied by Likharev and Yakobson [1] within time-dependent Ginsburg-Landau theory. Our calculations on the same model confirm their results for small $\gamma = (\tau/\tau_j)(a/2\xi)^2$. In the opposite limit their *Ansatz* for $\psi(x, t)$ is not valid for $V \gtrsim V_j$, and hysteresis appears, i.e., $I_{\min}(V \neq 0) < I_m$ for $\gamma \gtrsim 10$. This effect and the reduction in the amplitude of voltage oscillations with increasing γ are similar to the effect of a parallel capacitance in the standard shunted junction model which is recovered when $\gamma = 0$.

Reference

- [1] K. K. LIKHAREV and L. A. YAKOBSON, Sov. Phys. JETP 41, 570 (1975).

Order parameter dynamics and current hysteresis in long superconducting filaments

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By proper adjustment of the boundary conditions used in our previously described investigation of weak links we can simulate the dynamics of long superconducting filaments at finite voltage. Computed I-V characteristics exhibit a plateau at a value $J_L < J_{m1}$ (max. supercurrent). Above J_L the filament is normal except near its ends. Below J_L the filament is mostly superconducting, but order parameter and voltage oscillations persist near its center down to J_{\min} . Both J_L and J_{\min} decrease as τ/τ_j increases.

Reference

- [1] K. K. LIKHAREV, Soviet Phys. JETP Letters 20, 338 (1974).

Dynamische Korrelationsfunktionen für lineare Spinketten

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Die dynamischen Spin-Korrelationsfunktionen für zyklische lineare Ketten, welche bis zu zehn Spins mit $s = \frac{1}{2}$ und $s = 1$ umfassen, werden durch explizite Diagonalisierung des Hamiltonoperators numerisch berechnet. Wir behandeln ferro- und antiferromagnetische Systeme und berücksichtigen auch anisotrope Austausch- und Kristallfeldterme. Die Resultate sind in guter Uebereinstimmung mit Neutronstreuquerschnitten von quasi-eindimensionalen Systemen und bestätigen gewisse analytische Voraussagen.

Korrelationen und lineare Antwort für stochastische Prozesse mit Gedächtnis (Nicht-Markoff-Prozesse)

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In der Theorie des stochastischen Verhaltens offener Systeme führt die Elimination der Badvariablen bekanntlich auf die Beschreibung durch einen stochastischen Prozess mit Gedächtnis (Nicht-Markoff-Prozess) [1]. Wir leiten für einen solchen Prozess einen Satz von kinetischen Gleichungen für die n -fachen Verbundwahrscheinlichkeiten her und diskutieren die Berechnung von Mittelwerten und Korrelationen. Ferner entwickeln wir die Theorie der linearen Antwort. Im Gegensatz zum Markoff-Fall lässt sich die Suszeptibilität nicht durch eine stationäre Korrelationsfunktion ausdrücken.

Literatur

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Size dependence of the lifetime of metastable states in the kinetic one-spin-flip Ising model¹⁾

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Using the Monte Carlo technique we have studied the relaxation of the order parameter after a sudden reversal of the applied field as a function of the size of the system. We considered square Ising models with nearest-neighbor interaction subjected to periodic boundary conditions. For large systems, the lifetime of metastable states is size independent. For small systems, however, the reciprocal relaxation time is proportional to the probability for the occurrence of a cluster covering the half system. In an intermediate range, this inverse relaxation time grows with the number of spins.

¹⁾ A more detailed description will appear in *Physica B* (Neth.) in February 1977.

Adsorption an einer porösen Oberfläche und die Theorie des Porosimeters

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Die Adsorption an einer porösen Oberfläche wird im Rahmen der Thermodynamik beschrieben. Eine Pore sei charakterisiert durch ihr Volumen V , die Fläche der Porenwand A_w und die Fläche der Oeffnung A_d , ein poröser Körper durch die Verteilungsfunktion $n(V, A_w, A_d)$. Wir berechnen das Gibbs'sche Potential der adsorbierten Flüssigkeit und zeigen, dass man aus Adsorptionsmessungen auf die Verteilungsfunktion n schliessen kann.

Beim Porosimeter wird eine nicht benetzende Flüssigkeit (z.B. Hg) unter Druck in die Poren gepresst. Registriert wird das Volumen der hineingepressten Substanz in Funktion des Druckes. Wir entwickeln die Thermodynamik des Porosimeters analog zur Thermodynamik der Adsorption und zeigen, dass man in beiden Fällen im wesentlichen das Gleiche misst. Die übliche Theorie des Porosimeters folgt als Spezialfall für tiefe, zylindrische Poren.

