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Limit on an exotic three-body decay of orthopositronium

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Abstract. – An exotic annihilation of e^+e^- at rest into a single γ and a weakly interacting neutral vector boson V^0 was searched for by measuring the related γ energy under the condition that no additional γ was detected by a hermetic detector. No evidence of such an annihilation was observed, resulting in the upper limits on the branching ratio of 4.2×10^{-6} ($m_{V^0} < 200\,\mathrm{keV}/c^2$) and 2.1×10^{-5} ($m_{V^0} < 2m_e = 1022\,\mathrm{keV}/c^2$). These upper limits were converted to those on the branching ratio of an exotic three-body decay ($\gamma\gamma V^0$) of orthopositronium (o-Ps), *i.e.* 6.2×10^{-6} ($m_{V^0} < 200\,\mathrm{keV}/c^2$) and 1.7×10^{-5} ($m_{V^0} < 2m_e$). As a result, this mode was definitely excluded as the source of the reported discrepancy in the o-Ps decay rate.

The decay rate of orthopositronium (o-Ps) measured in the vacuum at an accuracy level of 200 p.p.m. [1] was found to be 1400 p.p.m. higher than that predicted by QED theory [2], with precision experiments performed in low-pressure gases [3], vacuum cavities [4], and SiO_2 powders [5] showing a consistent discrepancy. As QED calculations only include corrections up to order α, it follows that higher-order corrections may reduce the apparent difference in the o-Ps lifetime. However, several estimations of such corrections have only accounted for 500 p.p.m. [6], being significantly less than the necessary level. Another possibility exists that o-Ps has a novel decay channel which could be proven by finding a corresponding branching ratio. In this respect, we consider o-Ps to be an attractive and promising hunting field to search for exotic phenomena, although cosmological considerations place an indirect but stringent constraint on the exotic decay of o-Ps [7].

Various exotic o-Ps decay modes have been investigated to explain this discrepancy, though none have thus far been found [8]-[13]. Although all the two-body exotic decay modes [8]-[10], the 4γ mode [11] and an invisible decay mode [12], [13] have experimentally been excluded, no experiment has been carried out to search for o-Ps decay into $\gamma\gamma V^0$, *i.e.* two photons and a weakly interacting neutral vector boson V^0 (fig. 1 *a*)). Because V^0 is parity-odd for charge-conjugation, this three-body final-state decay would be the lowest-order contribution. However, this mode is difficult to search experimentally, since the γ-rays from this three-body decay would not show any peak. V^0, if it exists, would also be emitted in a prompt e^+e^- annihilation at rest into γV^0 (fig. 1 *b*)), which can be detected as the monochromatic γ-ray. This *letter* reports on our search for $e^+e^- \to \gamma V^0$, whose upper limits can be converted to

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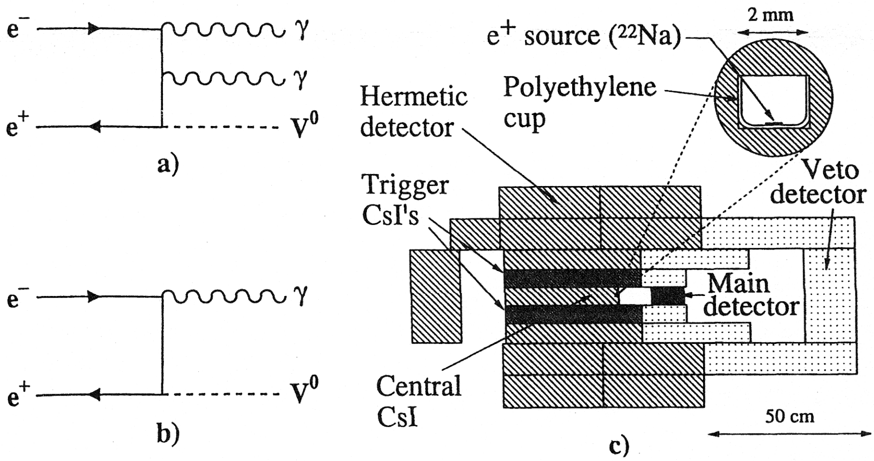

Fig. 1. – *a*) Exotic three-body decay of o-Ps. *b*) Exotic annihilation of e^+e^- at rest. *c*) Cross-sectional diagram of the employed experimental set-up.

those of the hypothesized o-Ps decay o-Ps$\rightarrow \gamma\gamma V^0$ in order to determine whether or not this mode is the origin of the reported discrepancy in the o-Ps lifetime.

