[1]R. Van Dyck, P. Schwinberg, and H. Dehmelt, Phys. Rev. Lett. 38, 310 (1977).
[2]J. Wesley and A. Rich, Phys. Rev. A 4, 1341 (1971).
[3]T. S. Jaseja, A. Javan, J. Murray, and C. H. Townes, Phys. Rev. 133, A1221 (1964).
[4]R. F. C. Vessot and M. W. Levine, Metrologia 6, 116 (1970).
[5]R. V. Pound and G. A. Rebka, Jr., Phys. Rev. Lett. 4, 274 (1960).
[6]H. Hay, J. Schiffer, T. Cranshaw, and P. Egelstaff, Phys. Rev. Lett. 4, 165 (1960).
[7]J. C. Hafele and Richard E. Keating, Science 177, 166 (1972).
[8]H. I. Mandelberg and L. Witten, J. Opt. Soc. Am. 52, 529 (1962).
[9]K. Brecher, Phys. Rev. Lett. 39, 1051, 1236(E) (1977).
[10]D. J. Grove and J. G. Fox, Phys. Rev. 90, 378 (1953).
[11]V. P. Zrelov, A. A. Tyapkin, and P. S. Farago, Zh. Eksp. Teor. Fiz. 34, 555 (1958) [Sov. Phys. JETP 7, 384 (1958)].
[12]D. S. Ayres, A. M. Cormack, A. J. Greenberg, R. W. Kenney, D. O. Cladwell, V. B. Elings, W. P. Hesse, and R. J. Morrison, Phys. Rev. D 3, 1051 (1971).
[13]T. Alvager, F. J. M. Farley, J. Kjellman, and I. Wallin, Phys. Lett. 12, 260 (1964).
[14]Z. G. T. Guiragossián, G. B. Rothbart, M. R. Yearian, R. A. Gearhart, and J. J. Murray, Phys. Rev. Lett. 34, 355 (1975).
[15]J. Bailey, K. Borer, K. Combley, H. Drumm, F. Krienen, F. Lange, E. Picasso, W. von Rüden, F. J. Farley, J. H. Field, and P. M. Hattersley, Nature 268, 301 (1977).
[16]M. P. Balandin, V. M. Crebenyuk, V. G. Zindv, A. D. Konin, and A. N. Pondmarev, Zh. Eksp. Teor. Fiz. 67, 1631 (1974) [Sov. Phys. JETP 40, 811 (1974)].
[17]C. Møller, *The Theory of Relativity* (Oxford Univ. Press, London, 1952).
[18]V. Bargmann, L. Michel, and V. L. Telegdi, Phys. Rev. Lett. 2, 435 (1959).
[19]G. W. Ford and C. W. Hirt, The University of Michigan Report, Contract No. Nonr 1224(15), 1961 (unpublished).
[20]S. Granger and G. W. Ford, Phys. Rev. D 13, 1897 (1976).
[21]J. Bailey, K. Borer, F. Combley, H. Drumm, F. Farley, J. Field, W. Flegel, P. Hattersley, F. Krienen, F. Lange, E. Picasso, and W. von Rüden, Phys. Lett. 68B, 191 (1977).
[22]J. Bailey, W. Bartl, B. vonBochmann, R. Brown, F. Farley, M. Giesch, H. Jöstlein, S. van der Meer, E. Picasso, and R. Williams, Nuovo Cimento A9, 369 (1972).
[23]J. E. Romain, Rev. Mod. Phys. 35, 376 (1963).
[24]C. W. Sherwin, Phys. Rev. 120, 17 (1960).

Instantons and the Hypothetical Light Boson

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The mass formula for the π, η, A (axion) system is analyzed taking into account explicit instanton effects. It is argued that small-size instantons serve as an independent source for the axion's mass which may render the axion more massive than the mixing of the bare axion with π and η.

There is a revived hope that the reality of Higgs particles can be finally established. Very recently Weinberg[1] and Wilczek[2] independently observed that the existence of an elementary pseudo Goldstone boson is an indispensable feature of a certain class of *CP-conserving* models which combine colored gauge interactions with unified electromagnetic and weak interactions. These models were defined earlier by Peccei and Quinn[3] by requiring an extended axial $U_{PQ}(1)$ symmetry for the entire Lagrangian. The symmetry is broken spontaneously by Higgs fields and explicitly by Adler-Bell-Jackiw anomalies. Hence the corresponding current A_μ^{PQ} is a source of pseudo Goldstone bosons—axions. In this note explicit effects of instantantons on properties of the axion will be exhibited.

