width $\Delta\omega \simeq 3 \times 10^4$ Hz on lines with $\lambda' \gg 1$ μ. $|\kappa_\omega|$ increases sharply in the resonance region, but it can be shown that the condition for "dragging" of the atoms coincides with the condition for strong saturation of κ_ω by the field. The strong saturation leads to a broadening of the homogeneous line [4], making it difficult for the effect to occur in the resonance region.

5. The narrow component can be observed in both absorption and emission. By scanning the frequency ω' of the monochromatic wave, directed strictly along the intense standing wave, it is possible to obtain the shape of the absorption line. If the absorption line consists of a number of lines, then the narrow component appears at the center of each line. The resolution of such a method is $R = (\omega'/\Delta\omega) \simeq 10^9 - 10^{11}$. In spontaneous or stimulated emission of atoms in the direction of the standing wave there should also arise one or several narrow components corresponding to the line structure.

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PROTON HALO AND ANTIBOUND STATES

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A hypothesis was recently advanced [1] that the proton has an aureole (halo) with dimensions $R_H \sim 5 - 8$ F and with small charge density, leading to an increase of the mean-square electromagnetic radius of the proton by an approximate factor 1.5 compared with the presently assumed value $R_p \simeq 0.8$ F.

This assumption eliminates the discrepancy between the theoretical and exponential data on the Lamb shift in hydrogen and deuterium atoms, and also reconciles the data obtained on the charge distribution in the Bi^{209} nucleus by two different methods, electron scattering and the μ-mesic atom. It does not contradict data on other processes, particularly e-p scattering.

On the other hand, as noted by the authors of [1], it is difficult to explain the existence of the halo without introducing a new light vector particle with mass ≈100 MeV, which, as is well known, has not been observed as yet.

In this note we show that this difficulty can be circumvented by postulating the existence of a virtual level ("antibound state") in the p-wave of the system of two virtual π mesons at $t = t_0$ $(0 < t_0 < 4\mu^2)$. This level should correspond to a zero of the S-matrix of the ππ scattering:

$$S(t) = 1 + 2i\rho T(t) \qquad (1)$$

($\rho = \sqrt{(t - 4\mu^2)/t}$ is a normalization factor), and consequently a pole on the second sheet of the Riemann surface of the $\pi\pi$-scattering amplitude T

$$T^{II}(t) = T^{I}(t)S^{-1}(t). \tag{2}$$

As will be shown below, this pole becomes effectively manifest in the form factors of the π meson and nucleon at $t < 0$, just as if a particle with mass $\sqrt{t_0}$ were to exist.

We consider first the form factor Φ of the π meson. Since

$$\Phi^{II}(t) = \Phi^{I}(t)S^{-1}(t), \tag{3}$$

it is obvious that Φ^{II} also has a pole at $t = t_0$. To take this pole into account we shall use a method developed in [2]. We write the Cauchy integral

$$\frac{1}{2\pi i}\oint \frac{\Phi(t')}{t'-t}dt', \quad \frac{1}{2\pi i}\oint \frac{\Phi(t')}{\rho^3(t')(t'-t)}dt' \tag{4}$$

along a contour enclosing both sheets of the Riemann surface of the function. This makes it possible to exclude from the dispersion relations the contribution of the elastic cut $4\mu^2 < t < 16\mu^2$ and replace it by the contribution from sheet II. Combining both expressions, we get

$$\Phi^{I}(t) = R_{-}(t,t_0)\frac{\lambda}{t_0 - t} + \frac{1}{\pi}\int_{16\mu^2}^{\infty}\frac{\mathrm{Im}\,\Phi(t')}{t'-t}dt' + \Delta, \tag{5}$$

where

$$\Delta = \frac{1}{2\pi i}\int_{16\mu^2}^{\infty}\{R_{-}(t,t'_{+})[\Phi^{I}(t'_{+}) - \Phi^{II}(t'_{+})] + R_{+}(t,t'_{+})[\Phi^{I}(t'_{-}) -$$

$$- \Phi^{II}(t'_{-})]\}\frac{dt'}{t'-t} - \frac{1}{2\pi i}\oint_{-\infty}\{R_{-}(t,t'_{+})\Phi^{II}(t'_{+}) + R_{+}(t,t'_{+})\Phi^{II}(t'_{-}) -$$

$$-\frac{1}{2}\frac{\rho^3(t)}{\rho^3(t'_{+})}(\Phi^{I}(t'_{+}) + \Phi^{I}(t'_{-}))\}\frac{dt'}{t'-t}, \tag{6}$$

$$R_{\pm}(t,t') = \frac{1}{2}[\pm 1 + \frac{\rho^3(t)}{\rho^3(t')}]. \tag{7}$$

The first term on the right side of (5) takes into account the contribution of the pole on sheet II (λ - residue at the pole). When $t < 0$, this term becomes complex. Since Φ^{I} is a real quantity when $t < 0$, this means that the imaginary part of the pole term should cancel completely the imaginary part coming from Δ and, the effective contribution of the antibound state is

$$\Phi_H = \frac{\lambda}{2(t - t_0)}, \tag{8}$$

i.e., it is similar to the contribution from the real level.