Figure 1 *c*) shows a cross-sectional diagram of the experimental set-up, where as indicated in the inset, a 0.4 μCi ^{22}Na positron source with a 0.5 mm spot diameter is deposited on the inner surface and centre of a small polyethylene cup having a 2.0 mm diameter, 1.5 mm depth, and 70 μm thickness. Most emitted positrons are stopped inside the cup and promptly annihilate into two photons or possibly a single photon and a V^0. In both cases the energy of one photon will be measured by a CsI scintillator ($56 \times 56 \times 100$ mm), called the "main detector," which is situated directly in front and adjacent to the source. This detector is surrounded by 8 CsI and 30 NaI scintillation counters (weight: 15.8 and 240 kg, respectively) that detect the background and Compton-scattered photons inside the main detector. Together they form the "veto detector." The source and polyethylene cup are embedded in another CsI scintillator ($56 \times 56 \times 350$ mm), called the "central CsI," which is an active photon counter used to detect the second photon from the two-photon annihilation. Upon each β^+ decay of the ^{22}Na source, a 1275 keV photon is simultaneously emitted, being detected by CsI scintillators ($56 \times 56 \times 420$ mm $\times$ 8 modules) placed around the central CsI, called "trigger CsI's." This photon signal is used to tag each β^+ decay. The central CsI and trigger CsI's are surrounded by 16 CsI and 64 NaI counters (weight: 105 and 635 kg, respectively), forming a thick hermetic photon detector, called the "hermetic detector," which together with the central CsI detects the second annihilation photon.

Data acquisition is initiated when a pulse from the main detector and one from the trigger CsI's have coincidence within $\pm$ 100 ns. The time difference of these two pulses is recorded by a time-to-digital converter (TDC) and their energy information is recorded by analog-to-digital converters (ADCs). The light pulses from the CsI and NaI scintillation counters of the veto and hermetic detectors and those of the central CsI are also fed to the ADCs. To completely detect these pulses, the gates of the ADCs are set to open 100 ns prior to receiving pulses from the annihilation photons and to close 3 or 1 μs later for the CsI and NaI counters, respectively. The energy calibration of the main detector and the measurement of its absolute peak efficiencies are performed using the peaks obtained from γ-ray sources of known strength placed at the source position, *i.e.* ^{109}Cd, ^{22}Na, and ^{152}Eu. The resultant energy resolution is 39.8, 45.8, and 75.0 keV FWHM at 344, 511, and 1275 keV, respectively. The energy calibration of all other photon counters is similarly performed, though sources of ^{109}Cd, ^{22}Na, and ^{60}Co are used.

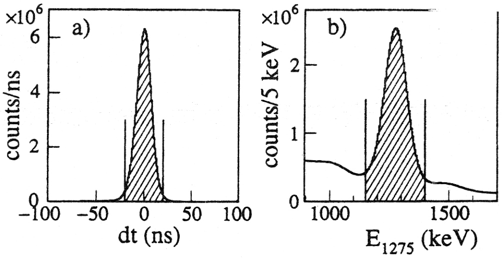

Fig. 2. – *a*) Time difference between the pulses from the main detector and the trigger CsI's. *b*) γ energy spectrum deposited at the trigger CsI's.

The energy gain of the counters and the pedestal of the ADCs are, respectively, monitored by the 511 keV peak from ^{22}Na and the zero-energy peak. Both peaks are stable to ± 0.5 keV throughout the total data acquisition period of 4.9×10^6 s. To reduce events triggered by background photons, the only events that are recorded are those in which the total energy of the veto detector is less than 5 keV. Off-line analysis of 1.3×10^8 recorded events is possible.