The Lagrangian of the theory is given by

$$\mathcal{L} = \mathcal{L}_{G,W} + \mathcal{L}_{Y} + \mathcal{L}_{H,L}, \tag{1}$$

where $\mathcal{L}_{G,W}$ describes the gauge interaction of N-flavor quarks through colored gluons (G) and weak bosons (W) according to standard schemes SU(3) and $U_R(1) \otimes SU_L(2)$, respectively; $\mathcal{L}_{H,L}$ describes weak interactions of both leptons and *two* Higgs doublets $\varphi_i = (\varphi_i^+, \varphi_i^0)$, $i = 1, 2$, according to the scheme $U_R(1)\otimes SU_L(2)$, and the Higgs potential in $\mathcal{L}_{H,L}$ is assumed to be symmetric with respect to global phase transformations performed on φ_1 and φ_2 independently; $\mathcal{L}_Y$ represents the $U_R(1)\otimes SU_L(2)$-invariant Yukawa couplings of quarks and leptons with Higgs fields. The part of $\mathcal{L}_Y$ which contains neutral Higgs components φ_i^0 can be written in the form[4]

$$\mathcal{L}_{YN}(u,d,s,c,\ldots) = \{(m_d\bar{d}_R d_L + m_s\bar{s}_R s_L + \ldots)(\varphi_1^{0*}/\lambda_1) + (m_u\bar{u}_R u_L + m_c\bar{c}_R c_L + \ldots)(\varphi_2^0/\lambda_2) + \text{c.c.}\}. \quad (2)$$

Here parameters $\lambda_i = \langle\varphi_i^0\rangle$, $i=1,2$, fix vacuum expectation values of the φ_i's determining the spontaneous breaking of the weak gauge symmetry as well as the aforementioned global phase transformations on the φ_i's.

These transformations may be generated by the following currents:

$$A_\mu^1 = \tfrac{1}{2}\bar{\psi}\gamma_\mu\gamma_5\tau_3\psi + i\varphi_1^{0*}\overleftrightarrow{\partial}_\mu\varphi_1^0 + i\varphi_2^{0*}\overleftrightarrow{\partial}_\mu\varphi_2^0 + \ldots, \quad (3)$$

$$A_\mu^{PQ} = \tfrac{1}{2}\bar{\psi}\gamma_\mu\gamma_5\psi - i\varphi_1^\dagger\overleftrightarrow{\partial}_\mu\varphi_1 + i\varphi_2^\dagger\overleftrightarrow{\partial}_\mu\varphi_2 + \ldots, \quad (4)$$

where contributions from the charged vector bosons and leptons have not been exhibited. The ψ represents N-flavor quarks. Evidently the above currents must be sources of two would-be Goldstone bosons. One of these bosons renders a weak neutral vector boson massive, whereas the other, the axion, acquires a mass through the anomaly responsible for the nonconservation of the corresponding current.[4]

The source of the axion's mass can be identified further. Indeed, following Weinberg[1] and Bardeen and Tye[5] we observe that in the absence of weak-electromagnetic interactions and the mass term $\delta\mathcal{L} \equiv \mathcal{L}_{YN}(u,d,s)$ of u, d, s quarks, one can construct a conserved counterpart of A_μ^{PQ} and, in addition, two Goldstone currents with respective quantum numbers, π^0 and η, as follows:

$$A_\mu^2 = \tfrac{1}{2}\bar{\psi}\gamma_\mu\gamma_5\psi - \tfrac{1}{3}N(\bar{u}\gamma_\mu\gamma_5 u + \bar{d}\gamma_\mu\gamma_5 d + \bar{s}\gamma_\mu\gamma_5 s) - i\varphi_1^\dagger\overleftrightarrow{\partial}_\mu\varphi_1 + i\varphi_2^\dagger\overleftrightarrow{\partial}_\mu\varphi_2, \quad (5)$$

$$A_\mu^3 = \tfrac{1}{2}(\bar{u}\gamma_\mu\gamma_5 u - \bar{d}\gamma_\mu\gamma_5 d), \quad (6)$$

$$A_\mu^4 = \tfrac{1}{6}\sqrt{3}(\bar{u}\gamma_\mu\gamma_5 u + \bar{d}\gamma_\mu\gamma_5 d - 2\bar{s}\gamma_\mu\gamma_5 s). \quad (7)$$

Thus we have arrived at currents $\{A_\mu^a(x),\ a=1,2,3,4\}$ which represent a complete set of the neutral Goldstone bosons. It is assumed that the ninth pseudoscalar boson η' is massive even for $\delta\mathcal{L}=0$.

