diluted with helium in a ratio 1:100 and higher, the character of the generation depends weakly on the amount of helium. Thus, the collisions between the BCl_3 molecules and air or helium molecules decrease in a suitable manner the relaxation time of the absorption in the cell.

The best results however were obtained by adding to the BCl_3 very small amounts of ammonia (NH_3). In the pressure range 1 - 7 Torr and in the presence of traces of ammonia in the cell, stable and perfectly regular generation of giant pulses is observed, with repetition periods from 40 to 100 μsec, depending on the pressure, and with an average power 1 - 2 W, equal to the continuous-generation power when the cell is evacuated. Unlike in [1], there is no continuous-generation background in the pulsed regime. When the pressure increases, the pulse repetition frequency decreases, apparently by virtue of the decrease in the losses introduced into the resonator. A pressure of 50 Torr corresponds to a repetition period of 2 msec, but in this case the generation amplitude is much smaller than at ~1 Torr. Filling the cell with pure NH_3 produced no Q-switching.

6. In interpreting the results we must bear in mind that ammonia has a strong absorption line at 950 cm^{-1} [3] and has, unlike BCl_3, a very small relaxation time.

In a cell containing BCl_3 with a slight admixture of NH_3, there occurs apparently a highly effective resonant transfer of vibrational energy by collision of the excited BCl_3 molecules with the unexcited NH_3 molecules, followed by rapid relaxation of the NH_3 molecules to the ground state. The populations of the ammonia molecules in the 950 cm^{-1} transition are not equalized by the laser emission field, since the relaxation time of the NH_3 is very short.

Obviously, in order for such a filter to operate efficiently, the number of NH_3 molecules should be much smaller than that of the BCl_3 molecules, in order that the absorption due to the ammonia not exceed the absorption due to the boron trichloride.

We note that absorption saturation in ammonia can be expected in the case of high-power lasers.

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MASS DIFFERENCE OF K_S^0 AND K_L^0 MESONS, MASS OF W MESON, AND SUM RULES FOR SPECTRAL FUNCTIONS

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A number of recent papers [1-3] have demonstrated the interesting possibility of obtaining convergent results in the calculation of the matrix elements of electromagnetic and weak second-order processes with the aid of sum rules of the spectral functions of the current propagators [4-6].

In particular, Glashow, Schnitzer, and Weinberg [2] obtained a finite expression for the matrix element of the $K_{2\pi}$ decay in second order of weak-interaction theory with intermediate vector boson. Comparison of the result with experiment yielded an estimate for the mass M of the W meson ($M \approx 7.6$ GeV).

Using a similar approach, we obtain in this paper a finite expression for the mass difference of the K_S^0 and K_L^0 mesons, $\Delta m = m_S - m_L$ (which is a fourth-order effect in the weak interaction with the W meson) and obtain for the same value of M

$$\tau_S \Delta m \approx -0.60, \tag{1}$$

which agrees well with experiment ($\tau_S \Delta m_{exp} = -(0.48 \pm 0.02)$).

The $K^0 \to \bar{K}^0$ transition amplitude $A_{K\to\bar{K}}(p_1, p_2)$, which we need here to determine Δm, is defined by the following expression:

$$(2\pi)^4 \delta^{(4)}(p_1 - p_2) A_{K\to\bar{K}}(p_1, p_2) = = -\frac{i}{4!} \int \prod_{i=1}^{4} d^4 x_i \langle \bar{K}^0(p_2) | T\{H_w(x_1) H_w(x_2) H_w(x_3) H_w(x_4)\} | \bar{K}^0(p_1) \rangle, \tag{2}$$

where

$$H_w(x) = g W_\mu(x)(J_{\mu\,2}^{\,1}(x) \cos\theta + J_{\mu\,3}^{\,1}(x) \sin\theta) + \text{h. c.}$$
$$J_\mu(x) = V_\mu(x) - A_\mu(x). \tag{3}$$