To eliminate the background contribution from accidental coincidence, we set a time window of 20 ns on the absolute time difference between the pulses from the main detector and the trigger CsI's (fig. 2 *a*)). The background contribution was also eliminated by setting an energy window of 1150 keV $< E_{1275} <$ 1400 keV on the energy spectrum of the trigger CsI's (fig. 2 *b*)). A total of 6.7×10^7 events remain after these cuts. The energy spectrum of these events is shown in fig. 3 *a*) as "spectrum A." The dominant peak at 511 keV is due to the two-photon annihilation. Spectrum B (fig. 3 *a*)) is obtained by setting an energy window of less than 15 and 5 keV for the central CsI and hermetic detector scintillation counters, respectively. The shape of the overall curve is due to the background photons which pass through the veto and hermetic detectors with no interaction and deposit their energy in the main detector and trigger CsI's by Compton scattering.

The following method is used to search for exotic peaks in the spectrum-B energy region from 350 to 511 keV . The region is first scanned at 5 keV intervals (inset of fig. 3 *a*)) and each scan point is considered to be the position of a peak. The bin of the energy spectrum is then set to be the standard deviation (σ) of the assumed peak shape (dotted line, inset of fig. 3 *a*)), with the centre of the bin being aligned with the scanned point. An exotic peak is then searched for by checking to see if the spectrum is consistent with an overall smooth curve plus a peak. The smooth curve is obtained at each peak position (solid line, inset of fig. 3 *a*)) by fitting the spectrum with polynomials of up to fifth order, while excluding the data within 5σ (5 bins) of the assumed peak position. The shape of the peak is assumed to be Gaussian with the width determined by the energy resolution of the main detector. The 511 keV peak of spectrum A is also fitted to such a Gaussian in order to check the validity of the assumed peak shape and to obtain the total number of e^+e^- annihilation events, $N_{e^+e^-}$, detected in our experiment. The peak is well fitted to the Gaussian with a width measured before, *i.e.* $\sigma = 19.5$ keV (FWHM $= 45.8$ keV), which is used when searching for the exotic peak at 511 keV in spectrum B. From this, the population of the peak in spectrum A is obtained as $N_{e^+e^-} = 4.1 \times 10^7$. Results indicate that no statistically significant peak was detected in spectrum B. The upper limits on the peak strengths thus obtained must be corrected for the estimated loss of genuine events, which occurs when source positrons escape from the polyethylene cup and deposit more than 15 keV of energy in the central CsI, *i.e.* they are outside its energy window. This correction is made by multiplying the upper limits by 1/0.30, where 0.30 is the probability that the positrons do not escape from the cup. We determined the probability value using a Monte Carlo simulation based on the geometries of the cup and central CsI. As some of the escaped positrons will fall within the central CsI's energy window, the non-escape probability of 0.30 is

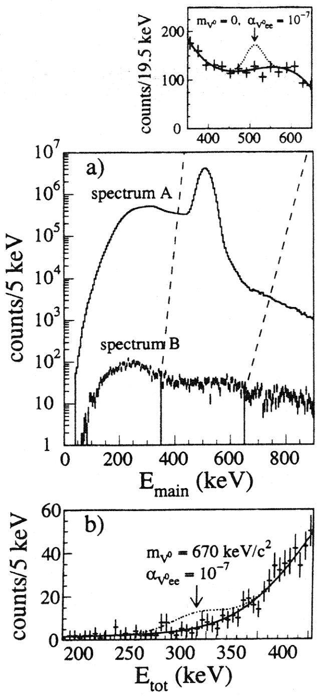

Fig. 3.