Henceforth I will ignore weak-electromagnetic interactions with respect to $\delta\mathcal{L}$. Therefore the $\delta\mathcal{L}$ can be identified as being responsible for both an explicit breaking of the above currents $A_\mu^{a\neq1}$ and a rendering of three out of four would-be Goldstone bosons massive.

Now calculations of the axion's mass will be sketched. The method is based on a standard current-algebra technique. Consider the correlation function of current divergences

$$\Delta_{ab}(k^2) = i\int d^4x\, e^{ikx}\langle 0|T[\partial^\mu A_\mu^a(x)\partial^\nu A_\nu^b(0)]|0\rangle. \quad (8)$$

In terms of the axial charges

$$Q_a^5 = \int d^3x\, A_0^a(x), \quad (9)$$

on has

$$\partial^\mu A_\mu^a = i[Q_a^5, \delta\mathcal{L}] \equiv iP^a(x), \quad (10)$$

$$\Delta_{ab}(0) = \langle 0|[Q_a^5,[Q_b^5,\delta\mathcal{L}]]|0\rangle, \quad (11)$$

which may be easily calculated. Notice that $P^a(x)$, $a=1,\ldots,4$, are given by various linear combinations of quark pseudoscalar densities, $\bar{q}\gamma_5 q$, as quantum fluctuations of Higgs fields are ignored $(\varphi_{1,2}=\lambda_{1,2})$.

The correlation function $\Delta_{ab}(s)$ is analytic in the complex s plane cut along the positive real semiaxis $L=[s_{th},+\infty]$ with $s_{th}=(3m_\pi)^2$. Hence it obeys the finite-energy sum rule

$$\Delta_{ab}(0) = F_{ac}M_{cd}^2F_{bd} + \bar{\Delta}_{ab}(0) \quad (12)$$

with

$$\bar{\Delta}_{ab}(t) = \frac{1}{2\pi i}\oint_C \frac{\Delta_{ab}(s)ds}{s-t}. \quad (13)$$

In Eq. (12) the contribution of one-particle intermediate states has been made explicit by the first term on the right-hand side where the M^2 is the mass matrix of pseudo Goldstone bosons and the matrix F has elements $F_{11}=-F_{21}=\lambda_1$, $F_{12}=F_{22}=\lambda_2$, $F_{13}=F_{33}=F_{44}=\sqrt{3}F_{14}=f_\pi\approx\frac{1}{2}\sqrt{2}m_\pi$

with remaining ones being zero. The contour C in Eq. (13) runs along the circle $|s|=s_0>s_{th}$ and lower and upper lips of the cut L. [I avoid writing a dispersion relation for $\Delta_{ab}(s)$ since at least one subtraction is needed. Clearly, the integral (13) is not necessarily positive.)

Equation (12) without the integral term reproduces previously reported results on the axion's mass and well known expressions for m_π and m_η.[1,5] I would like to emphasize the following points. Firstly, the derivation of Eq. (12) assumes that the pseudo Goldstone states can be well approximated by linear combinations of their bare counterparts $\mathrm{Im}\varphi_{1,2},\pi,\eta$. Secondly, it is a usual practice to ignore terms like $\overline{\Delta}_{ab}(0)$ as a higher-order quantity in the chiral-symmetry-breaking $\delta\mathcal{L}$. Now I will argue that instantons, without affecting the first approximation, render the second one inadequate.[6]

I begin by observing that terms FM^2F and $\overline{\Delta}(0)$ in Eq. (12), respectively, determine long- ($x^2 \gtrsim d_c^2 \sim s_{th}^{-1}$) and short- ($x^2 \lesssim d_c^2 \sim s_{th}^{-1}$) distance contributions to the configuration-space representation (8) for $\Delta(k^2=0)$. It should be further noted that the threshold enhancements are absent in the integral (13).[7]

I will proceed from the premise that $\overline{\Delta}(0)$ describes a short-distance propagation of a light quark-antiquark pair which is amenable to perturbation analysis. In particular, I will assume that the spontaneously broken character of the chiral symmetry can be ignored in evaluating $\overline{\Delta}(0)$. Thus $\overline{\Delta}(0)$ receives contributions from two different components which arise from perturbative expansions about the trivial and instanton solutions of the classical gauge-field equations.