We can consider similarly the form factor F of the nucleon, which by virtue of the relation

$$F^{II}(t) = F^{I}(t) - 2i\rho G\Phi^{II}(t) \tag{9}$$

(G - amplitude of $\pi\pi \to N\bar{N}$) also has at $t = t_0$ a pole on sheet II, which gives at $t < 0$ a contribution in the form

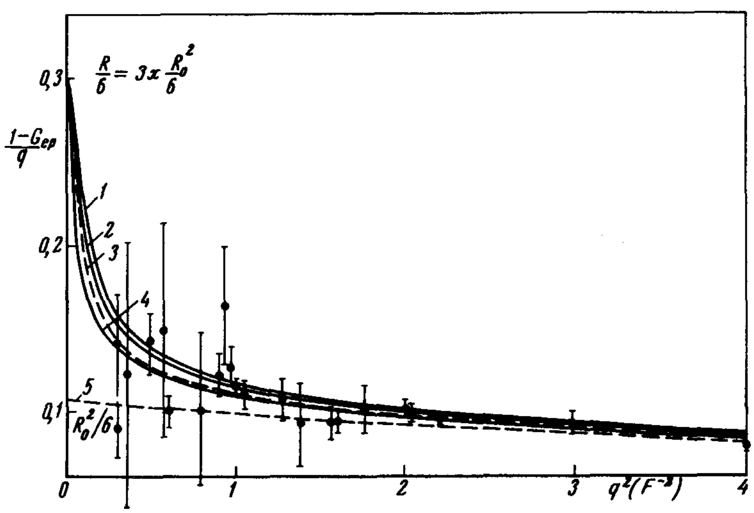
$$F_H = \frac{\lambda_N}{2(t - t_0)}. \tag{10}$$

When t_0 is close to zero and λ_N is small, the contribution F_H is significant only at small momentum transfers $q^2 \approx 0$. Denoting all the remaining contributions to the dispersion relations by F_0 (F_0, in particular, includes the contribution of the ρ meson, which is usually used to approximate F), we find that when $q^2 \gtrsim 1\ F^{-2}$ we have $F \simeq F_0$. It is not our task to discuss the different models for F_0. For convenience in the comparison, we shall use for F_0 the expression given in [1]:

$$F_0 = (1 - \epsilon)(1 + q^2 \frac{R_0^2}{12})^{-2}, \tag{11}$$

where R_0 is the radius of the "body" of the nucleon, i.e., that part which contains practically the entire charge of the proton ($R_0 \simeq 0.8$ F), and ϵ characterizes diffraction of the charge contained in the halo.

Expression (10) describes the form factor of the halo, and the radius R_H of the halo and ϵ are connected with λ_N and t_0 by the relation:



Curves for the function $(1 - F)q^{-2}$:
1 - $\epsilon = 0.06$, $R_H = 4.75$ F;
2 - $\epsilon = 0.02$, $R_H = 7.75$ F;
3 - $\epsilon = 0.0128$, $R_H = 10$ F;
4 - $\epsilon = 0.013$, $R_H = 10$ F;
5 - $\epsilon = 0$.

$$R_H = \frac{6}{t_0}, \quad \epsilon = -\frac{\lambda_N}{2t_0}.$$

t_0 and λ_N are free parameters. The figure shows plots of the function $(1 - F)q^2$, constructed with the aid of (10) and (11) for certain values of the parameters λ_N and t_0. For comparison, the same figure shows the experimental data obtained from e-p scattering for the quantity $(1 - F)q^2$ and the curves (dashed) with the aid of which these data were approximated in [1].

Thus, the existence of the virtual level in the p-wave of a two-pion system can lead to the occurrence of a proton halo (and just as well to a pion halo). We note that this level, if it does exist, should become manifest also in the processes $\pi\pi \to \pi\pi$, $\gamma\pi \to \pi\pi$, $\pi N \to \pi N$, $\gamma N \to \pi N$, etc., a study of which, in principle, could yield information concerning this level in addition to data on e-p scattering.

In conclusion we note that virtual levels can exist in principle in a three-pion system, leading to the existence of a halo of the isoscalar part of the nucleon.

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PHOTOCURRENT IN A SEMICONDUCTOR WITHOUT A SYMMETRY CENTER

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We report in this note that, unlike usual photoconductivity, a photocurrent is produced in a semiconductor without a symmetry center also in the absence of a constant field. Indeed, in a crystal without a symmetry center, a connection is possible between the current and the electromagnetic field; this connection is quadratic in the field and is described by a cross-conductivity tensor $\sigma_{\mu\alpha\beta}(\omega_1, \omega_2)$. Then the photocurrent is given by the tensor [1]

$$\sigma_{\mu\alpha\beta}(\omega, -\omega) = \frac{e}{2\hbar}\sum_{nk} f_n(k)\nabla_k^\mu \chi_{\beta\alpha}^{nk}(\omega). \tag{1}$$

Here μ, α, and β - projections on the x, y, and z axes, e and h - electron charge and Planck's constant, $f_n(\vec{k})$ - the Fermi function normalized to the total number of electrons, n and $\hbar k$ - number of the band and the quasimomentum, $\chi_{\beta\alpha}^{nk}(\omega)$ - linear susceptibility of the Bloch electron in the state $|n\vec{k}\rangle = u_{nk}\exp(i\vec{k}\vec{r})$ with energy $\epsilon_n(\vec{k})$, and ω - frequency.

The nonlinear mechanism leading to the aforementioned tensor is the combined motion of the electron, namely the interband-intraband motion [1]. Since $f_n(\vec{k})$ is even, the tensor (1) vanishes when $\alpha = \beta$ as $\vec{k} \to -\vec{k}$, regardless of the crystal symmetry type.