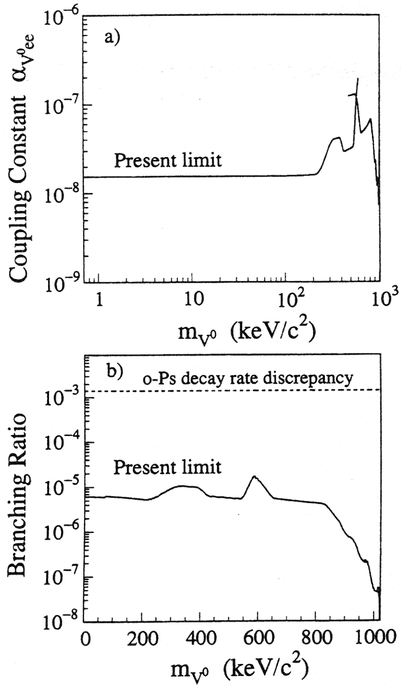

Fig. 4.

Fig. 3. – *a*) Energy spectrum of the annihilation photon measured by the main detector. Spectrum A: No constraint on the energy of the central CsI and hermetic detector. Spectrum B: The energy windows of the central CsI and hermetic detector scintillation counters are set to less than 15 and 5 keV , respectively. Inset: Magnified view of spectrum A around 511 keV (see text). The dotted line indicates the peak from $e^+e^- \to \gamma V^0$ with $m_{V^0} = 0$ and $\alpha_{V^0 ee} = 10^{-7}$: a signal that is 6.7 times stronger than the upper limit at a 90% c.l. *b*) Total energy spectrum of annihilation photons measured by hermetic detector in our previous experiment [12]. The dotted line indicates the peak from $e^+e^- \to \gamma V^0$ with $m_{V^0} = 670$ keV and $\alpha_{V^0 ee} = 10^{-7}$: a signal that is 2 times stronger than the upper limit at a 90% c.l.

Fig. 4. – *a*) Resultant upper limits at a 90% c.l. on the electron-coupling constant of the V^0-boson. *b*) Resultant upper limits at a 90% c.l. on the branching ratio of the decay o-Ps $\to \gamma\gamma V^0$.

a lower limit and corresponds to conservative upper limits on the peak strengths. The upper limits of γV^0 events with peak energy k, $n_{\gamma V^0}(k)$, can be converted into corresponding upper limits on the branching ratio, $B_{\gamma V^0}(k)$, *i.e.*

$$n_{\gamma V^0}(k) = N_{e^+e^-} B_{\gamma V^0}(k) \epsilon_{\text{pile}} \epsilon(k) / 2\epsilon(511), \tag{1}$$

where $\epsilon(k)$ is the absolute peak efficiency of the main detector for a γ-ray with energy k. It should be noted that the resultant $B_{\gamma V^0}(k)$ is only dependent on the ratio of $\epsilon(k)/\epsilon(511)$, and not on their absolute values. Analysis of the randomly triggered events under real data acquisition conditions showed that background and electronic noise pile up with genuine γV^0 events, causing an inefficiency of 20%; thus, this necessitates including $\epsilon_{\text{pile}} = 0.80$ in eq. (1).

To search for exotic single-photon annihilation events in which the photon energy is below 350 keV , we use our data obtained by a previous experiment [13] that searched for an invisible decay of o-Ps and an invisible final state from e^+e^- annihilation at rest by measuring the

TABLE I. – *90% c.l. upper limits on the coupling constant and branching ratio of the decay* o-Ps $\to \gamma\gamma V^0$.

m_{V^0} (keV/c^2)	Upper limits	
	α_{V^0ee}	B(o-Ps$\to \gamma\gamma V^0$)
< 200	1.5×10^{-8}	6.2×10^{-6}
580	1.2×10^{-7}	1.7×10^{-5}
1000	1×10^{-8}	5.3×10^{-8}
1020	1×10^{-7}	2.9×10^{-8}
$2m_e$	—	2.8×10^{-6} (a)

(a) From ref. [13].

total energy spectrum of photons E_{tot}. Figure 3 *b*) shows the E_{tot} spectrum for 4.4×10^6 e^+e^- annihilation events. We searched this spectrum for exotic peaks using the same method used for spectrum B, though the bin of the E_{tot} spectrum is always set to be 5 keV . Once again, however, no statistically significant peak was detected. The resultant upper limits of γV^0 events were also similarly converted into corresponding upper limits on the branching ratio.