The naive perturbative (parton) component is not expected to give a significant contribution. Otherwise, since it does not distinguish different densities $P^a(x)$, mass formulas for π and η would be at once invalidated. Furthermore, a direct estimate leads to $\overline{\Delta}(s)\sim m_{u,d,s}^2 s\ln(-s/s_{th})$ also suggesting $\overline{\Delta}(0)\ll\Delta(0)$. Thus we are left with the instanton component. It will be evaluated through the Euclidean configuration-space representation for $\overline{\Delta}(0)$ [cf. Eqs. (8) and (13)],

$$\overline{\Delta}_{ab}(0) = -\int_{x^2<d_c^2} d_E^4x \langle P^a(x)P^b(0)\rangle_E, \tag{14}$$

where $d_c^2\sim s_{th}^{-1}$ determines the infrared cutoff and the sign $\langle\cdots\rangle_E$ has been used for the Euclidean functional average.

We recall that an instanton with a size ρ describes a tunneling between topologically distinct vacua separated by a potential barrier of size ρ.[8-12] Apparently a quark-antiquark pair cannot propagate distances $\lesssim d_c$ traversing a barrier of size $\rho > d_c$. Therefore the parameter d_c in Eq. (15) seems to provide a natural infrared cutoff $\rho_c\sim d_c$ on the size of instantons which effectively couple to the light quark-antiquark pair.

After this observation the evaluation of Eq. (14), in principle, can be performed exactly. However, some approximation needs to be made, since solutions to the *massive* Dirac equation with an instanton potential are not available. I will use 't Hooft's tunneling amplitude,[9] generalized to the color SU(3) gauge group. The amplitude is sufficiently accurate for light quarks, $m_q\rho_c\ll 1$. A contribution from heavy quarks, such as charm c, etc., will be ignored since they effectively decouple from instantons ($m_c\rho_c\gg 1$).

Unfortunately, the above approximations seem to be less applicable to strange quarks $m_s\rho_c\lesssim 1$. For the purposes of orientation, two different evaluations, which are obtained under two extreme assumptions, $m_s\rho_c\ll 1$ and $m_s\rho_c\gg 1$, will be presented. In the latter case the perturbation $\delta\mathcal{L}$ is taken to be $\mathcal{L}_{YN}(u,d)$ rather than $\mathcal{L}_{YN}(u,d,s)$ [see Eq. (2).]

Explicit calculations show that all matrix elements $\overline{\Delta}_{ab}(0)$ *except* $\overline{\Delta}_{22}(0)$ identically vanish. This is not an unexpected result, since the instanton-generated effective Lagrangian induced by chirally asymmetric sources, such as $\delta\mathcal{L}$, is SU(3)⊗SU(3) symmetric. By itself the above result ensures that mass relations for π and η remain intact in their ordinary form. Furthermore the mixing of the axion with π and η is unaffected by the $\overline{\Delta}_{22}(0)$ term in Eq. (12). This can be verified directly. However the $\overline{\Delta}_{22}(0)$ term introduces a significant shift in the axion mass m_A:

$$m_A^2 = \left[\frac{\sqrt{2}G_F m_\pi^2 f_\pi^2}{\sin^2 2\alpha}\right] N^2 \left\{\frac{m_u m_d}{(m_u+m_d)^2} + \frac{1}{(m_\pi f_\pi)^2}\mathfrak{D}_{a(b)}(\rho_c)\right\}. \tag{15}$$

Here $\tan\alpha=\lambda_1/\lambda_2$, and $\mathfrak{D}_{a(b)}(\rho_c)$ determines the density of instantons of size $\rho<\rho_c$:

$$\mathfrak{D}_a(\rho_c) = \int_0^{\rho_c} (m_u m_d \rho^2) D(\rho\,|\,N_f=2)d\rho/\rho, \quad m_s\rho_c \gg 1, \tag{16a}$$

$$\mathfrak{D}_b(\rho_c) = \int_0^{\rho_c} (m_u m_d m_s \rho^3) D(\rho\,|\,N_f=3)d\rho/\rho, \quad m_s\rho_c \ll 1, \tag{16b}$$

TABLE I. The dependence of the axion's mass m_A on the infrared cutoff ρ_c. The parameter $x_{a(b)}$ is defined in Eq. (19).