To calculate $A_{K\to\bar{K}}(p_1, p_2)$ we shall use the "light" boson approximation ($p_1, p_2 \to 0$) and the well known commutation of the currents $V_\mu(x)$ and $A_\mu(x)$ of the algebra of $SU(3) \otimes SU(3)$ symmetry. Unlike in [2], there arise here "two-current" Green's functions containing a T-product of four currents. Retaining in these Green's functions only the disconnected diagrams (with respect to the currents, see [4]), and using spectral representations [2,6] for the resultant current propagators, we obtain after simple transformations for $A_{K\to\bar{K}}(0, 0)$:

$$A_{K\to\bar{K}}(0,0) = \frac{ig^4 \sin^2\theta \cos^2\theta}{(2\pi)^4 F_K^2} \{ \int d^4 q [\frac{1}{M^4}(\int \mu^{-2} d\mu^2 \sigma(\mu^2) - F_\pi^2 + F_K^2 + F_\kappa^2)^2 + \frac{3}{(q^2 - M^2)^2}(\int \frac{d\mu^2 \sigma(\mu^2)}{(q^2 - \mu^2)})^2]\}, \tag{4}$$

where $\sigma(\mu^2) \equiv \rho_V(\mu^2) + \rho_A(\mu^2) - \rho_V'(\mu^2) - \rho_A'(\mu^2)$ and $\rho_i(\rho_i')$ are the corresponding spectral functions of the propagators of the non-strange (strange) currents ($i = V, A$). We see that, unlike [2], the first spectral sum rule suffices to eliminate the divergence in (4), so that we get ultimately, integrating with respect to q in (4),

$$A_{K\to\bar{K}}(0,0) = -\frac{3g^4 \sin^2\theta \cos^2\theta}{16\pi^2 F_K^2} \int \frac{d\mu_1^2 \, d\mu_2^2}{\mu_1^2 - \mu_2^2} \sigma(\mu_1^2) \sigma(\mu_2^2) \times$$

$$\times\left[\frac{\mu_2^2 - \mu_1^2}{(M^2 - \mu_1^2)(M^2 - \mu_2^2)} + \frac{2\mu_1^2}{(M^2 - \mu_1^2)^2}\ln\frac{M^2}{\mu_1^2}\right]. \quad (5)$$

Now, saturating the spectral functions $\rho_V(\mu^2)$, $\rho'_V(\mu^2)$, $\rho_A(\mu^2)$, and $\rho'_A(\mu^2)$ with the single-particle states $\rho(770)$, $K^*(890)$, $A_1(1080)$, and $K_A(1320)$, and taking as usual the common coefficient $2m_\rho^2 F_\pi^2$ for the corresponding δ-functions in $\rho_i(\mu^2)$, as well as the experimentally known values $F_K = 1.28F_\pi = 220$ MeV and $\cos\theta \sin\theta = 0.22$, we arrive at the result (1).

We emphasize that in our calculation we do not encounter the problem of the so-called σ terms, since it can be readily seen that they do not appear in our case.

As is well known, the unsatisfactory result obtained when the electromagnetic mass difference of the K^+ and K^0 mesons is calculated by an analogous method, is attributed just to the unaccounted-for σ terms.

It is of interest to trace in (5) the transition to the local limit of weak interaction ($M^2 \to \infty$, $\sqrt{2}g^2/M^2 = 10^{-5}/M_p^2 = G$). In this case it can be readily seen that the matrix element $A_{K\to\bar{K}}$ diverges like ln M. It becomes finite if we use the second ("less reliable") sum rule $\int\sigma(\mu^2)d\mu^2 = 0$, which played an important role in [2] for the elimination of the divergences from the matrix element of the $K_S \to 2\pi$ decay with finite M. In our case the use in (5) of the condition $\int\sigma(\mu^2)d\mu^2 = 0$ with $M \to \infty$ and fixed G leads to the finite value $\tau_S\Delta m \approx -0.36$.

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EXCITATION OF MOSSBAUER LEVELS WHEN HEAVY IONS ARE DECELERATED

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As is well known, short-period nuclei produced upon bombardment of different targets in accelerators as a result of nuclear reactions (d, p), (n, γ), etc. or Coulomb excitation, have recently come into use as sources of emission in Mossbauer transitions, in addition to the long-period parent isotopes.

We wish to call attention here to the possibility of direct or cascade excitation of nuclear levels in the accelerated nuclei themselves (beams of heavy ions) as a result of their Coulomb excitation as they are decelerated in the medium. Such a manner of exciting the Mossbauer levels would afford unique possibilities for a direct investigation of the mechanism of stopping and relaxation of the heavy ions penetrating into matter, their chem-