The upper limits on the branching ratio of γV^0 annihilation $B(e^+e^- \to \gamma V^0)$ were subsequently obtained over the entire m_{V^0} region below $2m_e = 1022\,\text{keV}/c^2$, where the mass of the V^0-boson, m_{V^0}, is related to the photon energy by $E_\gamma = m_e(1-(m_{V^0}/2m_e)^2)$. At a 90% confidence level they are 4.2×10^{-6} ($m_{V^0} < 200\,\text{keV}/c^2$) and 2.1×10^{-5} ($m_{V^0} < 2m_e$), which can easily be converted to upper limits on the coupling constant of the V^0-boson to the electron, α_{V^0ee}, *i.e.*

$$\alpha_{V^0ee} = \frac{\alpha}{2}\left(1-\left(\frac{m_{V^0}}{2m_e}\right)^2\right)^{-1} B(e^+e^- \to \gamma V^0), \tag{2}$$

where $\alpha = 1/137$. Table I and fig. 4 *a*) show these upper limits for $m_{V^0} < 1020\,\text{keV}/c^2$. Conversion of the upper limits on α_{V^0ee} to those on the o-Ps branching ratio of $\gamma\gamma V^0$ decay $B(\text{o-Ps} \to \gamma\gamma V^0)$ is also possible using the lowest-order Feynman graphs of this three-body decay, *i.e.*

$$B(\text{o-Ps} \to \gamma\gamma V^0) = \frac{3\alpha_{V^0ee}}{\alpha(\pi^2-9)} \int_0^{1-r} dx_1 \int_{1-r-x_1}^{1-r/(1-x_1)} dx_2 \left[\frac{x_0^2}{(x_0-2r)^2} \times \right.$$
$$\left. \times \left\{\left(\frac{1-x_0}{x_1x_2}\right)^2 + \left(\frac{1-x_1}{x_2x_0}\right)^2 + \left(\frac{1-x_2}{x_0x_1}\right)^2 + \frac{F}{x_0^2x_1^2x_2^2}\right\}\right]_{x_0=2-x_1-x_2}, \tag{3}$$

$$F = 2r^4 + 2r^3(x_1x_2+4x_1+4x_2-2) + 2r^2(x_1^2x_2+x_1x_2^2+6x_1^2+6x_1x_2+6x_2^2-8x_1-8x_2+3) +$$
$$+ r(x_1^2x_2^2+8x_1^3+10x_1^2x_2+10x_1x_2^2+8x_2^3-20x_1^2-28x_1x_2-20x_2^2+20x_1+20x_2-8), \tag{4}$$

where $r = (m_{V^0}/2m_e)^2$. Figure 4 *b*) and table I show the resultant upper limits on $B(\text{o-Ps} \to \gamma\gamma V^0)$. Note that the results for $m_{V^0} \ll m_e$ agree with those obtained using the asymptotic form of the eq. (3): $B(\text{o-Ps} \to \gamma\gamma V^0) = 3\alpha_{V^0ee}/\alpha$. The upper limits on $B(\text{o-Ps} \to \gamma\gamma V^0)$ do not even come close to explaining the reported discrepancy in the o-Ps decay rate. If the V^0-boson has a mass between $1020\,\text{keV}/c^2$ and $2m_e$, then the $\gamma\gamma V^0$ decay of o-Ps should appear as an invisible decay of o-Ps due to the total energy of the two photons in the final state being less than 2 keV. However, since this decay was not observed in the previously examined signal region of invisible decay ($E_{tot} < 50\,\mathrm{keV}$) [13], the upper limit must be $B(\text{o-Ps} \rightarrow \gamma\gamma V^0) < 2.8 \times 10^{-6}$ ($1020\,\mathrm{keV}/c^2 < m_{V^0} < 2m_e$). When these limits are considered in conjunction with those in fig. 4 *b*), it becomes obvious that, regardless of the mass of the V^0-boson, the $\gamma\gamma V^0$ decay of o-Ps is not the source of the reported discrepancy in the o-Ps decay rate.

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