$[(300\text{ MeV})\rho_c]^{-1}$	$m_A^{(a)}\sin 2\alpha/N$ (MeV)	$m_A^{(b)}\sin 2\alpha/N$ (MeV)	x_a	x_b
1.0	82	50	9.7	3.0
1.5	80	49	9.3	2.8
2.0	71	43	7.1	2.0
2.5	59	37	4.5	1.2
3.0	47	33	2.5	0.7

with

$$D(\rho|N_f)=2\left(\frac{2}{\pi^2}\right)\left(\frac{1}{\rho_4}\right)\left\{\frac{8\pi^2}{g^2(\rho)}\right\}^6\exp\left\{-\frac{8\pi^2}{g^2(\rho)}-\sum_{t=\frac{1}{2},1}N(t)\alpha(t)+0.29N_f\right\}, \tag{17}$$

where $\alpha(\frac{1}{2})=0.146$, $\alpha(1)=0.443$, and the number of isospinor and isovector gluon multiplet is given by $N(\frac{1}{2})=2$ and $N(1)=3$, respectively.[13]

In what follows all estimates will be carried out using the simple formula $8\pi^2/g^2(\rho)=-(11-\frac{2}{3}N_f)\ln(\rho\mu)$ with $\mu=300$ MeV corresponding to $\alpha_s(\rho)\equiv g^2(\rho)/4\pi=0.3$ at $\rho^{-1}=3$ GeV.[14] The infrared cutoff will be allowed to vary in the range $0.3<\rho_c\mu<1$. Further, quark masses will be taken to be $m_d/m_u=1.8$, $m_s/m_d=20.1$, and $m_s\approx 150$ MeV.[15]

Equation (15) assumes the validity of the dilute-gas approximation for the instanton distribution,[16] i.e.,

$$\epsilon\equiv\int_0^{\rho_c}\frac{\pi^2}{2}\rho^4\frac{d\mathfrak{D}_{a(b)}(\rho)}{d\rho}d\rho/\rho\ll 1, \tag{18}$$

which is well satisfied ($\epsilon<0.02$) because of the smallness of quark masses, $m_{u,d}\rho_c\ll 1$.

Table I shows the estimates for m_A^2 and the ratio x of the second term to the first term in Eq. (15),

$$x_{a(b)}=\frac{(m_u+m_d)^2}{m_um_d}\frac{1}{(m_\pi f_\pi)^2}\mathfrak{D}_{a(b)}(\rho_c). \tag{19}$$

Accepting $\rho_c^{-1}\sim 500$ MeV as a reasonable infrared cutoff one concludes that the axion's mass may be a few times larger than the mass component $m_A(x_{a(b)}=0)$ generated by the mixing of the bare axion with π and η.

I believe that the presented method of analysis is not restricted to the axion problem and may have other applications.

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[1]S. Weinberg, Phys. Rev. Lett. 40, 223 (1978).

[2]F. Wilczek, Phys. Rev. Lett. 40, 279 (1978).

[3]R. D. Peccei and H. R. Quinn, Phys. Rev. Lett. 38, 1440 (1971), and Phys. Rev. D 16, 1791 (1977).

[4]S. Weinberg, Phys. Rev. Lett. 17, 657 (1976).

[5]W. A. Bardeen and S.-H. H. Tye, "Current Algebra Applied to the Properties of the Light Higgs Boson" (to be published).

[6]For an analogous situation see V. Baluni and D. J. Broadhurst, Nucl. Phys. B84, 178 (1975), and references therein. Here the dispersion integral for the correlation function of strangeness-changing vector-current divergences was shown to be an important source of physical information.

[7]Propagators of pseudoscalar densities do not develop nonanalytic dependence on the chiral-symmetry–breaking parameters which usually arises from a threshold region; see H. Pagels and A. Zepeda, Phys. Rev. D 5, 3262 (1972).

[8]G. 't Hooft, Phys. Rev. Lett. 37, 8 (1976).

[9]G. 't Hooft, Phys. Rev. D 14, 3432 (1976).

[10]R. Jackiw and C. Rebbi, Phys. Rev. Lett. 37, 172 (1976).

[11]C. G. Callan, R. F. Dashen, and D. J. Gross, Phys. Lett. 63B, 334 (1976).

[12]C. G. Callan, R. F. Dashen, and D. J. Gross, Phys. Rev. D (to be published).

[13]In deriving Eq. (17) I take into account that there are two isodoublet gluon zero modes in addition to those for the SU(2) case and that the instanton occupies the phase-space volume $SU(3)/U(1)\otimes Z_2$ in the manifold of the SU(3) group. The resulting coefficient $(2/\pi^2)$ is somewhat different from that quoted in Ref. 12. The author wishes to thank Professor C. G. Callan, Professor R. F. Dashen, Professor D. J. Gross, Professor

J. Goldstone, and Dr. A. Guth for very beneficial discussions concerning Eq. (17).

[14]The charge $\alpha_s(\rho)$ is defined by the Pauli-Villars regularization. Its relation to the dimensionally regularized charge $\alpha_D(\rho)$ is $(\alpha_s/2\pi) \approx (\alpha_D/2\pi) + 9.5(\alpha_D/2\pi)^2$. Evidently α_D is not a suitable expansion parameter. This was recognized earlier; see V. Baluni, Phys. Rev. D 17, 2092 (1978).

[15]See, e.g., S. Weinberg, in *A Festschrift for I. I. Rabi*, edited by Lloyd Motz (New York Academy of Sciences, New York, 1977).

[16]This easily follows from the fact that instantons reduce the vacuum "energy density" by $\mathfrak{D}_{a(b)}(\rho_c)$ [cf. Eq. (15)]. Observe that in distinction to Ref. 12 the condition (18) is applied to the gas of instantons which are liberated by nonzero quark masses appearing in the Lagrangian (2) rather than constituent quark masses generated dynamically. The dynamically generated component of instantons was assumed to have a negligible effect for $\overline{\Delta}(0)$.

Discrete Energy Transfer in Collisions of Xe(nf) Rydberg Atoms with NH_3 Molecules

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In collisions with NH_3 molecules it is observed that Xe(nf) Rydberg atoms are selectively excited to discrete, more highly excited states. These states are identified by the technique of field ionization and are found to be displaced in energy from the initial Rydberg states by amounts equal to the rotational spacings of the NH_3 molecule. These measurements demonstrate, for the first time, the conversion of rotational energy to electronic energy in collisions.

Flannery[1] and Matsuzawa[2] have predicted that in collisions between Rydberg atoms and polar molecules rotational de-excitation of the molecules could provide the energy necessary to further excite or ionize the Rydberg atoms. To investigate this prediction we have made an experimental study of collisions between xenon Rydberg atoms Xe(nf) with $26 \leq n \leq 40$ and ammonia molecules.

The rotational term energies for NH_3 are given by

$$E_J = BJ(J+1) + (A-B)K^2, \tag{1}$$

where J and K are rotational quantum numbers and A and B are constants. In dipole-allowed de-excitations ($J \to J-1$, $K \to K$) the energies released, ΔE_J, are approximately

$$\Delta E_J = E_J - E_{J-1} = 2BJ. \tag{2}$$

If these energies are transferred to Xe(nf) Rydberg atoms, with term values T_n, the resulting states will have term values given approximately by

$$T_n' = T_n + 2BJ, \quad J = 1,\ 2,\ 3, \ldots. \tag{3}$$

Clearly if $2BJ$ exceeds $|T_n|$, T_n' will be positive and ionization will occur, while for $2BJ$ less than $|T_n|$ further excitation will result.

For the particular case of Xe(27f) atoms the possible reactions are

$$\mathrm{Xe}(27f) + \mathrm{NH}_3(J)$$
$$\to \mathrm{Xe}^+ + \mathrm{NH}_3(J-1) + e^-, \quad J > 7; \tag{4a}$$
$$\to \mathrm{Xe}(n'l') + \mathrm{NH}_3(J-1), \quad J \leq 7. \tag{4b}$$

These processes are illustrated in Fig. 1. The arrows have lengths $2BJ$ with $J = 1, 2, 3, \ldots$) and widths which are proportional to the room-temperature populations of the upper rotational levels involved.

Collisions of the type (4b) lead to the production of seven discrete Rydberg states (or groups of states) which can be separately detected and identified by the technique of field ionization[3] since each of them will have its own characteristic critical field. This paper describes the first investigation of collisions of this type.

The apparatus has been described elsewhere[4] and only a few details will be discussed here. Xe $^3P_0$ atoms produced by electron impact excitation are excited, using a pulsed laser, to a selected Rydberg nf state in a region into which NH_3 target gas can be admitted. Approximately 7 μsec after each laser pulse, the electric field in the excitation region is increased from 0 to 1100 V/cm in ~2 μsec. As the field strength grows the different groups of Rydberg states present are successively ionized, and the resulting

